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STUDIES OF THERMAL ENERGY CONFINEMENT SCALING
IN PDX PLASMAS: $D^0 + H^+$ LIMITER DISCHARGES

By

S.M. Kaye et al.

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STUDIES OF THERMAL ENERGY CONFINEMENT SCALING

IN PDX PLASMAS: $D^0 \rightarrow H^+$ LIMITER DISCHARGES

S.M. Kaye, R.J. Goldston, M. Bell, K. Bol, M. Bitter R. Fonck,
B. Grek, R.J. Hawryluk, D. Johnson, R. Kaita, H. Kugel,
D. Mansfield, D. McCune, K. McGuire, D. Mueller, M. Okabayashi,
D. Owens, G. Schmidt, P. Thomas¹

Plasma Physics Laboratory, Princeton University
Princeton, New Jersey 08544

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ABSTRACT

Experiments were performed on the PDX tokamak to study plasma heating and β scaling with high power, near-perpendicular neutral beam injection. The data taken during these experiments were analyzed using a time-dependent data interpretation code (TRANSP) to study the transport and thermal confinement scaling over a wide range of plasma parameters. This study focuses on results from experiments with D^0 injection into H^+ plasmas using graphite rail limiters, $a = 40$ to 44 cm, $R = 143$ cm, $I_p = 200$ to 480 kA, $B_T = 0.7$ to 2.2 T, and typically $\bar{n}_e = 2.5$ to $4.2 \times 10^{13} \text{ cm}^{-3}$. The results of this study indicate that for both ohmic and neutral beam heated discharges the energy flow out of the plasma is dominated by anomalous electron losses, attributed to electron thermal conduction. The ion conduction losses are well described by neoclassical theory; however, the total ion loss influences the power balance significantly only at high toroidal fields and high plasma currents. Therefore, except for these cases, the total thermal energy confinement time

¹ JET Joint Undertaking, Abingdon, Oxon, U.K.

for the neutral beam heated discharges follows the scaling of electron energy confinement. While the confinement is found to have little dependence on toroidal field, it is a strong function of plasma current, increasing with increasing I_p . In contrast, ohmic confinement times are found to scale as $\bar{n}_e a^{1/2}$. Because of the dominant effect of the electrons, the total confinement is found to depend heavily on the variation of the electron thermal diffusivity, χ_e , in the outer region of the plasma ($a/2 < r < 7a/8$). The values of χ_e in this region of the plasma show an approximately inverse dependence on the local value of the poloidal magnetic field.

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1. INTRODUCTION

Energy confinement studies of neutral beam heated plasmas in PDX have been carried out over the full range of operating parameters: $200 < I_p < 480$ kA, $0.7 < B_T < 2.2$ T, $1.7 < q_{cyl} < 7$, $2.5 < \bar{n}_e \lesssim 7 \times 10^{13} \text{ cm}^{-3}$, and $0 < P_{inj} < 7.2$ MW. There were two main experimental efforts. The first was to study the effectiveness of perpendicular neutral beam heating in discharges at high toroidal field ($B_T = 2.2$ T, $I_p = 200$ to 480 kA) [1-3]. These studies were typified by high temperatures in both ions [$T_i(0) = 6.5$ keV] and electrons [$T_e(0) = 2.2$ keV]. The second was a $\langle \beta_T \rangle$ -scaling experiment carried out over a wide range in q ($B_T = 0.7$ to 1.5 T, $I_p = 250$ to 400 kA, $q_{cyl} = 1.9$ to 3.7) [4,5]. Both studies were carried out for the most part using circular cross sections of H^+ plasmas limited by graphite rails with deuterium neutral beam injection (NBI) to maximize absorbed power, and with $a = 40$ to 44, $R = 143$ cm.

1.1 PDX experiments

The high field (2.2 T) study was conducted to test the efficacy of near perpendicular neutral beam injection (NBI) in producing high temperature plasmas at low collisionality. The results of these injection experiments on PDX were compared to those of PLT, a tokamak with a comparable discharge characteristics but with a near tangential injection system and $P_{inj} \approx 2.6$ MW. It was found that the central ion heating efficiency, a measure of the central ion energy confinement, was approximately inversely proportional to the ion thermal diffusivity in the core. The central ion heating efficiency for PDX was found to increase with increasing current and to be comparable to that of PLT at comparable currents. This study also examined the variation of total (beam + plasma) stored energy with absorbed beam power (injected power less shine-through) using the measured $L_{i/2} + \beta_{pol}$ signal. For the current

range from 200 to 480 kA, the total stored energy was found to increase linearly with absorbed beam power over the full range from ~ 1 to 6 MW [1].

In the more recent $\langle \beta_T \rangle$ -scaling experiments diamagnetic $\langle \beta_T \rangle^{\text{dia}}$ values of $\sim 3.0\%$ were obtained at low toroidal fields (0.9 T) [6]. From detailed scans in power, current, and toroidal field scans, it was found that for a given toroidal field and absorbed power, $\langle \beta_T \rangle$ increased nearly linearly with plasma current [4]. Furthermore, at high toroidal fields (> 1.5 T), $\langle \beta_T \rangle$ increased linearly with injected power. The confinement time, however, showed a modest decrease with increasing power ($\propto P_b^{-1/3}$) from one to four beams. However, at lower toroidal fields (0.7-1.2 T), $\langle \beta_T \rangle$ was observed to saturate when the value of $\langle \beta_T \rangle q_{\text{CY1}}(a)$ exceeded $\sim 4.5\%$. The degree of saturation versus input power was found to be independent of current for a given toroidal field, but the degree of saturation increased as the toroidal field was lowered. Associated with the saturation of $\langle \beta_T \rangle$ were large amplitude $m=1$ modes observed on the soft X-ray detectors. At large enough amplitudes these MHD spikes, dubbed "fishbone" activity, were found to be associated with large drops in neutron emission and a depletion in the fast charge exchange efflux spectrum; modeling studies matching the charge exchange efflux depletion and the drops in the neutron rate indicated that during the most severe fishbone activity as much as 40% of the beam heating efficiency was lost [5]. Cases used for the study to be presented here were those in the "unsaturated" $\langle \beta_T \rangle$ regime, i.e., where either no fishbone activity was observed or where the effect of the fishbone activity on the beam ions was minimal, and therefore a classical slowing down treatment was sufficient to describe the beam ion behavior. This generally restricted the cases to ones in which the injected beam power was ≤ 3.5 MW for $B_T \leq 1.2$ T.

1.2 Survey of tokamak confinement during neutral beam heating

Neutral beam heating experiments with limiter plasmas have been carried out in a variety of tokamaks (ASDEX [7,8], DITE [9,10], D-III [11-13], ISX-B [14-18], JFT-2 [19,20], PLT [2,21-25], TFR [26,27], and T-11 [28]). Because reviews of transport and energy confinement scaling have recently appeared in the literature [29,30], only the gross features of the studies on the above-mentioned machines will be discussed here. A result common to most of these experiments is that the energy confinement during neutral beam injection (NBI) at densities above $3 \times 10^{13} \text{ cm}^{-3}$ and currents below 600 kA is degraded by a factor of 1.5 to 4 from that observed during ohmic discharges. This result holds for discharges with beam injection powers, P_{inj} , at least twice that of the ohmic power, $P_{inj} > 2 P_{OH}$. During some low density ($\bar{n}_e \leq 2.3 \times 10^{13} \text{ cm}^{-3}$) experiments on PLT [22,23] and TFR [27], little change in energy confinement was observed with the onset of NBI; however, the confinement was relatively poor during the ohmic phase at these low densities. At higher densities on PLT [23], the gross confinement degraded with the onset of NBI, although an improvement in central confinement was observed in some cases.

Recently, results from ASDEX [31-33], D-III [34-38], and PDX [39,40] indicate that under certain operating conditions the energy confinement during NBI may be comparable to that of ohmic plasmas. This plasma condition, commonly referred to as the H-mode, is associated with divertor operation and seems related to details of edge conditions. Furthermore, the ISX-B group has recently reported improved confinement in limiter discharges with a small amount of neon injection [41] or without titanium gettering [42]. These improved confinement modes of operation are beyond the scope of the present paper, and they will be reported in a future publication.

Limiting the discussion to the results that exhibit degraded confinement, a common reported feature of these experiments is the strong dependence of the energy confinement, τ_E , on plasma current, with τ_E varying as $I_p^{1-1.5}$ for fixed $\bar{n}_e$ [4,11,13,15,33]. However, little or no scaling of τ_E with toroidal field [4,11,15,18] or plasma density is observed. In addition, the results of many of the experiments, notably those from ASDEX, DITE, D-III, and ISX-B indicate that τ_E decreases with increasing beam power ($\tau_E \propto P_{inj}^{-(1/3-2/3)}$). For DITE, ISX-B, JFT-2, and T-11, this was viewed as resulting in a "soft" β_{pol} limit. The results from D-III indicate that plasma shaping is an important aid to confinement; τ_E is found to increase by a factor of ~ 3 as κ , the plasma elongation, was increased from 1 to 1.6 at fixed α . The improved confinement with κ was due mostly to the increased current that the shaped plasmas were able to maintain; at fixed I_p the confinement was found to vary as $\kappa^{1/2}$.

In most cases (ASDEX, ISX-B, D-III, and DITE at low densities) the degradation in τ_E with NBI was associated with a degradation in electron energy confinement. Since radiation losses in most of these experiments were only a small fraction of the total electron energy loss, the mechanism assumed to be responsible for the bulk of the energy loss was an anomalous electron thermal transport process taken to be electron thermal conduction. For instance, Becker et al. deduced that the electron thermal diffusivity, χ_e , had to increase by a factor of two to three over the ohmic value to explain the confinement degradation in NBI plasmas [8]. On the other hand, at high densities in the DITE tokamak, electron conduction losses were significant, but the ion energy losses (conduction, convection, and charge exchange) were found to be dominant.

The reason for the enhanced electron thermal diffusivity that is deduced from most NBI tokamak experiments is uncertain. One theory of anomalous electron transport was partially successful in accounting for results observed in ISX-B. It was found that the experimentally deduced χ_e in the "confinement region" of the plasma discharge, the region from $r(a = 1)$ to $r = 0.8a$, agreed approximately in magnitude with that predicted from the high- β_{pol} resistive ballooning mode calculation of Carreras and Diamond [43]. ISX-B discharges were modeled successfully with a predictive transport code using an electron thermal diffusivity equal to one-half the INTOR value ($\chi_e^{INTOR} = 5 \times 10^{17}/n_e \text{ cm}^2/\text{sec}$) plus twice the χ_e given by the Carreras-Diamond model.

The DITE group inferred from their measurements that there was no inherent difference in the physics governing the main transport loss process in OH and NBI plasmas. They deduced that the gross energy confinement time scaled simply as T_e^{-4} for both OH and NBI plasmas [10].

In this study, results of the transport analysis of zero to four beam power scans on PDX over a wide range of toroidal fields and plasma currents are reported. The thermal energy confinement is found to be degraded in NBI discharges over that in ohmic discharges; however, once degraded, τ_E shows little further decrease with increasing beam power beyond the two beam level. This is consistent, within errors, with the magnetics measurements in which the gross τ_E (beam and thermal) was found to scale weakly with injected power ($\tau_E \propto P_{inj}^{-1/3}$). In addition, τ_E is found to scale nearly linearly with I_p , but to have little explicit dependence on B_T . While the behavior of the ions can be well explained by neoclassical theory, the total thermal energy confinement in the PDX NBI discharges is found to be influenced primarily by electron confinement. The degradation of τ_E in the NBI phase and the scaling

of r_E with I_p is due to changes of the electron thermal diffusivity outside the $q = 1$ surface.

