Electron Pressure Profiles in High-Density Neutral Beam Heated Plasmas in the Large Helical Device

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In the Large Helical Device (LHD), electron pressure profiles in gas-fueled high-density discharges tend to have a similar shape, as if these were frozen. This frozen profile is insensitive to variations in the magnetic field strength and moderate changes in the neutral beam heat deposition profile. At the same time, however, the absolute value of the electron pressure itself increases with the heating power, the electron density, and the magnetic field strength. In this study, a reference model for the electron pressure is proposed which consists of the frozen profile and parametric dependences derived from experimental observations. It is possible to define an operational regime where this typical profile appears by comparing the electron pressure profiles with this model. In the standard configuration, at which the maximum plasma stored energy in LHD has been obtained, the frozen profile appears in the plateau to the Pfirsch-Schlüter regimes. As the collisionality decreases to the collisionless regime, the electron pressure becomes smaller than the prediction of the model and the deterioration is significant in the plasma core region. This tendency is enhanced in the configuration with the outward-shifted magnetic axis. The global energy confinement time, \( \tau_E \), in the high-collisionality regime has a weaker density dependence together with the mitigated power degradation, scaling as \( \tau_E \propto \bar{n}_e^{0.28} P^{-0.43} \) (\( \bar{n}_e \) and \( P \) are the line-averaged density and the heating power, respectively), compared with the International Stellarator Scaling 95, where \( \tau_E \propto \bar{n}_e^{0.51} P^{-0.59} \).

Keywords:
LHD, high-density, plateau regime, frozen profile, stiffness, energy confinement

1. Introduction

The confinement property of high-temperature plasmas in the high-density region is one of the important physics subjects in fusion reactor oriented studies, since high-density operation is favorable for achieving a higher fusion triple product. On the other hand, many examples of confinement degradation in the high-density region have been reported from both of tokamak and helical plasma experiments [1-4]. Use of empirical scaling laws for the energy confinement time, \( \tau_E \), is effective to identify the data deviating from the typical expectation, possibly due to a degradation or improvement of confinement. For example, the international stellarator scaling 95 (ISS95) [5] has been widely used to discuss the energy confinement of stellarators including the Large Helical Device (LHD) [6-9]. The energy confinement time predicted by the ISS95 scaling, \( \tau_{EISS95} \), has a favorable positive density dependence,

\[
\tau_{EISS95} = 0.079 a^{2.21} R^{0.65} P_{tot}^{-0.59} \bar{n}_e^{0.51} B_t^{0.83} \tau_{2/3}^{0.4},
\]

where \( a, R, P_{tot}, \bar{n}_e, B_t, \) and \( \tau_{2/3} \) denote the minor radius (m), major radius (m), total heating power (MW), line-averaged electron density (\( 10^{19} \) m\(^{-3} \)), magnetic field strength (T), and rotational transform at \( \rho = r/a = 2/3 \), respectively [5]. At a fixed heating power, therefore, the plasma-stored energy is expected to increase with the density. Indeed, this tendency is recognized in gas-fueled LHD plasmas in the moderate density region [10]. It is also observed, however, that the strong positive density dependence is lost in the high-density region during the density ramp-up by gas puffing. The energy confinement time in such a case is lower than the prediction of scalings with significant density dependence [3,4].

