Theory of Internal Transport Barrier of Helical Systems

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Abstract

Based on the electric field bifurcation model, a possibility to realize the internal transport barrier in helical plasmas is discussed. Condition to relaize the electric domain interface and reduction of turbulent transport is analyzed. Critical heat flux for inducing the bifurcation and transport barrier is obtained for the Heliotron/torsatron plasmas.

1. Introduction

In the high temperature plasmas which are confined in present helical devices, the anomalous energy transport dominates over the neoclassical energy transport. The improved energy confinement, if it exists, could be realized when the anomalous transport is suppressed. Based on the turbulent transport model, the improved confinement has been analyzed in the geometry of the Heliotron configuration [1]. The possible formation of internal transport barrier has been theoretically predicted [2]; the electric field domain interface [3] could be established, and transport reduction at the interface takes place. The electric field domain interface has been identified in CHS plasma, and the internal transport barrier has been found in the CHS plasma quite recently [4]. These developments of the confinement theory and experiment strongly suggest to explore the possible formation of internal transport barrier of the helical plasma. In this article, we study the electric field bifurcation and establishment of the internal transport barrier in helical plasmas. The critical heat flux, above which the interface and internal transport barrier appear, is obtained.

2. Flux and gradient

A theory of the strong turbulence in helical plasma has been developed based upon the method of self-sustained turbulence [5]. A theoretical formula of the turbulent transport coefficient is obtained as

$$\mu_N = \left(1 + h\,\omega_{E1}^2\right)^{-1} G_0^{3/2} s^{-2} c^2 \omega_p^{-2} v_A / qR$$

where $G_0 = a^2 \beta_p' d\ell n|B|/dr$ is the normalized pressure gradient which drives turbulence, $\omega_E = qRv_A^{-1}B^{-1}\, dE_r/dr$, $h = 1/2G_0$ and other notation is standard. This expression explains the L-mode transport in Heliotron plasmas. (See [1] for details.) The critical condition of the gradient of E_r to cause the transport barrier is written as $dE_r/dr \simeq \sqrt{G_0} Bv_A/qR$. This condition can be satisfied in the core plasma at the electric field interface.

In this article, we are interested in the transport barrier formation in the case of ECH plasma. For this case, the energy flow in the ion channel is smaller than that in the electron channel. For the analytic simplicity, we do not take into account the ion energy balance nor the ion temperature gradient. Under this simplified circumstance, we study the electron energy balance and particle balance. The particle flux and electron energy flux are given as

$$\Gamma_{tot} = \Gamma^{NC} - D_a n' \qquad (1)$$

$$q_{e,\,tot} = q_e^{NC} - \chi n T_e' \qquad (2)$$

where the index NC indicates the neoclassical component, and D_a and χ_a represent anomalous transport coefficients. In this article, off-diagonal elements in the anomalous transport matrix are not kept for the simplicity of the argument. In helical plasmas, the bipolar component of particle flux is dominated by the neoclassical flux. Radial electric field structure is governed by the nonlinear diffusion equation for the normalized electric field $X = eaE_r/T_e$ as

$$\nabla \cdot \mu_{eff} \nabla X = \Gamma_i^{NC} - \Gamma_e^{NC} \,. \tag{3}$$

The structure of the radial electric field, which is a consistent solution of Eqs.(1)-(3), is given as a function of the heat flux in this article. Equation (3) is approximated as

$$\Gamma_i^{NC} - \Gamma_e^{NC} = 0 \tag{3'}$$

except in the vicinity of the electric field interface. We use Eq.(3') and determine the location of the interface according to the Maxwell's construction [5].