2. PDX DESCRIPTION

The PDX is an extensively diagnosed neutral beam heated tokamak which can be run in either limiter or in various divertor configurations. In this paper only the results obtained with circular cross section, limiter plasmas will be discussed. During high field, high power experiments carried out in 1981, the discharges were run with uncooled top and bottom poloidal rail graphite limiters, while during the high- $\langle\beta_T\rangle$ experiments carried out in 1982, only the top limiter (with water-cooling) was used due to technical difficulties. A more comprehensive description of the PDX tokamak is given in Meade et al. [44].

The neutral beam injection system on PDX was built by the joint Princeton Plasma Physics Laboratory - Oak Ridge National Laboratory Heating Project, and consists of four duopigatron neutral beam injectors and beamlines. The injection systems, each with a 50 keV ion source, provided up to 7.2 MW of D⁰ power and up to 5.5 MW of H⁰ power, not accounting for duct reionization losses, for pulse lengths of up to 300 msec. The four beam lines inject at a tangency major radius of 35 cm giving an angle 14° from perpendicular at the center of the vessel. The beam injection systems are routinely operated under automatic computer control providing both reliability and reproducibility of beam line operations. Beamline calorimetry with direct computer processing of temperature sensor waveforms permits an accurate energy calibration to be made within 17 sec after conditioning shots, which occur at a 60 sec repetition rate between the 5 min PDX shot repetition rate. Reionization losses in the ducts connecting the beamlines to the torus were found to be ~ 12% [45].

The extensive diagnostic capability of PDX allowed a detailed magnetic and kinetic description of the plasma. Electron density and temperature profiles were obtained from a 56-point Thomson scattering system while line-averaged electron densities were measured using both a 2-mm microwave and a far-infrared laser interferometer. The central ion temperature was measured by a perpendicular mass resolving passive thermal charge-exchange analyzer and from the Doppler broadening of the TiXXI K_{α} line by a crystal spectrometer. Estimates of impurity radiation were obtained from a 19-channel bolometer array and a multiwavelength vacuum ultraviolet spectrometer. Z_{eff} estimates were generally based on visible bremsstrahlung radiation, supplemented by measurements of the charge-exchange recombination emission from fully stripped oxygen and carbon. Two techniques to determine the volume-averaged toroidal β , $\langle\beta_T\rangle$ were used. The first technique was to measure plasma diamagnetic effects from the voltage induced in the toroidal field coils [46]. The volume-averaged, perpendicular plasma β so obtained is denoted by $\langle\beta_T^{dia}\rangle$. The second technique was to infer an "equilibrium" β value from the magnetically measured $k_i/2 + \beta_{pol}$ signal, the current assumed to have the form

$$J \propto \left[1 - \left(\frac{r}{a}\right)^2\right]^{\alpha_{cyl}/\alpha_0 - 1}; \quad \alpha_0 = 0.8, \quad \alpha_{cyl} = \frac{aB_T}{RR_p}.$$

The β deduced in this fashion is denoted by $\langle\beta_T^{eq}\rangle$. The two $\langle\beta_T\rangle$ measurements agreed to within $\sim 10\%$. The systematic discrepancy was consistent with the calculated velocity space anisotropy of the near perpendicular beam particles. Beam slowing down spectra were measured by the horizontally scanning, high energy electrostatic charge-exchange analyzer of the Fast Ion

Diagnostic Experiment (FIDE). The MHD activity was measured by both the horizontal soft X-ray array and Mirnov coils. This list of PDX diagnostics is not complete, but it represents those which contributed most heavily to the confinement scaling study presented here.

2.1 . Transport analysis code input and calculations

The data that were used for this study were analyzed by the kinetic transport analysis code TRANSP [47,48]. The TRANSP code is a 1-D time-dependent interpretation code that assumes a circular concentric cross sectional plasma and that relies heavily on experimental input. The data inputs to the code include electron density and temperature profiles, line-averaged electron density, radiated power, central ion temperatures, plasma current and position, surface voltage, Z_{eff} , neutral beam power, injection voltage, and beam geometry.

Because electron transport affects the plasma energy confinement most critically, the electron temperature and density profiles constitute two of the most crucial inputs to the transport analysis code. The profiles are obtained from a Thomson scattering system capable of measuring 56-point T_e and n_e profiles at one time during each shot. The measured profiles undergo several levels of preparation before being input into the code. Profiles at the same laser firing time from several reproducible discharges are averaged together to yield a representative profile at that time. These averaged profiles are then spatially smoothed using a sliding weighted-average filter. Because of Thomson scattering measurement uncertainties in the outer portion of the discharge, the outer regions of the profiles are further smoothed so that the T_e values ~ 2 cm outside the plasma, both on the outside and the inside, are constrained to be 30 eV. This additional smoothing

results in, at most, a 20% change in the fitted T_e values over the outer 20% of the plasma. An example of averaged (unsmoothed) T_e and n_e profiles, and of the same profiles after the smoothing process is shown in Fig. 1. Loop voltage, plasma current, and line-averaged density traces are also smoothed and averaged in a similar fashion.

A "slice and stack" method, in which T_e is assumed to be constant on a flux surface, is used to symmetrize the T_e profile. The plasma radius, r , and shift of each flux surface, relative to the nominal major radius, R_0 , is determined from the positions of equal inner and outer T_e . In most cases, this experimentally determined shift, $\Delta R(r)$, is in agreement with the theoretically calculated Shafranov shift [49] to within 10 to 15%. The density at a given plasma radius is calculated by averaging the densities at points on opposite sides of the surfaces defined by the theoretical Shafranov shift. The maximum departure of the symmetrized density profile from the original asymmetric one (e.g., Fig. 4) is 5 to 20%, n_e always being higher on the large major radius side. Although the density asymmetry is not taken into account in the code calculations, which assume circular concentric flux surfaces, its overall effect on the confinement calculation should be relatively small. In the particular case of beam deposition, the effect of the density asymmetry is approximately offset by the effect of the differing path length across a flux tube for the circular concentric plasma assumed in the code and that of the true shifted flux surface plasma.

For most of the results to be reported here, a set of Thomson scattering T_e and n_e profiles were obtained only at one time in the equilibrium phase of the discharge, near the end of the beam pulse. For these cases, the transport code was run to equilibrium to model only that discharge time. For some cases a full time evolution was obtained in which T_e and n_e

profiles every ~ 50 msec across the discharge (including ohmic times) were concatenated to form a detailed time history of the electron temperature and density. Such analyses generally required at least ten to fifteen reproducible discharges at fixed plasma and beam conditions.

The calculations in the code are divided into five major sections; magnetic diffusion, thermal neutral penetration, beam heating, particle balance, and energy balance. In the magnetic diffusion routine the poloidal magnetic diffusion equation is solved under the constraint that the calculated surface voltage match the value inferred from measurements. To do this, either the Spitzer [50] or neoclassical [51] resistivity model is adopted; Z_{eff} and consequently the resistivity are adjusted dynamically in the code to yield the correct time-varying value of the surface voltage. The Z_{eff} calculated from the magnetic diffusion routine is not necessarily used to determine the plasma composition because of its sensitivity to precise profile shapes and to inductive effects which may be spuriously introduced by concatenating Thomson scattering profiles from different times in different discharges. The calculated Z_{eff} can be used, by comparing it to the spectroscopically measured value, to make a judicious choice of the resistivity model for a particular run. The high field discharges ($B_T = 2.2$ T) were typically fit better by the neoclassical model, while at lower fields ($B_T \leq 1.5$ T) the Spitzer model fit the data better. However, the thermal confinement time is insensitive to the choice of resistivity model, except for possible systematic errors in the measured electron temperature.

Two sources of thermal neutrals are considered in the neutral penetration model, a cold source ($T_0 = 3$ eV) to represent the gas puff, and a warmer source of recycling neutrals with $T_0 = T_1(o)/30$. Thirty percent of the escaping neutrals are reflected as "warm" neutrals. The neutral density

radial profile, $n_0(r)$, is calculated using a simple but accurate neutral penetration model [52,53], while the absolute magnitude of the neutral density is determined from particle balance using an assumed edge recycling time, $\tau_p(a)$. For the results to be presented, $\tau_p(a)$ was assumed to be ~ 30 to 40 msec, consistent with the ionization source rate determined from toroidally averaged edge H_α emissions.

The beam deposition code [46] is a full Monte-Carlo calculation which follows the orbits of a constant census of 1000 particle guiding centers. Beam neutral ionization mechanisms that are included in the code are bulk electron, ion, and impurity impact ionization, and charge exchange with thermal, beam, and impurity ions [54-56]. Beam ion orbits and losses by charge exchange are modeled using Monte-Carlo techniques. Beam ion energy is deposited into the thermal population by classical binary collisions between the beam ions and either the thermal ions or electrons. Recently, finite Larmor radius effects have been added to the beam code. This effect is significant for the lower field ($B_T \lesssim 1.2$ T) cases.

The particle balance is simply the solution of the continuity equation; by knowing both the measured $\partial n / \partial t$ and the thermal and beam source terms from published reaction cross sections [54-56], the particle outflow rate is calculated. The particle outflow is directly coupled to the convective energy loss ($\propto 5/2 nTv$) in the ion and electron energy balances. A more complex approach to the particle outflow term has been used in the recent work of Strachan et al., [57]. Also included in the energy balances are the energy source terms from the beams and ohmic heating, and losses such as radiation, charge exchange, and conduction. Electron-ion coupling is assumed to be classical. The convective loss from the core for the cases to be presented is relatively insensitive to the choice of $\tau_p(a)$ since the major source of

ionization there is due to beam neutrals. Ion conduction is assumed to be some multiple of the neoclassical value, to match the fit to the thermal charge-exchange spectrum and the magnetically measured $\langle \beta_{Ti} \rangle$; for this study either the Hinton-Hazeltine [58] or Chang-Hinton [59] model was used as the base value. All the electron energy source and loss terms are known to some degree except for thermal conduction; hence, the conduction loss, from which the electron thermal diffusivities are deduced, is output. The code contains no detailed impurity model; fully stripped oxygen is assumed to be the only impurity present, with the impurity density determined from an input composition Z_{eff} that either calculated from the visible bremsstrahlung measurements or constrained to be consistent with the resistivity calculations.

As yet, no attempt has been made to model the effect of sawtooth oscillations on the distribution of energy transport in the PDX discharges. Sawteeth, with typical periods of 15 to 20 msec, were seen often in the data to be reported; however, in the worst cases the relative change in electron energy, due to sawteeth, within the $q = 1$ surface was 10 to 15% [60]. When compared to a typical central electron energy replacement time of ~ 10 msec, this is a modest effect. To offset any bias induced by these oscillations, Thomson scattering profile measurements were made at different times during the sawtooth period for several reproducible discharges. The resulting averaged T_e and n_e profiles were then taken to be representative of the nonperturbed temperature and density.

2.2 Dataset

The PDX database used for the magnetics scaling studies [3] consisted of over 1000 discharges selected on the basis of shot quality and diagnostic

coverage. Figure 2 is the distribution of data in I_p and B_T parameter space for the subset of the discharges contained in the database used in the heating and $\langle \beta \rangle$ -scaling studies [1-4]. Each circled cluster of points represents a fully documented, controlled power scan (0 to 4 beams) at fixed B_T , I_p , and, to the degree possible, $\bar{n}_e$. The data shown in the figure were taken from circular discharges limited by poloidal graphite rail limiters with $R = 140$ to 144 cm, $a = 40$ to 44 cm, and $\bar{n}_e = 2.5$ to $7 \times 10^{13} \text{ cm}^{-3}$; the highest densities occurred during four beam NRIs at the highest currents. The line-averaged density was held, if possible, at $\sim 4 \times 10^{13} \text{ cm}^{-3}$ for the lower field ($B_T \leq 1.5$ T) discharges and 2.3 to $3.4 \times 10^{13} \text{ cm}^{-3}$ at 2.2 T. For these limiter plasmas no controlled density scan was performed. Deuterium injection into hydrogen target plasmas was used in order to maximize the absorbed power. While only equilibrium values of β_T , $\langle \beta_T^{eq} \rangle$, are available for the 2.2 T points, both $\langle \beta_T^{eq} \rangle$ and $\langle \beta_T^{dia} \rangle$ are available for the remainder of the points. The family of data points is bounded on two sides; the 10% orbit loss boundary occurs near $I_p = 200$ kA, and it represents the boundary below which $> 10\%$ of the injected beam ions are lost from the plasma due to classical bad orbits, not including finite Larmor radius effects. The $q_{cyl} = 1.7$ boundary represents a limit on the machine operation below which plasma disruptions prohibited successful discharges. Within this operating regime $\langle \beta_T^{dia} \rangle$ values of $> 3.0\%$ and β_{pol}^{eq} values of up to 1.8 were achieved.