In this study, confinement properties of gas-fueled neutral beam (NB) heated LHD plasmas in the high-density region are investigated. It will be shown that the significant positive density dependence as given by ISS95 declines in the high-density region and that electron pressure profiles tend to have a similar shape. The typical profile, which is insensitive to variations in the experimental conditions, will be called a “frozen profile”. In such a case, the profile of the scale length of the electron pressure gradient converges to a...
typical curve. This is analogous to the stiffness observed in
tokamaks, where the temperature gradient scale length is
robust to the change in the heating deposition profile. This
inevitably occurs if the thermal diffusivity, $\chi$, has a nonlinear
positive dependence on the temperature, $T$, as $\chi \propto T^\alpha$ ($\alpha > 0$), for example, since the electron temperature scarcely
increases with the heating power due to the simultaneous
increase in the thermal diffusivity. In other words, profile
stiffness implies the non-linear dependence(s) of the thermal
diffusivity. Studies of profile stiffness have been intensively
pursued in tokamaks [1,2,11-16]. In these studies, stiffness in
the temperature profile has been mainly discussed in relation to
the ion temperature gradient mode [11-13] or the electron
temperature gradient mode [14-16]. So-called transient
transport analysis enables one to identify the nonlinear
temperature gradient dependence of $\chi$ [14-16]. The ratio
$\chi^{\text{HP}}/\chi^{\text{PB}}$, indicates that a strong temperature gradient
dependence appears ($\chi^{\text{HP}} \gg \chi^{\text{PB}}$) when the temperature
gradient exceeds a critical value [16]. The ratio $\chi^{\text{HP}}/\chi^{\text{PB}}$
is called the “stiffness parameter”, since such a nonlinear
dependence of $\chi$ can cause the stiffness. As for the
stellarators, on the other hand, studies in Wendelstein 7-AS
[18-21] have shown that no systematic discrepancy between
$\chi^{\text{HP}}$ and $\chi^{\text{PB}}$ is found [19,21] and therefore strong temperature
gradient dependence can be ruled out. The transient transport
analysis has been also applied to LHD plasmas with an
electron internal transport barrier (e-ITB) [22-24]. It has been
shown that the electron thermal diffusivity increases with $T^{-1/6}$
outside of the e-ITB [23], which is similar to the gyro-Bohm
model, while the dependence on the temperature gradient
seems negligibly small [24]. Note that the experimental
regime is limited to low-density in these studies, since low
collisionality is required to achieve the e-ITB [22]. In the
present study, the electron pressure profile is highlighted to
investigate the global energy confinement property in NB
heated plasmas in LHD. The energy confinement time is
readily obtained from the electron pressure profile, while
the temperature profile alone is not enough to discuss the global
energy confinement. However, the pressure gradient is mainly
determined by the temperature gradient in our case, since our
database is extracted from gas puff experiments where the
density gradient is small except in the plasma edge region. It
also should be noted that electron heating is dominant in
LHD, where high-energy neutral beam injection is adopted.

2. Apparatus

LHD is the largest super-conducting helical device in the
world [6-9]. The major radius of the toroidal vacuum vessel is
3.9 m, and the maximum magnetic field strength at the
plasma center is ~ 3 T. This study deals with gas-fueled
dehydrogen plasmas in a standard configuration of $R_{ax} = 3.6$ m,
or an outward-shifted configuration of $R_{ax} = 3.75$ m, where
$R_{ax}$ denotes the major radius of magnetic axis in the vacuum
configuration. Each of the magnetic configurations has unique
properties, e.g. neo-classical transport is reduced in $R_{ax} = 3.6$
m compared with the outward-shifted configurations [25] and
especially the trajectories of the high-energy trapped particles
are well aligned with the magnetic surfaces [26], while better
MHD stability is expected in the outward-shifted configurations that have a magnetic well [27]. Furthermore,
the largest confinement volume of ~ 30 m$^3$ is obtained in $R_{ax} = 3.6$ m and the maximum plasma stored energy of 1.3 MJ
has been achieved in this configuration [28].

The main heating system in LHD is negative-ion based
neutral beam (NB) injection, which consists of three beam
lines. The beam energy ranges from 140 to 180 keV, which is
much higher than the typical electron temperature of a few
keV in NB heated plasmas. Therefore, electron heating is
dominant (typically, 80–90 % of total heating power is
deposited to electrons). Although the heating power of each
beam line is fixed, combinations of the three beam lines
enable a power scan experiment, where the total port-through
power of NB, $P_{NB}$, is varied. The NB heating power, $P_{NB}$, is
estimated from $P_{NB}$ and direct heat load measurements of
NB shine-through power on the armor plates [29]. Note that
$P_{NB}$ is smaller than $P_{NB}$ especially in the low-density region
where the fraction of NB shine-through is large. Furthermore,
due to the deviation of the high-energy beam trajectory from
the magnetic surfaces, prompt loss of beam ions is expected
other than the beam shine-through, which is not included in
$P_{NB}$. This prompt loss is calculated by the FIT code [30] to
give an NB heat deposition profile, $P_{dep}(\rho)$, where $\rho = r/a$ is
the normalized radius and a is the averaged minor radius of
the last closed flux surface.

