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Multi-scale transport simulation with ion temperature gradient driven drift wave turbulence

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Abstract

The gyro-fluid model proposed by Ottaviani et al. is extended, in which the ion parallel flow and ion neoclassical viscosity are taken into account. The heat source is also introduced. Using this model, the ion temperature gradient driven drift wave turbulence is analyzed. The reversed shear profile is employed, and internal transport barrier formation and collapse is investigated during the heating process. The multi-scale interaction between the thermal transport and the ITG turbulence is observed. In the saturation phase, the system approaches to marginal state, where the energy supply from source and intermittent transport with the avalanche process are balanced.

Key words: Gyrofluid, Gyrokinetic, Plasma turbulence, Profile relaxation, Heat source, Avalanche, ITG turbulence, Internal transport barrier, Intermittent transport, Reversed shear

1. Introduction

To achieve thermonuclear condition in a tokamak, it is necessary to confine the plasma for a sufficient time. Confinement is limited by thermal conduction and convection processes but radiation is also a source of energy loss. In the absence of instabilities, the confinement of tokamak plasma is theoretically determined by Coulomb collisions, so called neoclassical transport process. Unfortunately, the transport which occurs in reality does not agree with the values predicted by neoclassical theory. Especially, the thermal transport by electrons can be up to two orders of magnitude higher than predicted one.

It is considered that the observed anomalous transport is due to micro-instability of plasma, which causes particles and energy to escape at a higher rate [1]. Micro-Instabilities provide a mechanism for the generation of plasma turbulence and zonal flow, therefore understanding anomalous transport is crucial issue for ITER. Recently, a subclass of gradient-driven turbulence, the ion-temperature-gradient (ITG) driven turbulence has received considerable attention by theorists [2, 3]. This is partly due to the fact that, in neutral beam heated tokamaks, the main thermal loss occurs through the ion channel, therefore, ITG turbulence may play an important role in ion energy confinement. The transport studies in the framework of the international thermonuclear experimental reactor(ITER) have demonstrated that a machine of the ITER type would perform better if the present data can extrapolated according to the so called gyro-Bohm scaling (reviewed in Section 2.2). Several numerical studies are performed focusing on the $\rho_a$ scaling of ion thermal transport where $\rho_a$ is the ratio of the ion Larmor radius to a plasma minor radius. Early results from global gyro-kinetic codes pointed to a Bohm scaling (Section 2.2) for the ion thermal conductivity, however, the latest simulations which include the self-consistent poloidal flows address the gyro-Bohm scaling if the system size increases [4]. The similar study has been also performed based on gyro-fluid model which also supports the gyro-Bohm scaling [5]. These studies show that the zonal flow plays an important role for the saturation of ITG turbulence. It is shown that the zonal flow is not damped by collisionless process, but only damped by collisional process (reviewed in Section 2.4). Therefore, the realistic model
for neoclassical viscosity should be incorporated in the model, which contributes to the damping of the zonal flow. In this thesis, we extend the gyro-fluid model proposed by Ottaviani et al. [5]. In a new model, the effect of neoclassical transport is taken into account through the ion neoclassical viscosity, which includes the ion parallel flow and ion neoclassical viscosity. We analyze ITG mode and ITG turbulence by using this extended gyro-fluid model. We also perform transport simulation for the ITG turbulence, where the heating source is introduced in the system. The multi-scale interaction between transport (profile evolution), i.e., large scale, and turbulence (profile relaxation), i.e., microscale, is investigated in the presence of zonal flow, i.e., medium scale.

This thesis is organized as follows. In Chapter 2, ITG turbulence in tokamak plasma is reviewed. Theoretical approach of ITG turbulence is explained and results by numerical simulations are shown. In Chapter 3, we extend the gyro-fluid model to where we incorporate neoclassical viscosity which is derived by introducing the neoclassical MHD ordering [6], and equilibrium profile effects are also taken into account [5]. The transport simulation is performed using this model. In Chapter 4, we show the simulation results. Finally, in Chapter 5, summaries and discussions are given.

