Ion Thermal Diffusivity in High Ion Temperature Mode in CHS Heliotron/Torsatron Plasmas

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1. Introduction

The advanced scenario of the wall conditioning and plasma fueling/heating in existing stellarator experiments resulted in achievements of new improved operation modes [1,2]. A high ion temperature mode (high T_i mode) has been observed for neutral beam heated plasmas in CHS Heliotron/torsatron [3]. The high T_i mode is characterized by a high central ion temperature (up to 1 keV) and is associated with a peaked electron density profile produced by neutral beam fueling with low wall recycling. The ion thermal diffusivity is significantly reduced (by a factor 2-3) in the high T_i mode plasma as compared to the L mode plasmas [3]. The strong magnetic field dependence of the central ion temperature and gradual increase of it with increase of the NB power (above the power of 0.35MW) has been observed.

It is an important issue to understand what causes the improvement of ion heat transport in the high T_i mode and in particular to compare the data with the theory prediction.

In this report the results of the neoclassical theory prediction for the ion thermal diffusivity for the high T_i mode plasmas of CHS are compared with experimental data. The ion thermal diffusivity was taken as a sum of tokamak-like symmetric term and asymmetric helical ripple term (electric field effects are not included). The helical ripple term was calculated on the basis of new approach developed in [4].

2. Helical ripple transport in CHS device

It is well known for stellarators that asymmetric neoclassical losses can exceed essentially symmetric losses at low collisionality regime (see, e.g., [5]). At present a number of analytical and numerical methods have been worked out to analyze this question. We used the method proposed in [4] for investigating the $1/\nu$ neoclassical transport. In this method the neoclassical transport coefficients are calculated by integration along the magnetic field line. Such integration takes into account the particles trapped within one magnetic field period and those that are trapped within several field periods as well. If the magnetic field is originally available in real-space coordinates calculations can be performed without a field transformation to magnetic coordinates. In accordance with [4] the energy flux density (averaged over a magnetic surface) can be presented in the form:

$$F_T = -\frac{\sqrt{8}}{9\rho^{3/2}} \frac{v_T^2 r_L^2}{nR^2} e_{eff}^{3/2} \int_0^\infty \frac{dz e^{-z} z^{7/2}}{A(z)} \frac{nT}{f_0} \frac{\nabla f_0}{r}, \quad (1)$$

$$e_{eff}^{3/2} = \frac{\rho R^2}{8\sqrt{2}} \lim_{L_s \rightarrow \infty} \left(\int_0^{L_s} \frac{ds}{B} \right) \left(\int_0^{L_s} \frac{ds}{B} |\nabla y| \right)^{-2} \int_{B_{min}^{(abs)}/B_0}^{B_{max}^{(abs)}/B_0} db' \sum_{j=1}^{j_{max}} \frac{\widehat{H}_j^2}{\widehat{I}_j}, \quad (2)$$

$$\hat{H}_j = \frac{1}{b'} \int_{s_j^{(min)}}^{s_j^{(max)}} \frac{ds}{B} \sqrt{b' - \frac{B}{B_0} \left(4 \frac{B_0}{B} - \frac{1}{b'} \right)} |\nabla \psi| k_G, \quad \hat{I}_j = \int_{s_j^{(min)}}^{s_j^{(max)}} \frac{ds}{B} \sqrt{1 - \frac{B}{B_0 b'}}. \quad (3)$$

Here, f_0 is the Maxwellian distribution (as a function of the particle energy w and the magnetic surfaces ψ), $v_T = \sqrt{2T/m}$ is the thermal velocity, $\rho_L = mc v_T / (eB_0)$ is the typical Larmor radius, $A(z)$ is the quantity related to the collision operator, R is the major radius of the torus, $\langle f_0 \rangle_r$ is the average normal derivative, $k_G = (\mathbf{h} \times (\mathbf{h} \cdot \nabla) \mathbf{h}) \cdot \nabla \psi / |\nabla \psi|$ is the geodesic curvature of the magnetic field line ($\mathbf{h} = \mathbf{B}/B$). The integral over z in (1) is general for any magnetic configuration and corresponds to the integration over w . The characteristic features of the specific magnetic field geometry manifest themselves through the factor $\mathcal{E}_{eff}$, where $\mathcal{E}_{eff}$ is the effective ripple modulation amplitude. This factor naturally takes into account contribution to the $1/v$ transport arising from all classes of trapped particles, i.e. particles trapped not only within one magnetic field period but also within several magnetic field periods.