In this study, the confinement property of NB heated
LHD plasmas is discussed based on the analysis of electron
pressure profiles, $p_e(\rho)$. The $p_e(\rho)$ is given as a product of the
electron temperature, $T_e(\rho)$, and the electron density, $n_e(\rho)$,
where $T_e(\rho)$ is measured by Thomson scattering [31] and $n_e(\rho)$ is
estimated from Abel inversion of multi-chord far-infra-red
interferometer signals [32]. Here the projection to $\rho$ has been
performed using equilibria calculated by the VMEC code
[33].

3. Frozen Profile

To study the parameter dependence of the energy
confinement in the high density region, NB power scan
experiments have been performed, where $P_{NB}$ is scanned from 3 to 10 MW while the magnetic configuration is fixed
to $R_{ax} / B_0 = 3.6$ m / 1.5 T (Fig. 1). At each $P_{NB}$, $\tilde{n}_e$ is
scanned by gas puffing. The electron stored energy, $W_{e}^{\text{esp}}$, increases with $\tilde{n}_e$ (Fig. 1(b)) in the low-density region of $\tilde{n}_e < 3 \times 10^{19}$ m$^{-3}$. In this region, $W_{e}^{\text{esp}}$ agrees well with the trend of ISS95. In the high-density region, however, $W_{e}^{\text{esp}}$ stops
increasing (see hatched density region in Fig. 1) and even
decreases at higher density near the operational density limit
(beyond the hatched density region) [34,35]. The hatched
density region of $3.5 < \tilde{n}_e(10^{19}$ m$^{-3}) < 5.5$ in Fig. 1 is
arbitrarily determined since the threshold density, where the
strong positive density dependence of $W_{e}^{\text{esp}}$ at a fixed heating
power is lost, is ambiguous and the operational density limit
varies with the heating power. Although the strong positive density dependence is lost in this region, the electron pressure (Figs. 1(b) and 1(c)) still increase with the heating power. One might notice that \( W_{\text{exp}} \) plotted by the squares and the diamonds, which correspond to \( P_{\text{NB}} = 7 \) and 6 MW, respectively, are similar at \( n_e \sim 5 \times 10^{19} \text{ m}^{-3} \), in spite of the different heating power. This will be discussed later in section 4.1. To discuss the change in confinement property, it is important to know the electron pressure gradient, \( d\rho_e/d\rho \), and its scale length, \( L_p = (-d(\ln \rho_e)/d\rho)^{-1} \). In Fig. 1(d), shown are the density dependence of \( L_p^{-1} \) at \( \rho = 0.6 \) and 0.8. An important characteristic emerges from this figure, i.e. \( L_p \) is constant within \( \pm 10\% \) in the hatched density region, in spite of 4 times difference in \( P_{\text{NB}} \) (2.5 – 10 MW) and, correspondingly, 2 times difference in local \( p_e \). In the low-density region, \( L_p \) is not constant.

Fig. 1 Summary of the NB heating power scan experiment consisting of data from seven discharges, where (a) the NB heating power, \( P_{\text{NB}} \), (b) the electron stored energy, \( W_{\text{exp}} \), (c) the electron pressure, \( p_e \), at \( \rho = 0.8 \), and (d) the characteristic length of the electron pressure gradient, \( L_p^{-1} \), at \( \rho = 0.6 \) (open symbols) and \( \rho = 0.8 \) (closed symbols), are shown from top to bottom. Circles, squares, diamonds, and triangles denote \( P_{\text{NB}} \sim 10, 7, 6, \) and 3 MW, respectively. Note that \( P_{\text{NB}} \) in (a) is smaller than the corresponding \( P_{\text{NB}} \) especially in the low-density region due the large fraction of the beam shine-through power. The frozen profile appears in the hatched density region.