2.3 Data input and scaling

The time history of a typical two beam PDX shot is shown in Fig. 3 ($a/R = 40/143$ cm, $B_T = 1.0$ T). The current and surface voltage (loop voltage with vacuum induction corrections) are given in the top left panel. For this discharge the maximum current was ~ 300 kA with a surface voltage of just

under 1 V. The line-averaged density (upper right panel) rose to $4.2 \times 10^{13} \text{ cm}^{-3}$ from $2 \times 10^{13} \text{ cm}^{-3}$, at 400 msec just before the onset of NBI. The beam pulse lasted for 300 msec with an injected beam power of 3.2 MW. Over 90% of the injected power was absorbed by the end of the beam pulse. The value of $\langle \beta_T^{eq} \rangle$ is deduced from the magnetically measured $\ell_i/2 + \beta_{pol}$ (lower left panel) and the assumed current profile shape discussed earlier.

The equilibrium energy confinement time, τ_E^{eq} , is determined from $\langle \beta_T^{eq} \rangle$, P_{abs} , and from the time rate of change of $\langle \beta_T^{eq} \rangle$. The τ_E^{eq} (lower right) drops precipitously at the start of NBI from a value of 40 to 45 msec to a value of 16 to 19 msec. Because of the uncertainties in the $\ell_i/2$ model, the value of τ_E^{eq} during the ohmic phase when $\ell_i/2 > \langle \beta_T^{eq} \rangle$, may not accurately reflect the ohmic confinement in all cases. For this discharge, the Shafranov q , defined to be

$$q_{SHAF} = q_{cyl} \left\{ 1 + \frac{a^2}{R^2} \left[1 + \frac{1}{2} (\ell_i/2 + \beta_{pol})^2 \right] \right\}$$

decreases to a value of ~ 2 by the end of the beam pulse. After the beams are shut off, the plasma disrupts promptly, probably due to the fact that this discharge is beyond its ohmic density limit.

The time evolution of the T_e and n_e profiles for the discharge conditions discussed above is shown in Fig. 4. The profiles shown in the figure are spline fits to the measured data points. With the onset of NBI, the temperature increased from an ohmic value of $\sim 600 \text{ eV}$ to a peak value of $\sim 1 \text{ keV}$ at the end of the beam pulse. As previously discussed, the skewing of the T_e profile towards the outside is consistent with the theoretically

calculated Shafranov shift. The density profile, on the other hand, broadens considerably shortly after the onset of NBI, with a pronounced asymmetry clearly observed well into the beam pulse. This asymmetry is greater than that expected from the Shafranov shift, and it cannot simply be explained by the combined effect of toroidal rotation and poloidal asymmetries in the beam ion distribution [61-64].

The electron temperature and density profile shapes were studied over a wide range of toroidal fields, plasma currents, and beam powers. The profile peakedness (central value/volume averaged value) is plotted in Fig. 5 for discharges at 1.5 T for four different currents as a function of total power ($P_{\text{tot}} = P_{\text{abs}} + P_{\text{oh}}$, P_{abs} being the power contained in the beam particles ionized in the plasma, and P_{oh} being the ohmic input). In this figure, the volume-averaged electron temperature has been density weighted. The peakedness of the density profile decreases with the onset of NBI ($P_{\text{tot}} > 1$ MW), but shows little dependence on either current or beam power level beyond this value. The temperature profiles show very little systematic dependence on P_{tot} , but they clearly exhibit a broadening at higher currents (lower q). At fixed I_p (not shown) the central electron temperature increases almost linearly with B_T , with little change in temperature for $r > a/2$, also consistent with broader profiles at lower q .

The dependence of temperature profile peakedness on q is summarized in Fig. 6, which displays data compiled during ohmic, two- and four-beam discharges over a range of currents and toroidal fields. Despite the differences in the heating, very little difference in the peakedness of the profiles at fixed q is seen between the ohmic and NBI plasmas, while $T_e(0)/\langle T_e \rangle$ more than doubles as q_{cyl} is increased from two to eight.

Central ion temperatures were deduced primarily from the mass-resolving charge-exchange analyzers, and for the higher field discharges ($B_T = 2.2$ T) also from Doppler broadening measurements of TiXIX K_α line radiation. Because of the presence of some hydrogen in the deuterium beams, portions of the fast hydrogen slowing-down spectrum were observed in the thermal charge-exchange measurement. The transport code was used to determine the slowing-down portion of the hydrogen spectrum by assuming a pure hydrogen beam in a pure deuterium plasma, and using the geometric factors for the charge-exchange signal. The calculated ion tail fit the measured spectrum quite well, and the data were re-analyzed by subtracting the hydrogen slowing-down spectrum from the raw data. The remaining spectrum was attributed to the thermal distribution. This process lowered the fitted temperature by typically 5 to 20%.

A time evolution transport run based on the discharge shown in Figs. 3 and 4 (1.0 T, 300 kA, two beams) was made using the charge-exchange spectrum to model the central ion temperature as a function of time. The time history of the temperature deduced from the measurements, along with the spectra fit limits, were input to the code. The ion neoclassical conductivity multiplier, using the Chang-Hinton model as a basis, was then determined in a feedback loop to bring the fits to the calculated and measured charge-exchange spectra into agreement. Figure 7 displays the measured and calculated fits and the neoclassical multiplier as functions of time. The small multiplier at 400 msec, the onset of NBI, reflects the rapid increase in the measured temperature at this time and the need for additional n_e and T_e profiles. The large multiplier at $t = 0.45$ s is due to the transient change in T_i at the start of NBI, and it reflects the fact that during this low density phase of the discharge the ion temperature was expected to rise above its equilibrium

value due to the lower plasma density. Later in the discharge, near equilibrium, the value of the multiplier settles down to a value of ~ 3 . For this case, the value of the central ion temperature given by the ion power balance was 1.3 KeV at 580 msec.

3. TRANSPORT ANALYSIS RESULTS

How well the transport code does in modeling both the thermal plasma and the beam distribution may be checked by comparing the $\langle \beta_T^{dia} \rangle$ that is calculated in the code with the value that is measured experimentally. This value of $\langle \beta_T \rangle$ contains contributions of stored energy from both the thermal plasma and beam source. The comparison is made in Fig. 8. Typical error bars in both the code calculations and the experimental measurements are shown. For $\langle \beta_T^{dia} \rangle < 2\%$, the agreement between the two values is within 5%. For these cases, the measured slowing-down spectra agreed with those calculated classically. Greater disagreement is seen for values $> 2\%$. At these higher values of $\langle \beta_T^{dia} \rangle$, beam ion loss due to the presence of internal kink mode (fishbone) oscillations was present [4]. Consequently, the classical picture of beam ion slowing down, which ignores fast ion loss due to fishbone oscillations [5], is no longer valid and leads to an overestimate of the beam stored energy and thus the calculated $\langle \beta_T^{dia} \rangle$. An heuristic model of beam ion loss due to fishbones has been developed based on the time variations of both the neutron flux and fast charge-exchange efflux observations [4,5], and does, in fact, bring the measured and calculated values of $\langle \beta_T^{dia} \rangle$ into agreement. Such a model was used for cases in which the disagreement between the calculated and measured values of $\langle \beta_T^{dia} \rangle$ was greater than 10%. These cases will not be included in later discussions, however; only cases for which the classical picture of beam slowing down is a valid description of the beam ion content will be discussed here.

Figure 8 shows the comparison of $\langle \beta_T^{dia} \rangle$ values only for discharges with toroidal fields up to 1.5 T, taken during 1982. The verification of the 2.2 T points is more difficult, in part because diamagnetic flux measurements were not available for operation at this field in 1981. In these cases, the $\ell_i/2 + \beta_{pol}$ calculated in the transport code during NBI was found to exceed the magnetically measured value by as much as 30%, and the calculated change in this parameter, going from the ohmic to the NBI phase often exceeded the measured change by $\sim 20\%$. The source of this discrepancy is unknown. Uncertainties in the magnetic value of $\Delta(\ell_i/2 + \beta_{pol})$ were due primarily to the uncertainty in $\Delta\ell_i/2$ caused by changes in I_p between the ohmic and NBI phases. In the TRANSP calculations, the possible sources of error are the ion temperature profile shape, which was not measured, and as yet unidentified nonclassical effects that would result in premature loss of fast ion energy. For the $B_T < 1.5$ T cases, there was little discrepancy between the measured and calculated $\ell_i/2 + \beta_{pol}$ values. Since both the thermal content of the bulk ions and the beam distribution depend on Z_{eff} , uncertainties in this value would lead to uncertainties in both β_{pol} and $\ell_i/2$. However, a factor of two increase in Z_{eff} , from two to four, reduced the calculated value of $\ell_i/2 + \beta_{pol}$ by only 10 to 15%. As will be seen later, some of the 2.2 T cases, especially at $I_p \approx 320$ kA, appear to have significantly improved energy confinement times relative to those either at lower B_T or other 2.2 T cases at higher currents. Since the measured $\ell_i/2 + \beta_{pol}$ of these anomalous cases did not show the same relative improvement, their higher confinement may, in part, be a reflection of the overestimate in the calculated stored energy.

3.1 Power balances

To understand the processes governing the scaling of the energy confinement time, it is essential to determine first the dominant energy loss channels in the plasma. A typical example of the volume integrated ion and electron power balance profiles is shown in Figs. 9a and 9b for a 1.5 T, 300 kA, three beam case. For this case, $T_e(0) \approx 1.25$ keV, $T_i(0) = 1.90$ keV, and $\bar{n}_e = 3.8 \times 10^{13} \text{ cm}^{-3}$. The energy source terms are plotted in the top half of the panels, and the loss terms are shown in the bottom half of the panels. For this case, the direct beam heating of the ions and electrons are approximately the same globally ($r \sim a$), although because of the high T_e values in the core of the discharge ($r < a/2$), the ion heating there is greater than that of the electrons. The beam ion thermalization, the deposition of energy by beam ions slowed down to $3/2kT_i$ (at which point the beam ion is considered to become a thermal ion), is small due to the relatively low T_i . It is seen in the figure that ion conduction is an important energy loss mechanism in the ion power balance out to 15 to 20 cm. For this case, a neoclassical multiplier of two, using Chang-Rinton as the neoclassical value, was needed to bring the fits to the measured and calculated thermal charge-exchange spectra into agreement.

Beyond the 15- to 20-cm radius, the coupling of the ions to the electrons, Q_{ie} , becomes the dominant loss term for the ions. The ion-electron coupling is comparable to the electron heating by the beam, P_{be} , over the outer region of the profile, with the ohmic input power being only $\sim 10\%$ of the total electron heating. The dominant loss term for the electrons is clearly the anomalous transport loss, which is attributed to thermal conduction. In fact, the electron conduction very substantially exceeds the ion transport loss across the plasma. Inside of $r = 3a/4$, the radiated power

is almost negligible. The flattening of the volume integrated conduction loss, and the increase in the integrated radiation loss for $r > 3a/4$ is due to a peaking of the bolometer signal in that region. Nevertheless, the volume integrated electron conduction loss is still about twice that of the radiation loss at the plasma edge, accounting for close to 70% of the total loss in this case.