2. Reviews
2.1 Linear threshold of ITG mode

2.1.1 Linear analysis by drift-kinetic equation

We start the ITG mode theory from the drift-kinetic equation [7] with monovalent ion:

$$\frac{\partial f_i}{\partial t} + \frac{E \times B}{B^2} \cdot \nabla f_i + v_i \frac{\partial f_i}{\partial v_i} + \frac{e}{M} E_i \frac{\partial f_i}{\partial v_i} = 0. \tag{1}$$

where $f_i$ is the ion distribution function, $M$ is the ion mass, and the $z$-axis is taken to be parallel to the magnetic field $B$. The use of the drift-kinetic equation implies that we neglect the ion polarization drift as well as corrections of order $k_i p_i$ to the ion $E \times B$ drift. We assume the equilibrium distribution function for ion is Maxwellian given by

$$f_0(x,v) = n_0(x) \left( \frac{M}{2\pi T_{i0}(x)} \right)^{3/2} \exp \left( -\frac{Mv^2}{2T_{i0}(x)} \right). \tag{2}$$

Where $n_0$ and $T_{i0}$ is the ion equilibrium density and temperature, respectively. Here the electrostatic perturbation $E = -\nabla \nabla$ is considered. Linearizing the drift-kinetic equation and integrating over the velocity space, ion density perturbation $n_i$ is obtained as

$$n_i = -\frac{n_{oe} \Phi}{T_{i0}} + \frac{eB}{T_{i0}} \int_{-\infty}^{\infty} F_{0i}(v_z) dv_z \left\{ -k_i \sqrt{v_{ti}} \left[ 1 - \frac{\eta_i}{2} \left( 1 - \frac{v_{ti}^2}{v_{ti}^2} \right) \right] \right\}. \tag{3}$$

where an ion diamagnetic drift is given by

$$v_{di} = \frac{T_{i0}}{n_{oe} B_0} \frac{dn_{i0}}{dx} \tag{4}$$

and a dimensionless measure of the ion temperature gradient $\eta_i$ is given by

$$\eta_i = \frac{d\ln T_{i0}}{d\ln n_{i0}} \tag{5}$$

$(\omega, k)$ is the real frequency and the wave number of the mode, respectively. The one-dimensional Maxwellian distribution is also introduced as;

$$F_{0i}(v_z) = n_{i0} \left( \frac{M}{2\pi T_{i0}} \right)^{1/2} \exp \left( -\frac{Mv_z^2}{2T_{i0}} \right). \tag{6}$$

For electrons, we assume the adiabatic response:

$$n_{e1} = \frac{n_{oe} \Phi}{T_{e0}}. \tag{7}$$

By setting $n_{e1} = n_{ei}$, the dispersion relation is obtained as

$$1 + \frac{T_{e0}}{T_{i0}} = D(\omega) \tag{8}$$

where

$$D(\omega) = \frac{1}{n_{i0}} \int_{-\infty}^{\infty} F_{0i}(v_z) dv_z \left\{ -k_i \sqrt{v_{ti}} \left[ 1 - \frac{\eta_i}{2} \left( 1 - \frac{v_{ti}^2}{v_{ti}^2} \right) \right] \right\}. \tag{9}$$

Using the Nyquist diagram technique, it is found that the unstable drift wave exists for

$$\eta_i > 2 + \frac{4}{\Lambda} \frac{T_{i0}}{T_{e0}} \left( 1 + \frac{T_{e0}}{T_{i0}} \right) \tag{10}$$

where

$$\Lambda = \frac{n_{e1} k_i^2 v_{ti}^2}{k_i^2 v_{ti}^2}. \tag{11}$$

In the limit of $\Lambda \gg 1$, the condition for instability will approach to a limiting case of

$$\eta_i > 2. \tag{12}$$

The diamagnetic drift speed is generally much less than the ion thermal speed, the ratio $k_i/k_i$ must be exceedingly small to give $\Lambda \gg 1$. In the cases where arbitrarily small $k_i$ values are not allowed, such as in a torus geometry and in the presence of magnetic shear, the stability is modified. An inclusion of shorter wavelength modes is found to lower the instability threshold for $\eta_i$ to values close to unity.