Radial profiles of \( L_p^{-1} \) in the high-density region, where the strong positive density dependence of \( \tau_E \) is lost, are shown in Fig. 2(a). The magnetic configuration is fixed to \( R_{\text{ax}} = 3.6 \) m, while \( B_0 \) is varied from 0.65 T to 2.75 T. Although each of these time slices has a different density \( (2 - 10 \times 10^{19} \text{ m}^{-3}) \), NB heating power \( (2 - 12 \text{ MW}) \), and NB deposition profiles as summarized in Table 1 and Fig. 2(b), the \( L_p \) profile scarcely changes (Fig. 2(a)). The collisionalities, \( \nu_{\text{ei}} \), in these time slices are high \( (0.17 \sim 1.07 \text{ at } \rho = 0.9, \text{ see Table 1} \) and in the plateau or the Pfirsh-Schlüter (P-S) regimes, where \( \nu_{\text{ei}} = \nu_{\text{ei}} R_{\text{ax}}/\left(\nu_T^2/(\nu T e(\nu_{\text{ei}})\right) \) (\( \nu_{\text{ei}}, \nu_T, \) and \( \nu_{\text{ei}} \text{ denote electron-ion collision frequency, electron thermal speed, and rotational transform, respectively.}

Fig. 2 Radial profiles of (a) the characteristic length of the electron pressure gradient at various magnetic field strengths, and (b) the NB heat deposition.
respectively, and \( v_p^* = 1 \) corresponds to the boundary between the plateau and the P-S regimes. In typical cases, \( v_p^* \) in the central region (\( \rho < 0.3 \)) is about one tenth of that at \( \rho = 0.9 \) and monotonically increases with \( \rho \). The conspicuous character observed in Fig. 1(d) can now be extended in that the \( L_p \) in the high-collisionality regime is insensitive to the magnetic field strength and the NB heat deposition profiles. Interestingly, a similar profile has been found in the high-collisionality regime using the following equations;

\[
\rho ( \rho ) = F \exp (-g(\rho)) ,
\]

\[
g(\rho) = 1.5 \rho^2 + 1.5 \rho^{10}.
\]

The factor \( F \) is primarily a function of the heating power, which will be determined in the next section. Here we call \( \exp(-g(\rho)) \) a frozen profile, which is depicted in Fig. 3. The volume integral of the frozen profile is defined as the base electron-stored energy, \( W_e^{\text{base}} \), i.e.

\[
W_e^{\text{base}} = \int_0^1 \exp(-g(\rho)) \cdot V' d\rho ,
\]

where \( V' = dV/d\rho \) is the radial derivative of the specific volume.

### 4. Parameter Dependences of the Model

#### 4.1 Heating power dependence

Even though the pressure profile is frozen, the electron pressure itself (and \( W_e^{\text{exp}} \)) increases with the NB heating power (see Fig. 1). A normalized electron stored energy, \( W_e^{\text{exp}} / W_e^{\text{base}} \), is fitted by \( f_p = 3.3 P_{\text{tot}}^{-0.55} \), as shown in Fig. 4(a), where \( P_{\text{tot}} \) (MW) = \( P_{\text{NB}} - dW_e^{\text{dia}}/dt \) and \( W_e^{\text{dia}} \) is the plasma stored energy estimated from diamagnetic signals. In our database, \( dW_e^{\text{dia}}/dt \) is small and less than 10 % of \( P_{\text{tot}} \) at most. Data extracted from the density region of \( 3.5 < n_e (10^{19} \text{ m}^{-3}) < 5.5 \), which corresponds to the hatched region in Fig. 1, are used to obtain the fitting function. Models of the electron pressure profile and the electron-stored energy (model A) are then given by;

\[
p_e^{\text{modelA}}(\rho) = f_p \exp(-g(\rho))
\]

\[
= 3.3 P_{\text{tot}}^{-0.55} \exp(-(1.5 \rho^2 + 1.5 \rho^{10})) ,
\]

\[
W_e^{\text{modelA}} = f_p W_e^{\text{base}} .
\]