The route by which the input energy ultimately reaches the wall remains unclear. Global power balance studies of these graphite rail limiter PDX discharges concluded that, on the average, 30% of the input beam power was measured by the bolometers [6]. This result was supplemented by spectroscopic measurements, which indicate that the radiated power was due to oxygen, titanium, and carbon, in approximately a 5:4:1 ratio. The total spectroscopic radiation accounted for only $\sim 1/2$ of the bolometer signal; the difference between the bolometric and spectroscopic results is presumably due to charge-exchange efflux of beam particles. Monte-Carlo calculations of the fast neutral efflux indicate that this level of signal can easily be explained in this manner. Later studies also indicated that from 10 to 30% of the input power for the 1.0 to 1.5 T discharges was deposited on the limiter, with the deposited power increasing linearly with input power at low currents (250 to 350 kA), but independent of input power, above the 2 MW level, at high I_p (400 to 450 kA) [65]. Thus, from 40 to 60% of the input power is typically unaccounted for.

The power scaling of the volume integrated ion and electron power balances at $r = a/2$ is shown in Figs. 10a and 10b for the 1.2 T, 300 kA discharges. In the core region of these discharges, the beams preferentially heat the ions, which typically receive a 30% greater fraction of the total beam heating power than the electrons. Calculated transport losses for the

ions account for 20 to 40% of the total energy loss, with the major part of the transport loss due to thermal conduction. For these cases, the ion neoclassical multiplier is < 1 at the one beam level and is from 1.5 to 2 at the higher powers. Thermal charge-exchange losses at all powers are negligible. The dominant global energy loss for the ions is again through ion-electron coupling, which increases as a function of beam power owing to the larger difference between the electron and ion temperatures at the higher beam powers. The central electron temperature increases only from 0.70 to 1.0 keV going from one to four beams, while $T_i(0)$ increases from 0.9 to 1.8 keV. The line-averaged density for these cases was 3.7 to $3.8 \times 10^{13} \text{ cm}^{-3}$. Q_{ie} accounts for 40 to 45% of the total power input into the electrons. The dominant loss term for the electrons at all power levels is through thermal conduction; the electron conduction loss is typically three to four times greater than the total ion transport loss. Convection and radiation losses are negligible for the electrons within $r = a/2$. Towards the plasma edge, radiation, and possibly convection, can be locally important, but even out to $r = a$, the dominant global loss process appears to be electron thermal conduction. The power balance scaling shown in Figs. 10a and 10b is typical of those at other currents and fields.

In Figs. 11 to 13 are shown the current scalings of the ion and electron power balances for ohmic, two beam, and four beam 1.2 T discharges. The ohmic power balance shown is volume integrated out to $r = 3a/4$, since the peak of the power deposition is within this radius for these cases. All of the plasma heating in the ohmic case (Fig. 11) is taken to be to the electrons through ohmic heating. The ion-electron coupling term, with the power flow to the ions, is an order of magnitude less than the Joule heating. In this ohmic case, and typically for other ohmic cases as well, the electron thermal

conduction is the dominant loss mechanism, accounting for 60 to 70% of the total loss.

The two and four beam power balances in Figs. 12 and 13 are volume integrated out to $r = a/2$, and similar behavior is seen at both power levels. Note that in the two beam case the scales for the electron and ion power balances are different. For both cases, the beam heats the ions preferentially. For the two beam case, ion transport accounts for 25% of the total energy loss at the lower currents, with most of this loss due to ion thermal conduction. Since the ion convection loss is seen not to vary much with I_p , the reduction in conduction loss at higher currents reduces the fraction of loss due to transport to $\sim 20\%$, with the bulk of the remainder of the loss through ion-electron coupling. The ion neoclassical multiplier for these cases was ~ 2 at 250 and 300 kA, and was < 1 at the higher currents. Similar to the cases shown in Fig. 9a, thermal charge-exchange losses account for only a few percent of the total ion loss. In the four beam case (Fig. 13), the ion conduction and convection losses remain approximately constant with I_p , with most of the energy flowing into the electrons. The ion neoclassical multiplier ranged from 0.5 to 2 with no systematic variation with current. The coupling term is about 50% of the total input into the electrons, independent of current. Electron conduction losses dominate the total power losses, accounting for 60 to 70% of the total input power to the bulk plasma. The deduced convection loss for both species is dominated by the beam source in this region of the plasma. The source of ions and electrons due to the ionization of thermal neutrals is dominant in the outer region of the plasma, where the thermal neutral densities are higher than on the inside by one to two orders of magnitude, and are sensitive to the choice of edge recycling time.

For comparison with the lower toroidal field cases (see Figs. 10a and 10b), the three beam, 2.2 T power balances are plotted in Figs. 14a and 14b as a function of current. One difference is that the ion energy source due to beam ion thermalization is no longer so insignificant, especially at the higher currents. This is because the ion temperatures are higher, so that the beam ions need to degrade less in energy before being considered thermal. Ion convection loss is also greater at these toroidal fields directly and indirectly because of the higher ion temperature; directly since convection $\sim 5/2(n_i T_i v_i)$, v_i being the ion radial flow velocity, and indirectly due to better thermal neutral penetration [edge $T_0 = T_i(0) \sim 30$ for the code calculations] and thus a greater thermal source of ions in the plasma core region. This latter effect, which is model dependent, is not the dominant one, however. If the neutral recycling temperature is assumed to be 5 eV, and independent of ion temperature, the neoclassical multiplier necessary to reproduce the fit to the measured charge-exchange spectrum is, at the highest $T_i(0)$, only $\sim 30\%$ greater than that with the higher neutral temperature model. Note also in Fig. 14a that the convection loss term is generally comparable to the beam ion thermalization both in these and the 1.2 T cases.

Much of the ion energy still flows into the electrons through the classical coupling process. The ion charge-exchange loss is higher in these cases, but it is still not important, accounting for only 10 to 15% of the ion energy loss. For the electrons, the thermal conduction loss dominates. At 320 kA, electron thermal conduction losses account for $\sim 60\%$ of the total input power, but at 200 and 480 kA, electron conduction is responsible for closer to 75% of the loss.

3.2 Ion heating and confinement

In this section, the ion heating results will be discussed and related to ion transport. The ion temperatures that will be shown are the central values calculated from the ion power balance in the transport code. For the calculations, an ion thermal diffusivity of typically one to three times the neoclassical value given by Chang and Hinton was necessary to bring the fits to the calculated and measured thermal charge-exchange spectra into agreement. For some cases, especially those at lower beam powers and toroidal fields, the required neoclassical multiplier was less than one. For these cases the central ion temperatures were low (~ 1 keV), and the relative uncertainty in the fits to the measured charge-exchange spectra could be up to $\sim 30\%$.

The central ion temperature, $T_i(0)$, is plotted as a function of total heating input power to the plasma for four different currents at a fixed toroidal field of 1.2 T in Fig. 15. The total heating power, P_{heat} , is defined as the absorbed (injected power less the shine-through) plus ohmic power, less the bad orbit and beam ion charge-exchange losses. For the discharges shown here, the line-averaged density was typically 3.5 to $4.0 \times 10^{13} \text{ cm}^{-3}$ except at the highest current, highest power case where $\bar{n}_e \sim 4.8 \times 10^{13} \text{ cm}^{-3}$. The ion temperature is seen to increase linearly with P_{heat} , and to exhibit a slight current dependence for $1.2 < P_{\text{heat}} < 2.7 \text{ MW}$. Above the 3 MW level, a stronger current dependence of $T_i(0)$ may be seen, and the ion temperature appears to begin to saturate with beam power. This modest saturation is also seen in $T_e(0)$ and $\langle \beta_T \rangle$, for this toroidal field, at the higher beam powers. At this toroidal field, even for the high power discharges, fishbone oscillations were weak, and no clear depletion of the fast ion charge-exchange spectrum was observed. Therefore it was possible to model the beam slowing down

classically so that the calculated $\langle \beta_T \rangle$ agreed to within 10% of the measured value. Consequently, it is not clear whether the saturation in $T_e(0)$ and $T_i(0)$ at the higher powers are due in part to the presence of fishbone oscillations or to some nonfishbone mechanism. At 1.5 T, neither fishbone oscillations nor any saturation in $\langle \beta \rangle$ as a function of beam power were observed.

In a previous analysis of the high field subset of the present database, the ion heating efficiency parameter, η_i was shown to reflect the central ion energy confinement in the neutral beam heated phase [1]. This parameter, η_i , is defined to be $\bar{n}_e \Delta T_i(0) / P_{\text{abs}}$, where $\Delta T_i(0)$ is the difference in central ion temperature between the NBI and a "baseline" ohmic phase, and $\bar{n}_e$ is the line-averaged density during NBI. It was shown further that η_i and consequently $\tau_{Ei}(0)$, the central thermal ion energy confinement time, for that subset depended primarily on the ion transport loss from the plasma. To test the relationship between $\tau_{Ei}(0)$ and η_i , in this expanded dataset, τ_{Ei} at $r = 10 \text{ cm}$ ($\sim a/4$) is plotted as a function of η_i in Fig. 16 for all cases in which the ion neoclassical multiplier is ~ 1 . Cases for which this multiplier is < 1 have been excluded due to the large relative uncertainties in $T_i(0)$. The ion confinement, as used here, takes into account thermal charge exchange as well as transport loss. It is clear from Fig. 16 that for most cases the ion heating efficiency increases with increasing central ion confinement.

The ion heating for the 1.2 T case is shown in Fig. 17, in which $\bar{n}_e \Delta T_i(0)$ is plotted as a function of P_{abs} . The ion heating efficiency parameters, the η_i 's, for these cases are typically between one and two; for the 2.2 T cases η_i was found to be as high as three to four for $I_p > 320 \text{ kA}$. For the higher field cases, $\bar{n}_e \sim 2.5 \text{ to } 3.5 \times 10^{13} \text{ cm}^{-3}$, while at 1.2 T $\bar{n}_e \sim 4 \times 10^{13} \text{ cm}^{-3}$. It is also seen in Fig. 17 that the heating efficiency appears

to increase with increasing current, due primarily to the small but systematic dependence of $T_i(0)$ on plasma current. Furthermore, the ion heating is almost linear with absorbed power, except at the highest powers, and especially for the 400 kA cases, where $T_i(0)$ saturates.

Ion thermal conduction, in most cases, is the dominant ion loss mechanism in the central portion of the discharge. Therefore, it is expected that the central ion confinement should be related to the ion thermal diffusivity in an inverse fashion. This relation is demonstrated in Fig. 18, in which τ_{Ei} at $r = 10$ cm is plotted as a function of the ion thermal diffusivity, χ_i , at that same radius. The ion thermal diffusivity is $\alpha \chi_i^{NC}$, where χ_i^{NC} is the neoclassical (Chang-Hinton) diffusivity and α is the necessary neoclassical multiplier. The general trend that is seen in Fig. 18 is for τ_{Ei} to decrease with increasing χ_i .

How well neoclassical theory explains the variation of χ_i can be determined by examining the variation of the neoclassical multiplier with local discharge parameters (i.e., is there any parametric dependence of χ_i over and above that given by neoclassical theory?). Results of regression fits to the neoclassical multiplier indicate that there is no systematic dependence of χ_i on any of the other chosen variables (B_T , B_p , T_i , n_i , P_{heat}). As previously mentioned, the value of α needed to model the fit to the measured charge-exchange spectrum was generally one to three. Because of the lack of parametric dependence of α , it is believed that the range of α 's needed primarily reflects uncertainties in the experimental measurements.