The neoclassical flux associated with the helical-ripple trapped particles has been discussed in [6] and is simplified as

$$q_a^{NC} = \frac{n}{a} D_e \gamma_e T_e \left(-\frac{an'}{n} - X - \eta_{22} \frac{aT_e'}{T_e} \right) \tag{4}$$

$$\Gamma_e^{NC} = \frac{n}{a} D_e \left(-\frac{an'}{n} - X - \eta_{12} \frac{aT_e'}{T_e} \right) \tag{5}$$

$$\Gamma_i^{NC} = \frac{n}{a} D_i \frac{1}{1 + C_i X^2} \left(-\frac{an'}{n} + \frac{T_e}{T_i} X \right) \tag{6}$$

where coefficients γ_e, η_{22} and η_{12} are numerical constants of order unity, $C_j = 1.5(\varepsilon/\varepsilon_h)^{1/2}(T_e/\nu_j rBea)^2$ and $D_j = C\varepsilon^2\varepsilon_h^{1/2} v_D^2/\nu_j$.

We employ the normalized variables which represent gradients; $N = -\sqrt{C_i}\,an'/n$, $X = \sqrt{C_i}\,eaE_r/T_e$ and $Y = -\sqrt{C_i}\,\eta_{22} aT_e'/T_e$. By use of these normalized variables, Eqs. (1)-(3') are simplified as

$$\Gamma_t = \Gamma + d_{anom} N \tag{8}$$

$$Q_t = Q + \chi_{anom} Y \tag{9}$$

$$\frac{1}{1 + X^2} \left(N + \tau X \right) = \frac{1}{\xi} \left(N - X + bY \right) \tag{10}$$

where $d_{anom} = 2\sqrt{C_i}\,D_a/D_i(1 + \tau)$, $\chi_{anom} = \chi_a/\eta_{22} D_e \gamma_e$, $\Gamma = \sqrt{C_i}\left(nD_e\right)^{-1} a\Gamma_e^{NC}$, $Q = \sqrt{C_i}\left(nD_e\gamma_e T_e\right)^{-1} aq_e^{NC}$, $Q_t = \sqrt{C_i}(nD_e\gamma_e T_e)^{-1} aq_{e,\,tot}$, $\xi = D_i/D_e$, $\tau = T_e/T_i$ and $b = \eta_{21}/\eta_{22}$. Note that the ion energy balance is not taken into account. This is also because, for the parameter of interest, $D_e \simeq D_i$, (i.e., $T_e >> T_i$), the ion energy flux is much smaller than that of electrons.

3. Critical flux for transition

For the transparency of the argument, one may take a limit of a strong electron temperature gradient,

$$Y >> |N| . \tag{11}$$

This approximation is suitable, e.g., the case of electron cyclotron heating [4]; density gradient is often observed to be flat in such cases.

Gradient parameters, N, X, Y , are solved as functions of the normalized heat flux Q_t . When the heat flux is small, there is a solution with a weak radial electric field. As the heating power becomes stronger and Q_t increases, transition to a stronger electric field takes place. When the heat flux exceeds the critical heat flux,

$$Q_t > Q_{t,c} \quad , \tag{12}$$

the transition takes place. One obtains the critical heat flux as

$$Q_{t,c} = \left(\frac{\xi\tau}{2} + 1 \right) \frac{1 + \chi_{anom}}{b} - 1 \quad . \tag{13}$$

This transition is either hard or soft, depending on the parameters. One sees that Q_t is a monotonous increasing function of X, and the transition is a soft transition, if the condition

$$\left(1 - \frac{b}{1 + \chi_{anom}} \right) > \frac{\xi\tau}{8} \tag{14}$$

holds. On the contrary, in the large $\xi\tau$ limit,

$$\left(1 - \frac{b}{1 + \chi_{anom}} \right) < \frac{\xi\tau}{8} \quad , \tag{15}$$

the derivative, $\partial Q_t/\partial X$, changes its sign in the region of intermediate values of X. The electric field, as a function of the total heat flux, shows a hysteresis. The hard transition takes place at the critical condition. Figure 1 illustrates the electric field as a function of the heat flux, showing the soft and hard transitions.