The NB heating power deposition profile changes with the experimental condition, as was shown in Fig. 2(b). Apparently, a higher \( p_e \) profile is achieved in the central heating cases [3], although the profile shape, or the \( L_p \) profile, is frozen. Such an example is shown in Fig. 5, where two time slices at different densities are chosen from a discharge. These were also indicated in Fig. 2, as typical examples of the frozen \( L_p \) profiles in different experimental conditions. Broken lines show the profiles when the NB heat deposition is centrally peaked, while the gray solid lines show the higher density case where the NB heat deposition profile is flat. Since the heat flux in the core region should be more important than that at \( \rho = 1 \), we have adopted an averaged NB heat flux, \( Q_{\text{NB}}^{\text{avg}} \) (MW/m²), as an index of the NB heating power instead of \( P_{\text{tot}} \), where \( Q_{\text{NB}}^{\text{avg}} \) is the line-averaged value of the approximated NB heat flux profile defined as below.

\[
Q_{\text{NB}}(\rho) = P_{\text{avg}}^{\text{N}}(\rho)/(4 \pi \rho a R_{\text{oh}}) ,
\]

\[
P_{\text{avg}}^{\text{N}}(\rho) = \int_0^\rho P_{\text{avg}}(\rho') \cdot V' d\rho .
\]
Although the $P_{\text{dep}}^{V}$ at $\rho = 1$, which equals $P_{\text{tot}}$ as long as the prompt loss of beam ions and $dW_{e}/dt$ are negligible, is similar for both cases, higher $p_e$ is obtained in the central heating case, where $Q_{\text{NB}}$ is 30% higher at $0.3 < \rho < 0.8$ (the $L_p^{-1}$ profiles in both cases are similar to $g'(\rho)$ within ±10%). Here we adopt the line-average of $Q_{\text{NB}}$ as representative of the heating power for simplicity. It is possible to chose $Q_{\text{NB}}$ at a specific $\rho$, such as $\rho \sim 0.6$, which corresponds to the centroid of the volume integration of the electron pressure. However, line-average would be better since use of $Q_{\text{NB}}$ at a specific $\rho$ underestimates the heating power in the case of off-axis electron cyclotron heating outside of $\rho$, for example, although such plasmas are not included in this study. As long as the heat deposition profile is smooth as shown in Fig. 5(b), $Q_{\text{NB}}^{\text{avg}}$ is similar to $Q_{\text{NB}}$ at $\rho \sim 0.6$. As $\bar{n}_e$ increases, the NB heat deposition profile changes from centrally peaked to flat, which can be seen in the $\bar{n}_e$ dependence of $Q_{\text{NB}}^{\text{avg}}/P_{\text{dep}}^{V}(\rho = 1)$, as shown in Fig. 6.

Returning to Fig. 1(b), where $W_e^{\exp}$ of 200 kJ is obtained within ±10% at $3.5 < \bar{n}_e(10^{19} \text{ m}^{-3}) < 5.5$ in spite of 30% difference in $P_{\text{NB}}$ (squares for $P_{\text{NB}} \sim 7$ MW and diamonds for $P_{\text{NB}} \sim 5.5$ MW in the figure), these two have different NB heat deposition profiles but similar $Q_{\text{NB}}^{\text{avg}}$ (within 5%). Therefore, use of $Q_{\text{NB}}^{\text{avg}}$ instead of $P_{\text{tot}}$ is effective to reproduce $W_e^{\exp}$. As shown in Fig. 4(b), $W_e^{\exp}/W_e^{\text{base}}$ can be fitted with $42(Q_{\text{NB}}^{\text{avg}})^{0.57}$, in the density region of $3.5 < \bar{n}_e(10^{19} \text{ m}^{-3}) < 5.5$. Then we have model B as