2.1.2 Linear analysis of gyro-fluid equation

We have review the ITG mode theory based on the gyro-fluid equation which will be discussed in the chapter 3 in this thesis. Coupled equations for ion density,
parallel momentum and temperature evolution are linearized as,
\[
\frac{\partial}{\partial t}(\Phi - \rho^2 \nabla^2 F) + \kappa \frac{\partial \Phi}{\partial \theta} + A \nabla \Phi = c_0 \mu \Phi F
\]  
(13)
\[
dV = -A \nabla F
\]  
(14)
\[
\frac{3}{2} \left( \frac{\partial T}{\partial t} + \kappa \frac{\partial \Phi}{\partial \theta} \right) - \left( \frac{\partial n}{\partial t} + \kappa \frac{\partial n}{\partial \theta} \right) = \frac{2}{5} A \nabla V + \frac{1}{5} F
\]  
(15)
where \( F = \Phi + p/r \) is the generalized potential, \( p \) is the pressure, \( r = T_i/T_e \), \( \omega_a = 2 \cos \theta/r \theta_b + 2 \sin \theta \theta_b \) is the normalized \( \nabla \Phi \) drift frequency, \( A = e_\alpha / \mu_r \) is the inverse aspect ratio normalized by \( \mu_r \), \( \mu_r = p_\alpha / a \), is a normalized ion Larmor radius. \( \kappa_r = -d \ln T_0/\partial \) and \( \kappa_n = -d \ln n_0/\partial \) indicate gradient scale lengths of ion temperature and density.

In order to study the linear stability from Eqs.(13)-(15) in a toroidal geometry, we use the ballooning transformation [9-11]. We replace differential operator as
\[
\nabla \rightarrow \frac{1}{q} \frac{\partial}{\partial \phi}, \quad \nabla^2 \rightarrow -k_0^2(1 + \sin^2 \theta), \quad \frac{\partial}{\partial r} \rightarrow -i k_0 \theta \theta
\]  
(16)
Substituting these relations (16), Eqs.(13)-(15) are rewritten as
\[
-i \omega \Phi - i \omega \rho \omega \rho_\theta^2(1 + \sin^2 \theta) = k_0 n_\theta \kappa_\theta + \frac{A}{q} \frac{\partial V}{\partial \theta} = 2 e_\alpha k_\theta (\cos \theta + \sin \theta \sin \theta) F
\]  
(17)
\[
-i \omega V = -\frac{A}{q} \frac{\partial F}{\partial \theta}
\]  
(18)
\[
\frac{3}{2} \left( -i \omega T + i k_0 \kappa_T \Phi \right) - \left( -i \omega \Phi + i k_0 \kappa_\theta \Phi \right) = \frac{2}{5} \frac{\partial V}{\partial \theta}
\]  
(19)
If we neglect the RHS in Eq.(19) and eliminate \( V, T, \) and \( \Phi \) for simplicity, then the ballooning eigenvalue equation for ITG mode is obtained as,
\[
\frac{1}{A^2} \frac{\partial^2 F}{\partial \theta^2} + \left( \omega - \kappa_\kappa k_\theta \right)
\]  
\[
\frac{2}{q} + 1 + \frac{2}{q} \frac{2}{2} \left( \frac{3}{2} \kappa_T - \kappa_\theta \right)
\]  
\[
+ \omega^2 k_\theta^2(1 + \sin^2 \theta) + 2 e_\alpha k_\theta (\cos \theta + \sin \theta \sin \theta) F = 0
\]  
(20)
Taking the strong ballooning limit, i.e., the mode is localized to the region of \( \theta \simeq 0 \), we expand \( \cos \theta \rightarrow 1 - \frac{1}{2} \theta^2 \) and \( \sin \theta \rightarrow \theta \), and Eq.(20) reduces to the Weber type equation
\[
\frac{1}{A^2} \frac{\partial^2 F}{\partial \theta^2} + \left( X + \lambda \theta^2 \right) F = 0,
\]  
(21)
where
\[
X = \frac{\omega - \kappa_\kappa k_\theta}{\frac{2}{q} + 1 + \frac{2}{q} \frac{2}{2} \left( \frac{3}{2} \kappa_T - \kappa_\theta \right)} + \omega^2 k_\theta^2 + 2 e_\alpha k_\theta,
\]  
\[
\gamma = \omega \rho \rho_\theta^2 + 2 e_\alpha k_\theta \left( s - \frac{1}{2} \right).
\]  
The instability occurs if the condition \( D = b^2 - 4ac < 0 \) is satisfied, that is
\[
\kappa \simeq \frac{3}{2} \frac{2}{k_\theta} \kappa - \kappa > \frac{3}{2} \frac{2}{k_\theta} \frac{b^2}{16 e_\alpha k_\theta^2 a}
\]  
(22)
This gives,
\[
\tilde{\kappa} < \frac{3}{2} \frac{2}{k_\theta} \left( \frac{\sqrt{\omega a} + \sqrt{\omega a + k_\theta^2 p_\alpha^2 \kappa_\theta}}{k_\theta \rho_\theta^2} \right)^2 < \frac{3}{2} \frac{2}{k_\theta} \left( \frac{\sqrt{\omega a} + \sqrt{\omega a + k_\theta^2 p_\alpha^2 \kappa_\theta}}{k_\theta \rho_\theta^2} \right)^2
\]  
(23)
The linear growth rate is obtained as
\[
\gamma = \sqrt{\frac{4 e_\alpha k_\theta^2}{3 a}} \sqrt{\kappa - \frac{3}{2} \frac{2}{16 e_\alpha k_\theta^2 a}}
\]  
(24)
From Eq.(23), it is found that the radial profile of \( \tilde{\kappa} \) is determined by \( \kappa_n \). It should be noted the strong ballooning approximation is invalid in the vicinity of the minimum in the reversed shear configuration, i.e., the magnetic configuration that magnetic shear has negative value in the inner region and positive value in the outer region.
2.2 Bohm Scaling and Gyro-Bohm Scaling