To summarize this section briefly, for the range of cases studied the ion heating efficiency parameter, η_i , was found to reflect the central ion energy confinement, and in particular the ion transport confinement time. The heating efficiency parameter increased with increasing current and increasing

toroidal field (decreasing density), although some degradation was observed at the highest beam powers. Because most of the central energy loss in the ion channel is through thermal conduction, the central ion confinement was found to relate directly to χ_i in an inverse fashion. Finally, the behavior of the ions could be well explained by neoclassical theory; the range of the neoclassical multipliers needed to model the fit to the measured charge-exchange spectra merely reflects the inherent uncertainties in the experimental data.

3.3 Electron heating and confinement

In this section the electron heating and energy confinement will be discussed in light of the profile effects discussed earlier (e.g., the broadening of T_e with decreasing q at fixed beam power). It will be shown that while there is only a slight variation of the NBI electron confinement time with toroidal field, its variation with plasma current is strong.

The electron energy confinement time at $r = 0.6a$ (25 cm) is plotted in Figs. 19a and 19b as a function of toroidal field for two different ranges of currents and for 2 to 2.5 MW of heating power. The electron energy confinement time is defined to be

$$\tau_{Ee} = \frac{U_e}{P_{\text{heat},e} - \partial U_e / \partial t}$$

where $P_{\text{heat},e}$ is the total heating power to the electrons (beam + ohmic + ion-electron coupling) and U_e is the stored electron energy. Confinement times at $r = 0.6a$ are presented because the T_e and n_e profiles are known with most accuracy within this radius.

For the lower current range, 200 to 320 kA, there is little systematic variation of τ_{Ee} with toroidal field within the error bars of the calculated values. The confinement times range from 6 to 10 msec, except for a single, anomalously good case at 320 kA and 2.2 T which exhibited a τ_{Ee} of 19 msec. For the cases shown in Fig. 19a, $T_e(0)$ ranged from 0.65 keV at 0.7 T to 1.5 keV at 2.2 T (2.0 keV for the higher confinement, 320 kA point). At higher currents, Fig. 19b, the electron confinement exhibits a slight, but by no means dramatic, tendency to increase with toroidal field, with τ_{Ee} increasing from 9 to 14 msec as B_T increases from 1.2 to 2.2 T. For this subset of cases, $T_e(0)$ ranges from 1.5 keV to 2.0 keV, increasing with toroidal field.

Two factors are important in understanding the scaling of electron energy confinement time with B_T . The first is the variation of an electron energy parameter, defined to be $\bar{n}_e T_e(0)/P_{heat}$, with toroidal field, and the second is the temperature profile peakedness, which was seen to increase with increasing q . In Fig. 20, the electron heating parameter, defined as $\bar{n}_e T_e(0)/P_{heat}$, is plotted versus toroidal field for both ranges of currents. It is obvious from the figure that the central electron heating is a strong function of toroidal field, for both current ranges, indicating an improvement of central electron confinement with increasing B_T . The results from Figs. 19a and 19b indicated that the electron confinement at 0.6a was, at best, only weakly increasing with B_T at the higher currents, and independent of B_T at lower I_p . These contrasting results indicate that profile shape is an important consideration, and that the increasing central electron confinement with B_T at constant current is offset at larger radii by the narrower T_e profile at higher q .

The effect of profile shape may be nearly eliminated by examining the variation of electron confinement with toroidal field (and thus plasma

current) at constant q . These results are summarized in Table 1, which lists the central electron heating efficiency, electron energy confinement, and the volume integrated stored energy and power to the electrons, all at $r = 0.6a$. The results given in the table indicate that at constant q the central electron heating increases with toroidal field, consistent with the results seen in Fig. 20. However, at constant q , with differences in profile shape virtually eliminated, the electron confinement increases with toroidal field and plasma current. It is seen from the table that while the power input to the electrons remains relatively constant, it is the increase in stored energy with B_T and I_p , due to the increase in T_e across the profile, that yields the higher confinement times.

The variation of electron energy confinement with power and current at fixed toroidal field is shown in Figs. 21a to 21c. For the 1.2 T discharges, the confinement time appears to degrade with the onset of NBI to values of 6 to 10 msec. The confinement degrades from the ohmic to the one beam power level, and then weakly to the two beam power level. Once at the two beam level, the confinement time shows little further degradation with increasing beam power. For ohmic discharges $\bar{n}_e = 2.5 \times 10^{13} \text{ cm}^{-3}$, while $\bar{n}_e = 4 \times 10^{13} \text{ cm}^{-3}$ during NBI. Little current dependence of confinement is seen at the one or two beam level, while at higher beam powers the electron energy confinement times are clearly higher at the higher currents. This last result is due to the broader T_e profiles, and thus greater stored energy, at the higher currents (lower q at fixed B_T).

These results hold too for the higher toroidal field cases. At 1.5 T, τ_{Ee} degrades from the 20 to 27 msec range for ohmic discharges to values from 8 to 14 msec during NBI. The confinement times again degrade continuously from the OH to the two beam level, except for the 400 kA cases, although the

degradation from the one to two beam level at this toroidal field is even less than at $B_T = 1.2$ T. The confinement times for $B_T = 1.5$ T are seen also to increase with increasing current. The marginal increase of the 1.5 T confinement times with respect to those at 1.2 T at the same currents was discussed earlier in this section, and it is due primarily to an increase in the central electron confinement.

At 2.2 T, the points exhibit more scatter. For the NBI discharges, the 200 kA cases, from 5 to 8 msec, lie well below those at higher currents. At 320 kA, the situation is more confused. The three beam discharge, ($P_{\text{heat}} = 2$ MW) exhibiting an electron energy confinement time of 19 msec, is anomalously better than any of the other cases. The discharge or operational conditions that led to such good confinement have not been identified. The uncertainty in the 2.2 T data set is greater than at other toroidal fields. Overall, there is little current scaling between the 320 kA and 480 kA cases, aside from the anomalous 320 kA case, as reported previously [2].

The ohmic confinement times do not appear to follow any simple separate I_p or B_T scaling. It is found, however, that τ_{Ee} scales roughly with $\bar{n}_e q^{1/2}$ (Fig. 22) up to a value of $\bar{n}_e q^{1/2} = 5 \times 10^{13} \text{ cm}^{-3}$. Beyond this value, τ_{Ee} degrades by $\sim 30\%$ due to relatively larger thermal conduction losses. The large values of $\bar{n}_e q^{1/2}$ for these two cases are due to higher densities; $\bar{n}_e \sim 3.5 \times 10^{13} \text{ cm}^{-3}$ for these cases. The linear portion of the scaling, due either to increasing $\bar{n}_e$ or q , or both, is consistent with ohmic scalings reported previously [66-75]. The degradation in confinement at high values of $\bar{n}_e$ has been documented on [66], and has been modeled qualitatively based on a degradation of electron confinement at high $\bar{n}_e$ [30].

To summarize this section, it is found that although the central electron heating efficiency increases with toroidal field for NBI plasmas at

fixed current, the electron energy confinement time at $r = 0.6a$ has only a weak, if any, dependence on B_T , due to T_e profile narrowing at higher q . At fixed q , the electron confinement improves with increasing toroidal field and plasma current. It was further seen that at fixed B_T , τ_{Ee} increased with increasing I_p due to a broadening of the T_e profile, indicating that it is I_p that is the most important parameter for gross electron confinement scaling. Electron confinement times degrade from the ohmic to the NBI phase of the discharge, but beyond the two beam power level little further degradation is seen. Finally, the ohmic confinement times themselves scale linearly with $\bar{n}_e q^{1/2}$ up to $\bar{n}_e = 3.5 \times 10^{13} \text{ cm}^{-3}$; beyond this value, the ohmic confinement degrades.

3.4 Thermal confinement times

The thermal confinement times to be presented in this section are determined by considering the stored energy in and input energy to only the thermal plasma. The total thermal energy confinement time is defined to be

$$\tau_E = \frac{U_e + U_i}{P_{\text{heat}} - \partial(U_e + U_i)/\partial t}$$

where $U_{e,i}$ are the total stored energies in the bulk electrons and ions respectively, and P_{heat} is the total heating power as defined previously. For discharges in equilibrium ($\partial/\partial t = 0$), the expression for τ_E may be simplified. In the limit in which the majority of power is lost through the electron channel, $\tau_{Ei} \gg \tau_{Ee}$, and the ions have little influence on the total confinement except to store additional energy. This is essentially the

situation for the lower toroidal field (< 1.5 T) cases studied. Comparable amounts of power are lost through the electron and ion channels in some of the 2.2 T discharges, and then ion confinement does influence the total thermal confinement.

In Figs. 23a and 23b, the total thermal confinement times for heating powers from 2 to 2.5 MW are plotted as a function of toroidal field for the two current ranges, 200 to 320 kA and 350 to 480 kA. The total thermal confinement times behave in a similar fashion with toroidal field as the electron confinement times. At low currents, little variation of τ_E with toroidal field is observed, with τ_E values ranging from 11 to 16 msec. The 31 msec confinement case at 2.2 T is, again, the anomalously good 320 kA case. For all but this case in this current range, the total thermal confinement is dominated by the electrons; at least 70% of the input power is lost through the electron channel for these discharges. For the 320 kA case, however, the ions are responsible for $\sim 40\%$ of the energy loss. Thus, τ_E for this case is influenced by both species. The greater relative energy loss in the ions for the 320 kA case is due in part to enhanced central ion convection relative to the electron and ion conduction loss (see Fig. 14), but the dominant effect on total confinement is the anomalously high electron energy content.

At higher currents, the total confinement time appears to increase only marginally with increasing toroidal field, following the electron confinement trend. The 2.2 T point is at a higher current than the others (480 kA as compared to 350 to 400 kA). At this field, approximately 35% of the total energy loss is through the ion channel. As discussed earlier, the ion confinement improved with both increasing current and toroidal field (decreasing density); consequently, the higher confinement time of the 2.2 T, 480 kA case, which is influenced by both the electrons and the ions, reflects

the increased confinement at higher currents for both species, and perhaps the increased confinement at lower densities for the ions.

At constant q , as in the case for the electrons, the total thermal confinement improves with increasing toroidal field and plasma current. For the three cases shown in Table 1, the total thermal confinement times are 11, 14, and 27 msec for the 1.2, 1.5, and 2.2 T cases, respectively. For the two lower field cases, the electron loss dominates the confinement, while at the highest field the confinement is influenced by both the ions and the electrons, as noted earlier. Since little dependence of τ_E on toroidal field is seen at fixed current, the increase in confinement at fixed q appears to be due to the effect of increasing plasma current.

The dependence of confinement on heating power and, with more detail, on plasma current is shown in Figs. 24a to 24c, where τ_E is plotted as a function of P_{heat} for the various current scans at each toroidal field. The 0.7 T and 1.0 T power scan results are not presented because of the presence of severe fast ion loss due to fishbone oscillations at the higher beam powers. At 1.2 T (Fig. 22a), the ohmic confinement times of 29 to 37 msec are substantially degraded with the onset of NBI; the one beam cases reflect the transition from the ohmic value to the ultimately degraded values for $P_{\text{heat}} > 2.0$ MW. At powers higher than 2.0 MW, only a slight degradation of τ_E with beam power is observed, consistent with the weak, $P^{-1/3}$ scaling deduced from the magnetic measurements [4,30,39,40]. For the NBI discharges, the thermal confinement increases with increasing current at fixed P_{heat} , due primarily to an increase in plasma stored energy reflecting the broader T_e profiles. As an example, P_{heat} in the two beam discharges varies from 2.2 to 2.6 MW, while the total plasma stored energy increases from 27 kJoules at 250 kA to 46 kJoules at 400 kA, reflecting the increase in τ_E from 11 to 20 msec.

The NBI confinement variation in the 1.5 T discharges (Fig. 24b) exhibits similar behavior as in the 1.2 T cases. Confinement times degrade from ohmic values of 40 to 50 msec to values of 11 to 15 msec at 250 kA and to values of 20 to 24 msec at 400 kA. The lack of any further degradation in τ_E with P_{heat} beyond the two beam level is observed for the 400 kA discharges.