$$p_e^{\text{modelB}}(\rho) = f_P^0 \exp(-g), \quad (9)$$

$$W_e^{\text{modelB}} = f_P^0 W_e^{\text{base}}, \quad (10)$$

$$f_P^0 = 42(Q_{\text{NB}}^{\text{avg}})^{0.57}. \quad (11)$$

Weak positive density dependence exists in $W_e^{\exp}/W_e^{\text{modelB}}$, as depicted in Fig. 4(c). This density dependence has the form of $f_n = 0.66 \bar{n}_e^{-0.28}$, in the density region of $3.5 < \bar{n}_e(10^{19} \text{ m}^{-3})$.
< 5.5. Using this, the third model is derived as

\[ P_{\text{e, modelC}}(\rho) = f_B f_{\text{Q}} P_Q \exp(-g) , \quad (12) \]

\[ W_{\text{e, modelC}} = f_B f_{\text{Q}} W_{\text{e, base}} . \quad (13) \]

These three models are composed of the frozen profiles. In other words, it is possible to determine the operational region where the frozen profile appears by comparing the electron pressure profiles with these models.

### 4.2 Magnetic field dependence

Models A, B, and C are compared with the experimental data of \( W_{\text{e, exp}} \) in Fig. 7, where datasets taken at \( B_0 = 1.5 \) and 2.75 T are shown (\( R_{\text{ax}} \) is fixed to 3.6 m). The abscissa is given by \( \nu_p^* \) at \( \rho = 0.9 \). At \( B_0 = 1.5 \) T, all models reproduce \( W_{\text{e, exp}} \) within ±10% in the high-density and high-collisionality region (hatched in Fig. 7). As the collisionality decreases (\( \nu_p^*(\rho = 0.9) < 0.1 \)), \( W_{\text{e, exp}} \) becomes smaller than the prediction of the models. This corresponds to the low-density region of \( n_e - 3 \times 10^{19} \text{ m}^{-3} \) in Fig. 1. Precisely speaking, model C slightly overestimates \( W_{\text{e, exp}} \) of \( B_0 = 1.5 \) T in the P-S regime (\( \nu_p^*(\rho = 0.9) > 1 \)). This might be due to the large fraction of the radiation loss, which increases up to ~30% of the total heating power in the same density region until the discharge is terminated by radiative collapse [35].

Qualitatively a similar tendency can also be recognized in the other dataset of \( B_0 = 2.75 \) T, although the ratio systematically increases with \( B_0 \). This suggests a hidden \( B_0 \) dependence in the models. The ratio \( W_{\text{e, exp}} / W_{\text{e, modelC}} \) in the high collisionality regime (\( \nu_p^*(\rho = 0.9) > 0.1 \)) is plotted against \( B_0 \) in Fig. 8. As shown in the figure, \( W_{\text{e, exp}} / W_{\text{e, modelC}} \) is fitted by \( f_B = 0.76 B_0^{0.74} \). A corrected version of model C is introduced as,

\[ p_{\text{e, modelD}}(\rho) = f_B f_{\text{Q}} P_Q \exp(-g) , \quad (14) \]

\[ W_{\text{e, modelD}} = f_B f_{\text{Q}} W_{\text{e, base}} . \quad (15) \]

This model well reproduces \( W_{\text{e, exp}} \) in the high collisionality regime at \( R_{\text{ax}} = 3.6 \) m, regardless of \( B_0 \). At different \( R_{\text{ax}} \), the density region in which \( \nu_p^*(\rho = 0.9) > 0.1 \) is not necessarily the same. For example, \( \nu_p^*(\rho = 0.9) = 0.1 \) corresponds to \( n_e \)
~ $3 \times 10^{19}$ m$^{-3}$ for $B_0 = 1.5$ T and $\bar{n}_e \sim 6 \times 10^{19}$ m$^{-3}$ for $B_0 = 2.75$ T, in Fig. 7. This suggests that the collisionality is more important than the density alone in determining the parameter regime where the frozen profile appears. This is why we adopted $\nu_p^e$ as an abscissa in Fig. 7 instead of the density. However, the reason why the profile becomes frozen in $\nu_p^e(\rho = 0.9) > 0.1$ is still an open question. The boundary between the plateau and the banana regimes lies at $\nu_p^e(\rho = 0.9) \sim 0.06$ in $R_{ax} = 3.6$ m and is slightly smaller than 0.1.

