Weak Temperature Dependence of the Thermal Diffusivity in High-Collisionality Regimes in LHD

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The strong positive dependence on the line-averaged electron density, $\bar{n}_e$, of the energy confinement time, $\tau_E$, as in the international stellarator scaling 1995 (ISS95) \cite{1}, where $\tau_E \propto \bar{n}_e^{0.51}$, disappears in the high-density regime in the Large Helical Device (LHD) \cite{2-4}. To investigate the physics behind this, it is important to clarify the parameter dependences of the thermal diffusivity. Here, we define an effective electron thermal diffusivity, $\chi_e^{\text{eff}}$, assuming that the density and temperature profiles of ions are equivalent to that of electrons and the heat conductions due to electrons and ions are also the same. When we plot $\chi_e^{\text{eff}}$ versus the electron temperature, $T_e$, at least two kinds of temperature dependences seem to appear, i.e., one is the gyro-Bohm type ($\chi_e^{\text{eff}} \propto T_e^{1.5}$) at the moderate density regime and another is the weak temperature dependence ($\chi_e^{\text{eff}} \propto T_e^{0.5}$) at the high-density (and therefore, low-temperature) regime \cite{4}. Other than $T_e$, $\chi_e^{\text{eff}}$ also depends on the magnetic field strength, $B_0$, which denotes the toroidal magnetic field strength on the magnetic axis. Systematic density and power scan experiments have been carried out at various $B_0$, to precisely investigate the $B_0$ dependence of $\chi_e^{\text{eff}}$ in the high-density regime (up to $<n_e> \sim 7 \times 10^{19} \text{ m}^{-3}$, where $<n_e>$ is the volume-averaged electron density). The magnetic configuration is fixed to $R_{ax} = 3.6$ m, where $R_{ax}$ is the major radius of the magnetic axis in the vacuum configuration. Negative-ion based neutral beam (NB) injection is applied and the total heating power is varied from 2 to 5 MW. The density is scanned by hydrogen gas puffing.

In Fig. 1, the relations between $\chi_e^{\text{eff}}$ and $T_e$ at $B_0 = 1.0 / 1.5 / 2.0$ T are depicted. The weak ($\chi_e^{\text{eff}} \propto T_e^{0.5}$) and gyro-Bohm like ($\chi_e^{\text{eff}} \propto T_e^{1.5}$) temperature dependences are recognized for $T_e$ dependence of $\chi_e^{\text{eff}}$ at various $B_0$. Solid and broken lines show the best fits for Eqs. (1) and (2). The inflection points are plotted by white crosses.
each dataset. To estimate the inflection point, $T_{e1}$, where the temperature dependence of $\chi_e^{\text{eff}}$ changes from $\chi_e^{\text{eff}} \propto T_e^{0.5}$ to $\chi_e^{\text{eff}} \propto T_e^{1.5}$, a set of models below is assumed:

$$\chi_e^{\text{eff}} = C_1 T_e^{0.5} / B_0^\alpha \quad (T_e \leq T_{e1}), \quad (1)$$  
$$\chi_e^{\text{eff}} = C_2 T_e^{1.5} / B_0^2 \quad (T_e > T_{e1}). \quad (2)$$

Below the inflection point, $\chi_e^{\text{eff}}$ increases with $C_1 T_e^{0.5}$ and decreases with an unknown $B_0$ dependence of an index $\alpha$. Note that $\alpha = 1.0 \pm 0.2$ was obtained in the former study [4]. Above the inflection point, we simply assume the gyro-Bohm model with a factor $C_2$. Three parameters of $C_1$, $C_2$ and $\alpha$ are determined at each normalized radius, $\rho (= r/a$, where $a$ is the averaged minor radius of the last closed flux surface), to give the minimum standard deviation, $\sigma$, of the experimental $\chi_e^{\text{eff}}$ compared with the model. As summarized in Fig. 2, the best solutions of $C_1$, $C_2$ and $\alpha$ are obtained at $0.6 \leq \rho \leq 0.85$, with small standard deviations of less than 2% (Fig. 2 (b)). From Fig. 2 (a), we conclude that $\alpha = 1.2 \pm 0.1$, which is consistent with the former result of $\alpha = 1.0 \pm 0.2$, while the ambiguity is reduced. $C_1$ is approximately constant (or, it slightly increases with $\rho$, at $\rho > 0.7$, see Fig. 2 (c)). This is also consistent with the former result. $C_2$ monotonically increases with $\rho$ (Fig. 2 (d)). The best set of $C_1$ and $C_2$ for a fixed $\alpha$ of 1.2, which are depicted by open squares in Fig. 2 (c) and (d), also gives reasonable fittings with small $\sigma$ (open squares in Fig. 2 (b)) and these are similar to that obtained for the best $\alpha$. The inflection point is calculated by $T_{e1} = (C_1/C_2) B_0^{-2\alpha}$. The radial profile of $T_{e1}$ is a decreasing function of $\rho$ (reflecting the $C_2$ profile) and increases with $B_0^{0.8 \pm 0.1}$ (Fig. 2 (e)).