If the turbulence is on a microscopic scale, for example on the scale-length of the ion Larmor radius $r_s$ measured at the electron temperature, the collisionless skin depth $c/\omega_{pe}$ with light speed $c$ and the electron plasma frequency $\omega_{pe}$, or resistive layer width $a/\sqrt{\tau_R/\tau_A}$ with the resistive diffusion time $\tau_R$ and the Alfvén time $\tau_A$, then the scale invariance technique can be applied to the equation for the fluctuations. This procedure leads to scalings for the fluctuations and the corresponding turbulent transport. For example, the potential fluctuation takes the form [12]

$$\frac{\phi}{T} = \frac{r_s}{L_n} f(\nu*, \beta, \cdots),$$

and the diffusion coefficient

$$D = D_{gb} f(\nu*, \beta, \cdots),$$

where the gyro-Bohm coefficient is

$$D_{gb} = \frac{\rho_s}{L_n} D_B,$$

with the Bohm coefficient $D_B = T/\epsilon_B$ and the density scale length $L_n = n/\rho_s$. And $f$ is the function of non-dimensional parameters such as the collision frequency $\nu^*$ and the plasma beta value $\beta$. On the other hand, if the fluctuations have a scale $l$ which is proportional to $a$, rather than $r_s$, but still satisfies $l \ll a$, then

$$\frac{\phi}{T} \sim \frac{l}{a},$$

and

$$D = D_{gb} f(\nu*, \beta, \cdots).$$

The calculation of specific forms for the function $f$ requires a model for the nonlinear saturation. A simple bound is given by the so-called mixing-length estimate, in which it is assumed that the instability drive is removed when the perturbations reach an amplitude such that the perturbed gradients equal the equilibrium gradient. Thus for drift waves driven by the density gradient, this estimate gives the density fluctuation $n_{k_1}$ of

$$k_L n_{k_1} \sim \frac{n_{k_0}}{L_n},$$

and since the density perturbations satisfy the Boltzmann relation, the saturated potential fluctuation is

$$\frac{\phi}{T} \sim \frac{1}{k_L L_n}.$$

The quasi-linear formula for electrostatic fluctuations is given by

$$\Gamma = \left( \sum_k \text{Im} \left( \frac{1}{\omega - k_L \nu_f} \right) k_L^2 \langle |\phi|^2 \rangle \frac{df}{d\omega} \right).$$

Inserting equation (2.19) into equation (2.20) gives

$$D \sim \frac{\gamma}{k_L^2},$$

where $\gamma$ is the imaginary part of $\omega$, which is the growth rate of the mode. This result can be conventionally interpreted as a balance between the linear growth of a mode and a stabilization from turbulent diffusion $k_L^2 D$.