5. Outward-shifted Configuration

Up to this point, plasmas in the standard magnetic configuration of $R_{ax} = 3.6$ m have been analyzed. In this section, another dataset of the outward-shifted configuration, $R_{ax} = 3.75$ m, is introduced. The two configurations at the same $B_0 (= 1.5$ T) are compared in Fig. 9. In the low collisionality regime, $W_e^\exp/W_e^\text{modelD}$ at $R_{ax} = 3.75$ m decreases as the collisionality decreases. This qualitative behavior resembles that of the standard configuration, although the degradation is much enhanced. In the high collisionality regime, $W_e^\exp/W_e^\text{modelD}$ increases to ~0.8. In other words, the energy confinement time in $R_{ax} = 3.75$ m is ~20% smaller than that in $R_{ax} = 3.6$ m even in this regime.

Electron pressure profiles at various collisionalities are compared with model D in Fig. 10, where time slices A1 – A4 are chosen from $R_{ax} = 3.6$ m, and B1 – B4 are chosen from $R_{ax} = 3.75$ m (see also Fig. 9). In the low collisionality regime at $R_{ax} = 3.6$ m (A1), the electron pressure is lower than model D, especially at the plasma core. The electron pressure increases as the collisionality increases (A2). In the high-collisionality regime (A3 and A4), good agreement between $p_e$ and $p_e^\text{modelD}$ is obtained, and therefore the $p_e$ profile is frozen. Qualitatively a similar tendency is also observed in the outward-shifted configuration, where $p_e$ is smaller than $p_e^\text{modelD}$ in the low collisionality regime (B1 and B2) and increases as the collisionality increases (B3 and B4). Independent of $R_{ax}$, $p_e$ at the core is smaller than the prediction of model D in the low collisionality regime and increases as the collisionality increases, while $p_e$ at the edge is well reproduced by model D even in the low collisionality regime.

6. Discussions

It has been shown that the electron pressure profiles in LHD converge to the model profile in the high-collisionality regime. In the low-collisionality regime, the models overestimate the electron pressure profiles. There are two possible causes. One is that the profile is not yet frozen, and another is that the power and/or the density dependence in our models are not applicable in the low-collisionality regime. Radial profiles of $L_p$, $L_n = (-d\ln n_e/d\rho)^{-1}$, and $L_T = (-d\ln T_e/d\rho)^{-1}$, are shown in Fig. 11, where time slices of A1 – A4 and B1 – B4 are chosen as Figs. 9 and 10. In the case of $R_{ax} = 3.75$ m, the $L_p$ profile is not frozen at low collisionality (B1 – B2), as shown in Fig. 11(a). In the case of $R_{ax} = 3.6$ m, the $L_p$ profile seems to be frozen within the scatter of the data, even in the low collisionality regime (A1 – A2), where $W_e^\exp$ is lower than the prediction of model D. Therefore, the energy confinement time predicted by model D has the form of

$$\tau_e \propto \nu_e^{0.28} P_{abs}^{-0.43} B_0^{0.74},$$

where it is assumed that the beam deposition profile is fixed and the radiation loss is negligibly small. It should be noted that the energy confinement time in the low-collisionality regime is well reproduced by the ISS95 scaling, which has stronger positive density dependence and stronger power degradation ($\tau_e^{ISS95} \propto \nu_e^{0.51} P_{abs}^{-0.59} B_0^{0.83}$) than our model.