If the thermal diffusivity can be expressed with experimental parameters, as in Eqs. (1) and (2), it is possible to construct a global confinement scaling. In the former study [4], a global confinement scaling is deduced from the relation of $\chi_e^{\text{eff}} \propto T_e^{0.5} / B_0$. Here, we reconsider it by

**Figure 2.** Summary of the fitting results. Closed symbols denote the results with the best $\alpha$ and open symbols denote the results with a fixed $\alpha$ of 1.2.
adopts $\chi_{e,eff} \propto T_e^{0.5}/B_0^{1.2}$, which well reproduces $\chi_{e,eff}$ in the outer region of $0.6 \leq \rho \leq 0.85$, as long as $T_e \leq T_{e1}$. To obtain a dimensionally correct expression, another dependence on the minor radius of $\chi_{e,eff} \propto a^{-0.5}$ should be introduced as below;

$$\chi_{e,eff} \propto (T_e/a)^{0.5}/B_0^{1.2}. \quad (3)$$

Assuming that $\tau_{E,mod} \propto a^2/\chi_{e,eff}$, Eq. (3) is transformed to

$$\tau_{E,mod} \propto P_{tot}^{-1/3} \times <n_e>^{1/3} B_0^{-4/5} a^{-7/3} R^{1/3}. \quad (4)$$

where $P_{tot}$ and $R$ denote the total heating power and the plasma major radius, respectively. This can also be expressed by non-dimensional parameters [5], i.e.

$$\tau_{E,mod} \propto \tau_{E,Bohm} \times a^0 \rho^{-0.8} \nu^{-0.3} \beta^{0.3}, \quad (5)$$

where $\tau_{E,Bohm}$, $\rho$, $\nu$, and $\beta$ are the Bohm confinement time, the ion gyro radius normalized by $a$, the collisionality, and the plasma beta, respectively. Since its non-dimensional form (Eq. (5)) is independent of $a$, Eq. (4) is dimensionally correct.

According to Eq. (4), the plasma stored energy should scale as

$$W_p^{HD} (kJ) = C \times <P_{dep}>^{2/3} \times <n_e>^{1/3} B_0^{4/5} a_{eff}^{-7/3} R_{ax}^{1/3}, \quad (6)$$

where units of $<P_{dep}>$, $<n_e>$, and $B_0$ are MW, $10^{19}$ m$^{-3}$, and T, respectively. To include the NB deposition profile effect, which becomes shallower in the high-density regime [2-4], the volume-average of the NB deposition profile, $<P_{dep}>$, is adopted as an index of the heating power. In our case, $<P_{dep}>/P_{tot}$ is $\sim 0.8$ at $<n_e> < 3 \times 10^{19}$ m$^{-3}$, and exponentially decreases with density in the higher density regime ($<P_{dep}>/P_{tot} \sim \exp(-<n_e>/12)$). Here, we define an effective minor radius, $a_{eff}$ (in m), by a product of $a$ and $\rho_{100eV}$, where $\rho_{100eV}$ is the average $\rho$ where $T_e$ ranges from 50 to 150 eV. This $T_e$ range is chosen because the reliability of our Thomson scattering system is assured at $T_e > 30$ eV. The effective minor radius is distributed within...
roughly ±5 % of 0.64 m, which corresponds to a in the vacuum configuration. As for the major radius, we adopt \( R_{\text{ax}} \) (in m), for simplicity. Due to the Shafranov-shift, the actual \( R \) increases to \( (R_{\text{ax}} + 0.1) \) m in the case of 1 % beta [6], for example. However, this has only negligible impact on the result (less than 1 (3) % for 0.1 (0.3) m of the Shafranov-shift).

The new scaling in Eq. (6) is compared with the datasets of the \( B_0 \) scan experiment in Fig. 3. The factor \( C = 140 \) is determined by the least square method using this data. Although the data with \( T_e > T_{e1} \) in the core region are also included in the figure, all of the data are well reproduced by \( W_{p}^{\text{HD}} \). This indicates the importance of the plasma property in the peripheral region, where \( T_e \) is relatively low and scarcely exceeds \( T_{e1} \) (note that \( \sigma \) rapidly increases at \( \rho > 0.8 \) and it becomes difficult to determine \( T_{e1} \), see Fig. 2 (b)), while the volume is large. In Fig. 4, an extended database is compared with the new scaling, where datasets of hydrogen gas-fueled plasmas obtained in the experimental campaigns in FY2003 (these are also used in Ref. 4) and FY2004 (including the \( B_0 \) scan experiment) are shown. Various experimental conditions are included in the database, i.e. \( B_0 = 1.0 / 1.5 / 2.0 / 2.75 \) T, \( P_{\text{NB}}^{\text{PT}} = 0.8 – 11.9 \) MW, \( <n_e> = (0.2 – 8.2) \times 10^{19} \) m\(^{-3}\). Again, \( W_{p}^{\text{HD}} \) is well reproduced by \( W_{p}^{\text{HD}} \).

It has been shown that \( Z_{e}^{\text{eff}} \) has a weak temperature dependence as \( Z_{e}^{\text{eff}} \propto (T_e/a)^{0.5} / B_0^{1.2\pm0.1} \), in the outer region of high-density LHD plasmas. The ambiguity in the \( B_0 \) dependence is reduced compared with the former study [4]. The inflection point, where the weak temperature dependence changes to the gyro-Bohm type ( \( \sim 1.2 \) keV at \( \rho = 0.6 \) and \( B_0 = 2 \) T, for example), increases with \( B_0^{0.8\pm0.1} \). Based on these observations and assuming a dimensional constraint, a new global confinement scaling that predicts the stored energy of high-density LHD plasmas at given experimental conditions has been obtained (\( W_{p}^{\text{HD}} \), in Eq. (6)). This new scaling matches well with the experiment. Compared with ISS95 (\( \tau_{E} \propto \bar{n}_{e}^{0.51} P_{\text{tot}}^{-0.59} \)) [1], our scaling has a weaker density dependence (\( \tau_{E} \propto <n_e>^{1/3} \)). However, it should be noted that our scaling is still favorable because the density dependence is positive and, especially, the power degradation is weak (\( \tau_{E} \propto <P_{\text{dep}}>^{-1/3} \).

References