The quantity $\mathcal{E}_{eff}$ is calculated by integration over magnetic field line length, s , (within the sufficiently large interval of $0 \div L_s$) and over the perpendicular adiabatic invariant of trapped particles (by means of the variable b' , $B_{min}^{(abs)}$ and $B_{max}^{(abs)}$ are the minimum and maximum values of B within the interval of $0 \div L_s$). The quantities $s_j^{(min)}$ and $s_j^{(max)}$ in the sum over j in (2) correspond to the turning points of trapped particles. The result (1) differs from the corresponding formula for the standard stellarator [5] by a simple replacement of the helical modulation amplitude $\mathcal{E}_h$ (along the magnetic field line) through the $\mathcal{E}_{eff}$ quantity. Therefore, for any magnetic configuration the $1/v$ transport coefficients can be obtained from the corresponding coefficients for the standard stellarator with the replacement of $\mathcal{E}_h$ by $\mathcal{E}_{eff}$. Formulas (1)-(3) are further applied to study the $1/v$ transport for two specific magnetic configurations of CHS: standard configuration corresponding to $R_{ax}=92.1$ cm and drift-orbit-optimized configuration ($R_{ax}=87.7$ cm)[6], where R_{ax} indicates the vacuum magnetic axis position in the $\varphi=0$ plane. The magnetic field produced by helical winding current was computed with the use of the Biot-Savart law. Calculations were performed over the interval L_s corresponding to 250 magnetic field periods. Figs. 1 and 2 illustrate results of these calculations for both configurations.

Fig.1 CHS standard configuration ($R_{ax}=0.921\text{m}$) (a-magnetic surfaces, b – $|\vec{B}|$,
c – rotational transform)

Fig.2 CHS drift-orbit-optimized configuration ($R_{ax}=0.877\text{m}$) (a-magnetic surfaces, b – $|\vec{B}|$,
c – rotational transform)

The results of ϵ_h and ϵ_{eff} computations for both configurations are shown in Fig.3.

Fig.3. Helical ripple in CHS

From the results it follows that for the standard configuration of CHS for the magnetic surfaces at moderate distances from the magnetic axis ($0.5 < r/a < 1$) the $\epsilon_{eff}^{3/2}$ value is greater (by a factor up to 2) than that for the standard stellarator. These results are used below for the comparison with experimental data. For the drift-orbit-optimized configuration the $\epsilon_{eff}^{3/2}$ value turns out to be approximately by factor of 10 less than for the standard stellarator. As pointed out in [6] this configuration is very close to the omnigenous configuration and must have improved confinement of locally trapped particles. Our results show that the $1/\nu$ neoclassical transport for this configuration is improved by approximately one order of magnitude in comparison with the standard stellarator.

3. Comparison of experimental data with theory prediction.

The high T_i mode for CHS was observed and studied up to now for standard magnetic field configuration. The ion heat transport analysis has been done at $t=90\text{ms}$ by using data of profile measurements (T_i – CX spectroscopy, n_e and T_e – YAG Thomson scattering).

The theoretical value of C_I was calculated as a sum of symmetrical term [5] and helical ripple term (formula (1)) by using T_i and n_e profile measurement data and results of calculation of ϵ_{eff} (Fig.3). This is justified because in the high T_i mode the ion collision frequency corresponds to plateau or low collisionality regimes of neoclassical theory. Fig.4 shows the radial distributions of normalized values of the ion collision frequency n_{ii}^* and characteristic frequencies of neoclassical theory for stellarator $n_b^* = e_t^{3/2}$ and $n_{sb}^* = (e_t e_h)^{3/2}$ ($n^* \equiv n/n_{ps}, n_{ps} = \frac{v_{Ti}}{R}$).

Fig.5 shows results of comparison of the theoretically predicted and experimental data for the high T_i mode discharge with $B=1.86\text{T}$. In this case in the most part of plasma ($r/a > 0.2$) experiment and theory are in a reasonable (within of factor 2) agreement.