The 200 kA, 2.2 T discharges (Fig. 24c) exhibit a dramatic degradation in confinement going from OH to NBI due to the extreme narrowing of the T_e profiles, relative to those in the higher current cases, even at the one beam level. For the higher powers, there is again little systematic variation of τ_E with P_{heat} . The energy confinement at 480 kA shows practically no degradation from ohmic to NBI because of the initially low ohmic confinement value at this current (26 msec), similar to the high current, low density results from PLT [21,22]. For 480 kA ohmic discharges $\bar{n}_e = 2 \times 10^{13} \text{ cm}^{-3}$, while $\bar{n}_e = 2.3$ to $3.4 \times 10^{13} \text{ cm}^{-3}$ for the NBI plasmas at the higher currents, 320 and 480 kA, the τ_E values are substantially higher than those at lower toroidal fields for the middle power range, and this, as discussed, is a result of the ion influence on confinement as well as the anomalously high electron stored energy content in some of the 320 kA cases. Little current variation in τ_E is seen from 320 to 480 kA; in fact, while $\tau_E = 30$ to 31 msec for 320 kA for one and two beams, $\tau_E = 25$ to 27 msec for 480 kA at the same power level. The current dependence is recovered for the 200 kA discharges, which exhibit a much reduced confinement time.

The scaling of the total ohmic confinement times with $\bar{n}_e q^{1/2}$ is shown in Fig. 25, and it reveals a trend similar to that seen in the τ_{Ee} scaling shown in Fig. 22. The total confinement time appears to increase linearly up to $\bar{n}_e q^{1/2} = 5 \times 10^{13} \text{ cm}^{-3}$; in fact, assuming that τ_E scales linearly with density, a regression analysis of the ohmic points with $\bar{n}_e q^{1/2} < 5 \times 10^{13}$

cm^{-3} yields the fit $\tau_E/\bar{n}_e \sim q^{0.51}$. The $q^{1/2}$ dependence explains the inverse scaling of ohmic τ_E with plasma current seen most clearly in the 1.2 and 2.2 T cases (Figs. 24a and 24c). The almost linear scaling of τ_E with $\bar{n}_e q^{1/2}$ is also consistent with the results previously referenced [30,66-75]. Degradation of the total confinement time is seen for $\bar{n}_e q^{1/2} > 5 \times 10^{13} \text{ cm}^{-3}$, due, as mentioned, to increased electron thermal conduction losses. Saturation in confinement at high density has been seen on both Alcator C [74] and D-III [66].

In summary, the trends of the total thermal confinement times follow closely those seen for the electrons alone. Little variation of τ_E with toroidal field is seen except at the highest currents and fields, where the ions, with confinement times that are greater than those of the electrons, influence the total thermal confinement. Beyond the two beam level, a weak dependence of τ_E on P_{heat} is observed, consistent with the results deduced from the magnetics measurements. The strongest dependence of τ_E is on I_p , and it is due to the improved electron confinement at the higher currents, with some influence by the ions at the highest toroidal field. Finally, the ohmic confinement scales linearly with $\bar{n}_e q^{1/2}$, except at high densities, consistent with observations on other tokamaks, although the range of $\bar{n}_e q^{1/2}$ presented here reflects primarily the range of q_{cyl} .

3.5 Electron thermal diffusivities

From the results presented in the previous section, it was concluded that for most of the cases the electron channel dominated the power balance, and consequently it was the electrons that influenced thermal confinement the most. In particular, it was shown that an anomalous electron transport loss, attributed to thermal conduction, was the most important single energy loss

channel for the thermal plasma within $r/a=0.5$, accounting for, in some cases, up to 80% of the total energy input. This result is also true in the global sense (i.e., when the power balance is volume integrated out to $r=a$). At the edge, the volume integrated radiated power is typically only ~ 30 to 40% of the global conduction loss, and the global electron conduction loss accounted for 65 to 75% of the total energy loss at lower fields ($B_T < 1.5$ T). At higher toroidal fields (2.2 T), the ion transport loss was a larger percent of the total than at lower fields. For the global energy balance in these cases, the ion transport loss is not insignificant, with the ions accounting for 25% of the total loss. However, the electron conduction loss was still the largest single loss term, accounting for $\sim 50\%$ of the global plasma energy loss. Therefore, to understand the parametric scaling of the total thermal confinement times, it is necessary to examine how the electron thermal diffusivity, χ_e , varies across the discharge for the range of plasma parameters studied.

The local values of χ_e at each radius were determined from the calculated volume integrated electron conduction loss at that radius and the measured local values of n_e and $\partial T_e / \partial r$. Since the discharge parameters (i.e., current, voltage, line-averaged density) were measured with good accuracy ($\pm 5\%$), the major sources of uncertainty in χ_e are systematic uncertainties in $T_e(r)$ and $T_i(r)$. Because of the smoothing of the T_e profiles in the preprocessing routines (described in Sec. 2), random errors within the error bars of the Thomson scattering measurements, which will modify $T_e(r)$ and thus $\partial T_e / \partial r$, yield a ± 15 to 20% possible error in the value of the deduced χ_e . If 50% of the bolometer signal is due to fast charge-exchange efflux, instead of practically all of it being due to radiation as assumed in the calculation of χ_e , then the χ_e will be underestimated by $\sim 20\%$ at $r=a/2$ and $\sim 40\%$ at the

plasma edge, where the bolometer signal is the largest. If a systematic error of $\sim 15\%$ is assumed for the T_e profile, T_e for all radii being too low or too high by 15% , then this uncertainty propagates into a ± 25 to 35% error in χ_e across the profile. It was previously argued that the range of neoclassical multipliers needed to model $T_1(0)$ represented uncertainties in the measured thermal charge-exchange spectra. The uncertainty in the neoclassical multiplier could lead to a $\pm 25\%$ uncertainty in χ_e at $r=a/2$. Related uncertainties in the calculation are the $T_1(r)$ and neutral density profiles, the latter assumed to be poloidally and toroidally symmetric. Taking all these effects into account, it is believed that the values of χ_e at $r=a/2$ are known to within a factor of ~ 1.7 . This uncertainty in χ_e is much larger near the center of the plasma where the electron temperature profile is flat and $\partial T_e / \partial r$ is not well known, and because of this, as well as the lack of a reliable sawtooth model inside $q=1$, the χ_e values within $q=1$ will not be stressed. At the very edge of the plasma, where the relative uncertainties in T_e are largest, χ_e is also poorly known. Hence the discussion of χ_e is restricted to the range outside the $q=1$ radius and within $r=7a/8$.

Figure 26 shows the deduced χ_e profiles in the 1.2 T, 300 kA power scaling case for which the power balances are shown in Fig. 10. For the sake of clarity, the χ_e profile in the three beam case has been omitted in the figure. The large dot on each χ_e profile represents the calculated position of the $q=1$ surface, which agrees to within several cm with the results from the soft X-ray array. The total thermal energy confinement times for the four cases shown are given in the top left-hand corner of the graph. The confinement time degrades from 36 msec to 24 msec to 11 msec going from the ohmic to the one and two beam discharges, and then remains relatively constant, decreasing to 9 msec, for the four beam case. The magnitude of the

one beam χ_e profile is greater than that in the ohmic case by up to a factor of two in the region of interest. The local maximum in the ohmic χ_e at $r=10$ cm is due to a local flattening of the ohmic T_e profile. Associated with the large decrease in confinement time, the magnitude of the two beam χ_e profile outside $q=1$ is about twice that of the one beam profile. Going from the two to the four beam case, where the τ_E 's are comparable, the χ_e 's are comparable in most of the region outside $q=1$. For all cases, χ_e increases dramatically to values $> 10^5$ cm²/sec in the outer portion of the plasma, being greater than the value given by Bohm by one to two orders of magnitude.

The dependence of χ_e with plasma current is shown in Fig. 27 for 1.2 T two beam discharges. Again, the total thermal energy confinement times are given in the top portion of the figure. At 250 and 300 kA, the energy confinement times are equal, and the magnitude of the χ_e 's for these cases are comparable outside the $q=1$ surface. As the current increases to 400 kA, q near the edge decreases, and the value of χ_e in the region of interest decreases. At the lower currents, χ_e near the edge is $\sim 3 \times 10^5$ cm²/sec, while at 400 kA it is 7×10^4 cm²/sec, indicating a reduction in χ_e near the edge by a factor of four. This reduction in edge χ_e is reflected in the doubling of τ_E going from 250 to 400 kA.

That τ_E responds to changes in the χ_e from $r(q=1)$ to $r = 7a/8$, and that the value of χ_e itself decreases with decreasing local q , are results seen in most of the cases studied. However, as previously discussed, the scaling of τ_{EE} and τ_E appeared not to depend on q as much as it did on plasma current. Whether it is q or poloidal field (plasma current) which controls the χ_e variation can be determined by examining τ_E and χ_e at fixed q and fixed beam power level for a range of B_T and I_p 's (no a or R scaling is available). The results are shown for $q_{cyl} \sim 3$ in Fig. 28. As the toroidal field and current

increases from 1.2 T, 250 kA to 2.2 T, 480 kA, τ_E increases from 11 to 27 msec. The slight increase in τ_E for 1.2 to 1.5 T reflects the reduced χ_e in the $r = 20$ to 32 cm region for these cases. As the toroidal field increases to 2.2 T, and the confinement time more than doubles from the 1.5 T value, χ_e beyond $r=30$ cm decreases by a factor of $\sim$ two. Another element in the increase in total confinement in this high field case is the stronger role played by the ions. At fixed q , then, the confinement increases (χ_e decreases) with increasing toroidal field and plasma current, while at fixed B_T , τ_E improves and χ_e decreases with decreasing q (increasing poloidal field). This indicates that it is not the q -variation (or more precisely the B_T/I_p variation since variations in a/R were not significant for these cases) that controls the variation of χ_e . Rather, this suggests that it is the variation in the local poloidal field that is important.

The deduced experimental values of χ_e that are output from the transport analysis code were compared to a large number of predicted and empirical scalings [43, 76-83]. Substantial disagreement was found in all these comparisons (see Appendix 1). One possible explanation of this disagreement is that while many of the theoretical and empirical χ_e scalings were developed for ohmic plasmas, relatively few were developed for plasmas with neutral beam heating. For instance, a χ_e scaling that has been used extensively for NBI tokamak modeling is the INTOR scaling which is based on Alcator A ohmic scaling, and in which $\chi_e = 5 \times 10^{17}/n_e$ (cm²/sec), reflecting the improved energy confinement in ohmic plasmas with increasing density. Until recently on ISX-B [41] and D-III [38], such a density scaling has not been observed with NBI. For the NBI PDX cases presented here, the deduced χ_e values are typically a factor of three to seven greater than that value given by INTOR scaling.

A theoretical scaling which has recently received much attention, and one which was specifically developed for high- β_p NRI plasmas, is that derived from the pressure driven resistive ballooning mode formulation of Carreras and Diamond [43]. The χ_e 's deduced from the ISX-B discharges were found to match reasonably well with a value of χ_e given by $2\chi_e^{C-D} + \chi_e^{INTOR}/2$. The best agreement was found at the highest β_{pol} , where resistive ballooning is believed to be most important. In Fig. 29, the χ_e 's at $r = 30$ cm deduced from the PDX experiments are compared to χ_e^{C-D} at that radius as a function of β_{pol} . It is seen in the figure that the agreement between the two χ_e values is relatively poor, with $\chi_e/\chi_e^{C-D} \sim 6$ to 16 even at the higher β_{pol} . The deduced χ_e typically exceeds χ_e^{C-D} by at least a factor of seven to 20 over the entire profile.