In this study, the confinement property of gas-fueled plasmas has been discussed. The density profiles sustained by gas puffing are basically flat except the edge region as was shown in Fig. 10. Indeed, the electron density gradient (Fig. 11(b)) is much smaller than that of the electron temperature gradient (Fig. 11(c)). Therefore, it is possible to attribute the frozen pressure profile to the property of the electron temperature gradient. To support this idea, it should be pointed out that even at low-collisionality where the $L_p$ profile is not frozen (B1 – B2), the $L_T$ profile seems to be frozen within the scatter of the data. However, the electron temperature profile alone is not enough to discuss the energy confinement property and it is necessary to estimate the local thermal diffusivity. After that, frozen temperature profiles can lead to a discussion of the temperature dependence of the thermal diffusivity, which is left for future study.

The frozen pressure profile appears in the low-temperature (high-collisionality) regime in LHD. At this point, our result is contrastive to the stiffness in L-mode tokamaks that appears in the high-temperature regime [16]. The electron temperature gradient mode (ETG) is one of the most probable
candidates causing stiffness in $L_T$ in tokamaks. In this scenario, $\chi$ increases with $T_e^{1.5}$ [36] (the collisionless trapped electron mode (TEM) also predicts a similar temperature dependence [37]), which is in good agreement with tokamak experiments [16]. The $T_e$ dependence of $\chi$ causing the frozen profile should be much weaker than this, however, to explain the weak density dependence in Eq. (16). Although it is unknown if the thermal diffusivity can be expressed by a single parameter, let us assume here that $\chi \propto T_e^\alpha$. Then the expected energy confinement time scales with

$$\tau_\text{E} \propto (n/P)^{\alpha/\alpha}$$,  \tag{17}$$

where $n$ and $P$ are the density and the heating power, respectively. Therefore, $\alpha \sim 0.39$ is necessary to obtain the density dependence in Eq. (16). On the other hand, the power dependence in Eq. (16) is reproduced by $\alpha \sim 0.75$. To reproduce both dependences simultaneously, $\alpha \sim 0.6 \pm 0.2$ might be the plausible temperature dependence of $\chi$. Anyhow, it is much weaker than the temperature dependences of $\alpha = 1.5$ in existing anomalous transport models such as the gyro-

Fig. 10 Radial profiles of the electron pressure, model D, and the electron density. Indices of (A1) – (A4) and (B1) – (B4) correspond to those in Fig. 9.
Bohm, ETG, and TEM, and even weaker than the Bohm model ($\alpha = 1$). Therefore, our result does not contradict the small stiffness factor ($\chi^{\text{bs}} \sim \chi^{\text{bs}}$), which rules out the strong temperature gradient dependence of $\chi$, found in former studies in stellarators [19,21].

7. Summary

In the standard configuration of $R_{sa} = 3.6$ m in LHD, the electron pressure profiles of gas-fueled NB heated plasmas dominated by electron heating can be reproduced within $\pm 10$% by the model equation of $p_e = f_{B} f_{n} f_{D}^{2} \exp (-\varepsilon(p))$, in the high-collisionality regimes ($\varepsilon_{p,6}(\rho = 0.9) > 0.1$). This model consists of the frozen profile, positive dependence on the heating power, the electron density, and the magnetic field strength. Use of $f_{D}^{2} = 42(Q_{NB}^{\alpha \varepsilon})^{0.57}$ is necessary to include the NB deposition profile effect. The global energy confinement time in the high-collisionality regimes has a weaker density dependence of $\tau_{E} \propto n_{e}^{-0.28}$ together with the mitigated power degradation of $\tau_{E} \propto P^{-0.43}$, compared with ISS95 scaling ($\tau_{E}^{\text{ISS95}} \propto n_{e}^{-0.51} P^{-0.59}$). The frozen profile discussed in this study is defined as having a gradient of the electron pressure (temperature) that is proportional to the local electron pressure (temperature) itself. This leads to an idea that the local heat diffusivity is a function of local parameter(s), such as the temperature or its gradient, which are closely connected to each other. Parameter dependences of the energy confinement time obtained here suggest a weak temperature dependence of the thermal diffusivity as $\chi \propto T_{e}^{0.6+0.2}$. As the collisionality decreases, the electron pressure profile becomes smaller than that predicted by the model. This degradation is significant at the plasma core region and enhanced in the outward-shifted configuration of $R_{sa} = 3.75$ m.

References

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