Fig.4 Normalized ion collision frequency and characteristic frequencies of neoclassical theory for stellarator (standard configuration; $n_e(0)=2 \cdot 10^{13} \text{ cm}^{-3}$, $T_i(0)=0.84 \text{ keV}$)

Fig.5 Ion thermal conductivity for High T_i mode in CHS at $B=1.86 \text{ T}$
(1 – experiment, 2 – total theory predicted, 3 – symmetrical term, 4 – helical ripple term).

However at lower central ion temperatures in both High T_i and L-modes the experimental value of C_I is much larger than the theory predicted one (Fig.6)

Fig.6 Ion thermal conductivity in CHS at B=1.76T
(1- L-mode ($T_i(0)$)=390eV); 2 – High T_i mode ($T_i(0)$)=640eV))

The neoclassical theory also describes pretty well a general trends in dependence of the central ion temperature on heating power for high T_i mode at B=1.86T (Fig.6) (the 1/3 of NBI power P_{NBI} was used as the ion heating power P_{NBI-i} , on this Figure) and on confining magnetic field (Fig.7)

Fig.7 Central ion temperature vs. ion heating power

Fig.8 Central ion temperature as a function of confining magnetic field
(black points-experiment, red squares-calculation)

Experimental data on the central ion temperature as a function of magnetic field at $P_{\text{NBI}}=0.7\text{MW}$ are shown on Fig.7 (black points) together with the best fit (black line). Red squares show results of theory prediction together with the best fit - Gompertz relation (dashed red line) and linear fit (solid red line).

4. Discussion

Comparison of experimental data on the ion thermal conductivity in CHS device with the neoclassical theory predictions showed that the “best” high T_i mode discharges (with the highest ion temperature) exhibit nearly neoclassical behavior at the most part of plasma column. Apparent discrepancy near the plasma center (the theoretically predicted value is larger than experimental one) may be a result of large uncertainties in determining of ion temperature gradient in this part of column and underestimation of experimental value of C_i . The neoclassical theory described rather well the slow increase of the central ion temperature with NBI power ($T_i(0)\sim P^{0.264}$ as compared with experiment $T_i(0)\sim P^{0.25}$) and linear increase with confining field. It looks like that in the high T_i -mode CHS entered marginally in the regime of the neoclassical ion heat transport at central ion temperatures exceeding ~ 0.6 keV.

As it was pointed out in [3], the central ion temperature is sensitive to wall conditions. If the low wall recycling is attained and particle balance is sustained by NBI central fueling, the well-peaked density and NBI power deposition profiles are obtained. Altogether with improved ion heat confinement, this results in getting of well-peaked ion temperature profiles with high central temperature.

It seems that plasma-wall interaction still plays an important role in a small low-aspect ratio helical device CHS. One can suggest that the high T_i mode can be obtained much easier in the LHD.

The phenomenon of improvement of ion heat transport in the High T_i mode of CHS still needs better understanding. In this view it is interesting to achieve this mode for drift-orbit-optimized configuration of the CHS where the neoclassical theory predicts the central ion temperature up to 1.5 keV with the same NBI power as it was put in the standard configuration ($\geq 1\text{MW}$).

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REFERENCES

- [1] A.Iiyoshi, Fusion Energy (16th Conf. on Fusion Energy, Montreal, 1996) IAEA, Vienna Vol.1, 113 (1997)
- [2] M.Kick et al., Fusion Energy (16th Conf. on Fusion Energy, Montreal, 1996) IAEA, Vienna Vol.2, 27 (1997)
- [3] K.Ida et al., 17th IAEA Fusion Conference, Yokohama (1998), IAEA-CN-69/EX2/2
- [4] V.V.Nemov et al., 26th EPS Conf. on Contr. Fus. and Plasma Phys.1999, Maastricht P4.066; to be published in *Physics of Plasmas*.
- [5]. A.A. Galeev and R.Z. Sagdeev, in *Voprosy Teorii Plazmy* (Atomizdat, Moscow, 1973), No. 7, p.205 (in Russian); A.A. Galeev and R.Z. Sagdeev, *Reviews in Plasma Physics*, edited by M.A. Leontovich (Consultants Bureau, New York, 1979), Vol.7, p.257.
- [6] S.Okamura, K.Matsuoka et al., Annual Report of NIFS (April 1996-March 1997), 225 (1997)