The empirical scaling of Coppi et al., [82] is based on the empirical ohmic scaling of Coppi and Mazzucato [76], but it is extended to take into account auxiliary heating. Comparisons between the deduced χ_e 's and those given by this model generally agreed to within a factor of two to five near $r \sim a/2$. However, it should be noted that the Coppi et al., χ_e scales as $1/T_e$; consequently, the factor of two to five discrepancy reflects the same difference between the measured electron temperature and that predicted by the scaling.

Since little agreement was found with the published theories, it can be justifiably asked whether there are plasma parameters that order χ_e at all. To answer this question, a multiple linear regression analysis has been performed, treating χ_e as the dependent variable and plasma parameters such as E_p , E_T , T_e , and n_e as independent variables. While one advantage of this approach is that it does not presuppose any physical mechanism, care should be taken. For instance, variables that correlate strongly cannot meaningfully be

used simultaneously as independent variables. In the cases studied, it was found that a strong empirical relationship existed between the measured T_e and the calculated poloidal magnetic field, B_p , in the outer region of the plasma. Consequently, either B_p or T_e was chosen as an independent variable, but not both simultaneously. In addition, as will be seen, fits using the same variables differed at different radii; therefore, the analysis was performed at each chosen radius separately, and then for the data at all chosen radii at the same time. In the latter fits, r , the plasma radius, was also used as an independent variable. For the regression fits, values were chosen at 5-cm intervals from $r = a/2$ to $3a/4$ ($r = 20, 25$, and 30 cm). No points inside the $q = 1$ surface were used in the fit, as mentioned, because no sawtooth model was employed and the values of B_p there are uncertain.

The highlights of the regression analysis are presented in Figs. 35a to 35d and are summarized in Table 2. At $r = 20$ cm, all the regression fits are generally poor, with the best fit given by $\chi_e = 10^{5.15} B_p^{-1.45} n_e^{-0.13}$ ($R^2 = 0.51$ where $R^2 = 1.00$ implies a "perfect fit" [84]). This fit is shown in Fig. 30a and is given in Table 2 under case # 1. Since the n_e dependence is a factor of ten less important than that of B_p , an equally good fit is given by $\chi_e = 10^{3.40} B_p^{-1.50}$ with $R^2 = 0.50$ (case # 2). At $r = 25$ cm, the fits to χ_e are slightly better; χ_e is fit equally well with the choice of either B_p or T_e as independent variables. Shown in Fig. 30b is the fit $\chi_e = 10^{7.49} B_p^{-0.99} n_e^{-0.26}$ with $R^2 = 0.61$ (case # 3), although the fit $\chi_e = 10^{11.31} T_e^{-0.53} n_e^{-0.39}$ (case # 4, T_e in eV) has the same coefficient of variation. At $r = 30$ cm, two equally good fits are shown in Figs. 30c and 30d, respectively; $\chi_e = 10^{15.34} B_p^{-0.83} n_e^{-0.82}$ with $R^2 = 0.83$ (case # 5), and $\chi_e = 10^{19.48} T_e^{-0.40} n_e^{-1.02}$ with $R^2 = 0.85$ (case # 6). The latter scaling is similar to the $(n_e T_e^{1/2})^{-1}$ scaling reported by the D-III [37,38] and ISX-B [85] groups. Note that for the

regression fits at the individual radii, χ_e shows a strong inverse dependence on poloidal field. This relationship was inferred from the χ_e variations with q , B_p , and I_p discussed earlier. Also note that the dependence of χ_e on n_e , although small at $r = 20$ and 25 cm, increases steadily with plasma radius. Little dependence of χ_e on toroidal magnetic field is seen at any of the three radii, consistent with the apparent lack of scaling of τ_E and τ_{Ee} on this parameter.

Because it is difficult to infer a global scaling of τ_E from a scaling of χ_e at one particular radius, especially in light of the range of fits seen in cases 1 to 6, fits to the data at all three radii simultaneously were performed. This yields a more globally representative scaling of χ_e in the region where it affects confinement the most. Using only B_p and n_e as independent variables, the resulting fit, $\chi_e = 10^{11.88} B_p^{-1.05} n_e^{-0.60}$ (case # 7) yields a coefficient of variation of $R^2 = 0.54$. The fit to χ_e is better if T_e is substituted for B_p with $\chi_e = 10^{14.17} T_e^{-0.65} n_e^{-0.58}$ (case # 8) with $R^2 = 0.72$. This result, with a weaker n_e dependence than was determined from the recent D-III and ISX-B results, is shown in Fig. 31a. Adding plasma radius to T_e and n_e as independent variables, the fit is not significantly improved, with $\chi_e = 10^{12.92} r^{0.50} T_e^{-0.54} n_e^{-0.56}$ (case # 9) with $R^2 = 0.73$. The results of the fits in cases 8 and 9 indicate that r is not an important scaling parameter when T_e and n_e are chosen as independent variables. The r dependence is significantly stronger when B_p is used instead of T_e . In case 10 (Fig. 31b), the regression fit gives $\chi_e = 10^{0.86} r^{1.94} B_p^{-1.56}$ with $R^2 = 0.70$. With n_e included, the dependence of r and B_p is not as strong, but the overall fit is better, with $\chi_e = 10^{6.57} r^{1.65} B_p^{-1.20} n_e^{-0.37}$ ($R^2 = 0.77$). This latter case, case # 11, is shown in Fig. 31c. Cases # 9 to # 11 reflect the strong dependence of χ_e on plasma radius, with χ_e increasing towards the

outside of the plasma, and also indicate a strong, inverse dependence on poloidal field.

At first glance, the favorable dependence of χ_e on T_e seen in cases # 8 and # 9, for instance, could imply that τ_E should improve with increasing beam power, since the local T_e typically increases with increasing beam power. However, $T_e(r)$ in the outer region also increases with plasma current. To separate the two effects, and concentrate on the variation of T_e due to variations in beam power, regression fits were also performed on cases only in the range from 250 to 300 kA, and the results indicated a dependence much weaker than that shown in Table 2. For instance, the equivalent of case # 8, for the limited current range, yielded $\chi_e \propto T_e^{-0.37} n_e^{-0.83}$ ($R^2 = 0.68$) while that for case # 9 yielded $\chi_e \propto r^{0.38} T_e^{-0.26} n_e^{-0.73}$ ($R^2 = 0.68$). These results suggest that when the variations in T_e due primarily to variations in beam power are considered, the resulting χ_e dependence on T_e is still favorable, but it is weak ($T_e^{-1/3}$).

On the other hand, the χ_e dependence on B_p , due primarily to changes in I_p , can be ascertained by holding beam power fixed. This attempts to eliminate the effect of beam power on T_e . When this is done for the two and three beam cases separately, it is found that the resulting regression fits to the scaling parameters are equivalent to those given in Table 2 for cases # 3, 5, 7, 10, and 11 to within the exponent range shown. However, there is a difference in the numerical coefficient of the scaling. For instance, for the equivalent of case # 11, the numerical coefficient for the two beam case is $10^{4.96}$ while that for the three beam case is $10^{7.27}$, reflecting the increase of χ_e with beam power. This result also indicates that the beam power (through T_e) has little effect on χ_e , and that the T_e scaling is purely derivative of a direct B_p scaling. The increase of T_e and n_e at larger radii

with increasing B_p can be understood as being due to reduced transport coefficients at higher B_p . The covariation of B_p with changes in T_e or n_e cannot be explained, however, by an explicit dependence of B_p on χ_e .

The deduced dependence of χ_e on n_e should not necessarily be taken to reflect a global scaling of τ_E on $\bar{n}_e$. Such a dependence could not be inferred from the PDX dataset because the experimental range of $\bar{n}_e$ was small, within a factor of two for most cases. However, because of profile effects, the local value of n_e from $r = 20$ to 30 cm varied by a factor of four to six from case-to-case. Also for these radii, B_p varied by a factor of ~ 3 and T_e by a factor of ~ 4 from case-to-case.

4. SUMMARY

In conclusion, detailed diagnostic coverage has enabled a comprehensive study of the transport processes governing the energy confinement of PDX plasmas. As also observed in other neutral beam heated tokamaks, the PDX plasmas confined by limiters showed a severe degradation in thermal energy confinement from the ohmic to the neutral beam heated phase of the discharge except at the highest plasma currents and lowest densities. For ohmic plasmas, τ_E scaled as $\bar{n}_e q^{1/2}$. It was found that, to a great extent, and especially for toroidal magnetic fields < 1.5 T, the global energy confinement during NBI was influenced predominantly by the electrons, with typically $> 70\%$ of the total energy lost through anomalous electron transport attributed to electron thermal conduction. Once degraded, the total thermal energy confinement only marginally decreased further with increasing beam power, consistent with the trend deduced from the magnetics measurements, which revealed only a weak scaling with power ($P^{-1/3}$) from ~ 1 to 6 MW of absorbed power.

Both the electron and total thermal energy confinement times were found to depend strongly on plasma current at fixed B_T , reflecting the increase in electron stored energy associated with broader electron temperature profiles at lower q . At fixed I_p , the electron heating was seen to increase with toroidal field, but τ_E at $r=0.6a$ showed only a marginal increase with increasing B_T . At fixed q , both τ_{Ee} ($r=0.6a$) and $\tau_E(a)$ improved with increasing toroidal field and plasma current. No meaningful density dependence of confinement could be inferred from the dataset because of the small range of densities sampled (2.5 to $4.2 \times 10^{13} \text{ cm}^{-3}$ for most discharges).

At higher toroidal fields (2.2 T), the ion channel influenced the total thermal confinement. For some of these cases as much as 25% of the input power was lost through bulk ion transport and charge exchange. Since ion thermal conduction was found to be well described by neoclassical theory, to within a factor of three owing to uncertainties in the $T_i(0)$ measurements, the improvement in τ_{Ei} with higher I_p and lower $\bar{n}_e$ was especially seen in the high field, high current discharges. This improvement led, in this analysis, to thermal confinements for these cases that were higher than those at comparable currents at lower toroidal fields ($I_p = 400 \text{ kA}$, $\bar{n}_e \approx 4 \times 10^{13} \text{ cm}^{-3}$). It has to be noted, however, that a discrepancy between the calculated and measured $2i/2 + \beta_{pol}$ signals exists at these high toroidal fields. Nevertheless, the electron conduction loss remained the largest single energy loss channel even for these discharges. Consequently, the variation of thermal confinement time was studied in terms of the variation of χ_e , the electron thermal diffusivity, with local plasma parameters.

The variations of thermal confinement time reflected the variations of χ_e in the outer region of the plasma ($r(q=1) < r < 7a/8$). Outside $q=1$, the variation of χ_e was tied to the variation of local poloidal field (or local

electron temperature), with χ_e decreasing with increasing B_p . A slight inverse correlation with local electron density near the plasma edge was also found.

The physical processes controlling χ_e and thus τ_E are still not known. Little agreement was found between the χ_e deduced from the PDX experiments and those predicted by various theories and scaling laws. In particular, the resistive ballooning mode formulation of Carreras and Diamond, which succeeded in explaining quantitatively the degradation of confinement with beam power for the high- β_{pol} ISX-B discharges, could not explain the variation of the deduced χ_e for the PDX cases. Results from multiple linear regression analyses indicate that any empirical scaling of χ_e on local plasma parameters should at least include an explicit dependence on poloidal magnetic field.

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APPENDIX 1

In Appendix 1 (Figs. 30 to 34) is shown comparison between the deduced electron thermal diffusivity, χ_e , and those predicted by theoretical or empirical scaling laws. The scalings are listed below either through references or the explicit expressions. The example shown is from a $R_T = 1.2$ T, $I_p = 250$ kA, three beam discharge. For this discharge, the position of the $q = 1$ and $q = 2$ surfaces are $r = 19$ and $r = 33$ cm, respectively.