As is discussed in §2, the turbulent transport could be suppressed at the interface,

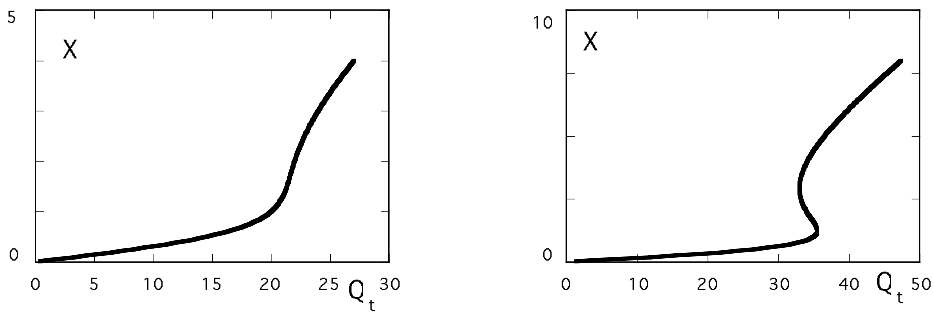

Fig.1 Radial electric field as a function of the heat flux (for $\chi_{anom} = 2$ and $b = 0.5$) . At a critical heat flux, the transition takes place. The case of soft transition ($\xi\tau = 5$, left) and that of hard transition ($\xi\tau = 10$, right).

where the condition $Q_t = Q_{t,\,c}$ is satisfied. Even in the region away from the interface, the neoclassical heat flux is reduced improving the energy confinement. The limiting relation between the temperature gradient and heat flux is also derived as

$$Y \simeq \left(1 - \frac{b}{\xi\tau + 1} + \chi_{anom}\right)^{-1} Q_t \qquad \text{(small } Q_t \text{ limit)} \qquad (16)$$

$$Y \simeq \left(\frac{1}{1 - b + \chi_{anom}}\right) Q_t \qquad \text{(large } Q_t \text{ limit)} \qquad (17)$$

Compared to the low flux limit, the derivative dY / dQ_t is higher in the high flux limit. This indicates that the energy confinement is improved in the high heat flux limit.

These analyses show that the transition of confinement takes place in the stellarator plasma. For a given flux, the electric field and temperature gradient could have multiple values if Eqs.(12) and (15) are satisfied. The transition from the lower electric branch to the higher electric field branch takes place at a critical heat flux. At the interface, $Q_t = Q_{t,\,c}$, the turbulent transport is suppressed. The result implies that the internal transport barrier can be established in helical plasmas. In the case that the heat deposition profile is strongly localized in the center. In this case, the heat flux is a decreasing function of the radius, except a central region where energy is absorbed. Consider the case that the condition $Q_t = Q_{t,\,c}$ is satisfied on one magnetic surface, $r = r_c$, the turbulent transport is suppressed at $r = r_c$, and the transport barrier is observed at $r = r_c$. Inside this surface, the neoclassical energy confinement is better than outside.

4. Summary and Discussion

In this article, we study the gradient of plasma parameters and radial electric field. It is shown that there exists a critical power flux above which the transition to the strong positive radial electric field is established. This explains the fact that the internal transport barrier appears if the heat flux exceeds a critical value.

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References

[1] K. Itoh, S.-I. Itoh, A. Fukuyama, H. Sanuki and M. Yagi.; Plasma Physics and Controlled Fusion **36** (1994), pp.123-129,
[2] K. Itoh, et al.; J. Plasma Fusion research Suppl. **74** (1998) 282. [In Japanese]
[3] S. -I. Itoh, et al., Phys. Rev. Lett. **67**, 2485 (1991). See a review, e.g., K. Itoh, S.-I. Itoh, Plasma Phys. Contr. Fusion 38 (1996) 1.
[4] A. Fujisawa, H. Iguchi, H. Sanuki, K. Itoh, et al.; Phys. Rev. Lett. **79** (1997) 1054, A. Fujisawa, H. Iguchi, T. Minami, Y. Yoshimura, H. Sanuki, K. Itoh, et al.; Phys. Rev. Lett. **82** (1999) 2669
[5] Itoh K, Itoh S-I, Fukuyama A; *Transport and Structural Formation in Plasmas* (IoP, Bristol, England, 1999).
[6] L. M. Kovrizhnykh; Nucl. Fusion **24** (1984) 435.