- 1) χ_e - electron thermal diffusivity deduced from TRANSP
- 2) Drake - Drake et al. [78].
- 3) Kaw - Kaw [81].
- 4) Car-Dia - Carreras et al. [43].
- 5) Hir-Mmol - Hirshman and Molvig [77].
- 6) Cop-Maz - Coppi and Mazzucato [76].
- 7) Cop-Bon - Coppi et al. [82].
- 8) Ohkawa - Ohkawa [79].
- 9&10) PP1, PP2 - Parail and Pogutse [80].
- 11-13) TFR1, TFR2, TFR3 - Girard [83] for $q > 1$, $1 < q < 2$, $q < 2$, respectively.
- 14) Bohm - $ckT_e/16eB$.
- 15) Intor - $5 \times 10^{17}/n_e$ cm²/sec.
- 16) $1/nT$ - $4 \times 10^{20}/n_e T_e$ (eV) cm²/sec.

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FIGURE CAPTIONS

- FIG. 1 Unsmooth (solid line) and smooth (dashed line) Thomson scattering T_e (left panel) and n_e (right panel) profiles. The smoothing was done with the weighted average sliding filter technique described in the text. The fluctuations in the unsmoothed profile reflect the statistical uncertainty level of the Thomson scattering measurement.
- FIG. 2 Range of I_p and B_T covered by the PDX heating and β -scaling experiments. Circled points represent fully documented, controlled power scans, $R = 143$ cm, $40 < a < 44$ cm. The $a_{cyl} = 1.7$ line is drawn for $a = 42$ cm.
- FIG. 3 Typical output from the O-D, between shots analysis program based on the measured $\ell_i/2 + \beta_{pol}$ vs time. Shown is the output from a $B_T = 1.0T$, $I_p = 300$ kA, $a/r = 40/143$ cm, two beam discharge.
- FIG. 4 Time evolution of the T_e (left panel) and n_e (right panel) profiles, measured by the Thomson scattering system for a two beam $B_T = 1.0T$, $I_p = 300$ kA discharge.
- FIG. 5 Profile peakedness as a function of P_{TOT} ($P_{abs} + P_{oh}$) for n_e (upper panel) and T_e (lower panel) for 1.2 T discharges. The profile peakedness is defined as central value/volume-averaged value, where the volume-averaged electron temperature is density weighted.

- FIG. 6 Temperature profile peakedness $T_e(0)/\langle T_e \rangle$ as a function of a_{cyl} for $B_T = 0.7$ to 2.2 T, $I_p = 200$ to 480 kA, and for OH and two and four beam discharges.
- FIG. 7 Time evolution of the fits to the measured (solid line) and calculated (dashed line) thermal charge-exchange spectrum (top panel), and multiplier of the Chang-Hinton neoclassical ion thermal diffusivity (bottom panel) for a two beam, 300 kA, 1.0 T discharge. The neoclassical multiplier necessary to bring the measured and calculated fits into agreement was determined by a feedback loop within the code.
- FIG. 8 Comparison between the $\langle \beta_T^{dia} \rangle$ as measured from the diamagnetic flux inducing a voltage across the PDX toroidal field coils, and the value as calculated in the TRANSP kinetic transport code.
- FIG. 9a Volume integrated thermal ion power balance as a function of plasma radius for a 1.5 T, 300 kA, three beam case. 9b) Same as 9a for electrons.
- FIG. 10a Ion power balance volume integrated out to $r = a/2$ as a function of total heating power for 1.2 T, 300 kA discharges. 10b) Same as 10a for electrons.
- FIG. 11 Current scaling of the ion (left panel) and the electron (right panel) power balance volume integrated out to $r = 3a/4$ for ohmic, 1.2 T discharges.

- FIG. 12 Current scaling of the ion (left panel) and the electron (right panel) power balance volume integrated out to $r = a/2$ for two beam, 1.2 T discharges.
- FIG. 13 Current scaling of the ion (left panel) and the electron (right panel) power balance volume integrated out to $r = a/2$ for four beam, 1.2 T discharges.
- FIG. 14a Current scaling of ion power balance volume integrated to $r = a/2$ for 2.2 T, three beam discharges. 14b) Same as 14a for electrons.
- FIG. 15 Central ion temperatures, $T_i(0)$, derived from the ion power balance fit to the charge-exchange measurements, as a function of total heating power for four different currents at a fixed toroidal magnetic field of 1.2 T.
- FIG. 16 Central thermal ion energy confinement time, τ_{Ei} , at $r = 10$ cm ($a/4$) plotted as a function of η_i , the ion heating efficiency parameter, for all cases with the ion neoclassical diffusivity multiplier > 1 .
- FIG. 17 Ion heating, $\bar{n}_e \Delta T_i(0)$, as a function of absorbed beam power, P_{abs} , for four different currents at $B_T = 1.2$ T. The dashed lines define the values of the ion heating efficiency parameter, defined as $\eta_i = \bar{n}_e \Delta T_i(0)/P_{abs}$, in units of 10^{13} keV/(cm³-MW).

- FIG. 18 Central ion confinement time, τ_{Ei} , at $r = 10$ cm ($a/4$) plotted as a function of the total ion thermal diffusivity, χ_i , at the same radius. χ_i is the product of the Chanq-Hinton ion neoclassical thermal diffusivity and the multiplier needed to bring the fits to the measured and calculated thermal charge-exchange spectra into agreement. Shown are cases for which the neoclassical multiplier > 1 .
- FIG. 19a Electron energy confinement time, τ_{Ee} , at $r = 0.6a$ plotted as a function of total magnetic field for 2 to 2.5 MW discharges in the current range 200 to 320 kA. 19b) Same as 19a for the current range 350 to 480 kA.
- FIG. 20 The electron heating, $\bar{n}_e T_e(0)/P_{\text{heat}}$, plotted versus toroidal field for 2 to 2.5 MW discharges over two ranges of plasma current; 200 to 320 kA and 350 to 480 kA.
- FIG. 21a Electron energy confinement time τ_{Ee} at $r = 0.6a$ plotted as a function of total heating power, P_{heat} , for 1.2 T discharges at four different plasma currents. 21b) Same as 21a for 1.5 T discharges at $r = 0.6a$. 21c) Same as 21b for 2.2 T discharges at three different plasma currents.
- FIG. 22 Electron energy confinement times, τ_{Ee} , at $r = 3a/4$ plotted as a function at $\bar{n}_e q^{1/2}$ for ohmic discharges.

- FIG. 23 Total thermal energy confinement times, τ_E , at $r = a$ plotted as a function of toroidal magnetic field for 2 to 2.5 MW discharges in the current range 200 to 320 kA. 23b) Same as 23a for the current range 350 to 480 kA.
- FIG. 24a Total thermal energy confinement times, τ_E , at $r = a$ plotted as a function of total heating power, P_{heat} , for 1.2 T discharges at four different plasma currents. 24b) Same as 24a for 1.5 T discharges. 24c) Same as 24a for 2.2 T discharges at three different plasma currents.
- FIG. 25 Total thermal energy confinement time, τ_E , at $r = a$ plotted as a function of $\bar{n}_e q^{1/2}$ for ohmic discharges. A regression analysis of the points with $\bar{n}_e q^{1/2} < 5 \times 10^{13} \text{ cm}^{-3}$, assuming that τ_E is linearly proportional to $\bar{n}_e$, results in the fit $\tau_E \text{ (msec)} = 8.4 \times 10^{-13} \bar{n}_e \text{ (cm}^{-3}\text{)} q^{0.51}$ with a coefficient of variation of $R^2 = 0.91$.
- FIG. 26 The deduced electron thermal diffusivity, χ_e , plotted as a function of plasma radius for 1.2 T, 300 kA discharges. Shown are the χ_e profiles for the ohmic, one beam, two beam, and four beam cases. The dots represent the position of the $q = 1$ surface for each respective case. The total thermal energy confinement time for each case is given in the top left-hand corner of the graph.

FIG. 27 χ_e profiles for 1.2 T, two beam discharges for four different currents. The positions of the $q = 1$ surfaces are denoted by dots.

FIG. 28 χ_e profiles at fixed q ($q \approx 3$), for two beam discharges at various B_T . The dots denote the positions of the $q = 1$ surface for each case.

FIG. 29 Ratio of deduced χ_e to that predicted by the theory of Carreras et al. [43], for plasma parameters at $r = 30$ cm ($3a/4$) as a function of β_{pol} .

FIG. 30a Deduced χ_e values at $r = 20$ cm as a function of regression fit given for case # 1 in Table 2. Shown in the figure is the fit and the coefficient of variation, R^2 . 30b) Same as 30a at $r = 25$ cm for case # 3. 30c) Same as 30a at $r = 30$ cm for case # 5. 30d) Same as 30a at $r = 30$ cm for case # 6.

FIG. 31a Deduced χ_e values at all three radii simultaneously as a function of the regression fit given for case # 8 in Table 2. 31b) Same as 31a for case # 10. 31c) Same as 31a for case # 11.

FIG. 32-36 Comparison of the experimentally deduced electron thermal diffusivity, χ_e , with various theoretical and empirical models (see text for references) for a 250 kA, 1.2 T three beam discharge.

TABLE 1

B_T	I_p	$T_e(0)$	$\bar{n}_e T_e(0)$	$\tau_{Ee}(0.6a)$	$U_e(0.6a)$	$P_{heat,e}(0.6a)$
1.2 T	250 kA	0.92 keV	1.6 keV/cm ³ -MW	8 msec	9 kJ	1.23 MW
1.5	300	1.00	1.8	10 msec	12 kJ	1.17 MW
2.2	480	2.0	2.7	14 msec	15 kJ	1.11 MW

Measured and calculated electron parameters for comparison of two beam discharges at $q_{cyl} \approx 3$.

TABLE 2

	Case	10	r(cm)	B _p (T)	T _a (eV)	n _e (cm ⁻³)	q ²
r = 20 cm	1	5.14 ± 1.86		-1.45 ± 0.31		-0.13 ± 0.13	0.51
	2	3.40 ± 0.21		-1.53 ± 0.30			0.50
r = 25 cm	3	7.49 ± 1.51		-0.99 ± 0.24		-0.26 ± 0.11	0.51
	4	11.31 ± 1.27			-0.53 ± 0.13	-0.39 ± 0.10	0.62
r = 30 cm	5	15.34 ± 1.76		-0.83 ± 0.25		-0.82 ± 0.12	0.83
	6	19.48 ± 1.27			-0.40 ± 0.10	-1.02 ± 0.10	0.85
all radii	7	11.88 ± 1.42		-1.05 ± 0.23		-0.60 ± 0.10	0.54
	8	14.17 ± 0.95			-0.65 ± 0.07	-0.58 ± 0.07	0.72
(r=20,25,30cm)							
	9	12.92 ± 1.14	0.50 ± 0.26		-0.54 ± 0.10	-0.56 ± 0.07	0.73
	10	0.86 ± 0.29	1.94 ± 0.20	-1.56 ± 0.17			0.70
	11	6.57 ± 1.18	1.65 ± 0.18	-1.20 ± 0.16		-0.37 ± 0.08	0.77

Results of multiple linear regression fits of the χ_e deduced in the transport analysis code. Cases # 1 to # 6 represent the results of the regression fits of each specified radius, while cases # 7 to # 11 represent the results when data at all three radii are used simultaneously.

Figure 1

Figure 2

#84X0068

Figure 3

#84X0090

Figure 4

Figure 5

Figure 6

#84X0067

Figure 7

82 X 1043

Figure 8

#84 X0093

Figure 9

Figure 10

Figure 11

Figure 12

Figure 13

Figure 14

Figure 15

84X0061

Figure 16

Figure 17

Figure 18

Figure 19

Figure 20

Figure 21

#84X0209

Figure 22

Figure 23

Figure 24

Figure 25

Figure 26

#84X0056

Figure 27

#84X0063

Figure 28

Figure 29

Figure 30

Figure 31

Figure 32

Figure 33

Figure 34

Figure 35

Figure 36