2.3 Zonal flow formation and turbulence suppression

Zonal flow is driven by turbulence and is known to strongly influence the level of turbulence and transport. Zonal flows are generated by the Reynolds stress and can be considered as a nonlinear modulational instability associated with the inverse cascade of the turbulence energy. The plasma density and temperature fluctuations can generate mean poloidal flow (zonal flow), which is driven by the divergence of the wave energy density flux or, equivalently, the gradient of the Reynolds stress. The properties of the turbulence required for flow acceleration are that the fluctuations be radially propagating waves, and there be a radial asymmetry across the fluctuation spectrum. The kinetic energy of the generated flow is extracted from expansion of free energy stored in the fluctuations. As zonal flow speed and shear increase, collisional and turbulent viscosity and shear-enhanced decorrelation of fluctuations can be expected to self-consistently limit mean flow evolution.
From Figure 1, we find that the turbulence transport level is strongly influenced by zonal flow [13]. The ITG evolves from a linear phase of exponential growth to a nonlinear stage in which zonal flows are generated. When the effective shearing rate, or root mean square shearing rate increases, the ITG turbulence and transport are significantly reduced. Zonal flow is slowly damped by the ion-ion collisions and becomes weaker. When the effective shearing rate is below the growth rate, the ITG turbulence grows again and drives zonal flow. These turbulence-zonal-flow interactions, modulated by collisions, result in a cyclic, bursting behavior of fluctuations, transport, and zonal flow. In the phenomenon named GAM which is one of the zonal flow, the experimental results are different from the predicted data calculated by the past model, but new models are recently proposed [3,14-18], and the reason for the differences is explained. In the next section, the poloidal flow damping for collision and collisionless process is discussed.

2.4 Neoclassical flow dumping

Recent advances in gyro-fluid simulation of ion temperature gradient mode in tokamaks have shown that the predominant saturation mechanism for the instability is the production of axisymmetric, primarily poloidal flows which vary with radius and serve to shear stabilize to the instability. The damping of such poloidal flows is, thus, critically important in determining the turbulence level to be expected. First, collisionless process of poloidal flow damping is explained. From gyro-kinetic equation, the poloidal flow equation is derived by [19]:

$$u_p = (1 + 1.6q^2/\epsilon^{1/2})^{-1}u_p(0),$$

(34)

where \(u_p\) is the poloidal flow velocity, \(q\) is the safety factor and \(\epsilon\) is the inverse aspect ratio. It is shown that the poloidal flow is not damped by collisionless process. At least near the marginal stability, where the nonlinear damping of poloidal flows should be negligible, and in the sufficiently collisionless regimes of interest (deep in the banana regime), the level of poloidal rotation should be larger, and the ITG turbulence level and transport should be considerably smaller than predictions made by gyro-fluid simulations which entail linear collisionless damping. With respect to the collisional process of poloidal flow damping, it is shown that the poloidal flow is damped by collisional process [20, 21]. Figure 2 shows the dependence of diffusion coefficient on ion-ion collision frequency. As the collision frequency increases, the diffusion coefficient becomes larger, so that the zonal flow becomes weaker, the turbulence transport is enhanced and the plasma confinement becomes worse. The collisional poloidal flow is quite different from the collisionless poloidal flow.

2.5 Negative shear effect

There has been an increased interest in the negative shear region in tokamak plasmas. A significant improvement in plasma performance has been shown to be achievable in the plasma configuration with negative shear. In a tokamak, it is well-recognized that the negative shear profile can be obtained when the bootstrap-current takes a high fraction of total current.

In the linear analysis using simplified fluid limit model [22] of \(\omega_{Di}/\omega < 1, k_1v_i/\omega < 1 \text{ and } k_2^2\rho_i^2 < 1,\) where \(\omega_{Di}\) is a curvature drift frequency, the linear growth rate of the toroidal ion temperature gradient driven mode is roughly estimated in the regime of \(\eta_i \gg \eta_T, \) i.e., characteristic length of ion temperature varying is much larger than its threshold of stability, as

$$\gamma \sim \left[\omega_{Di}\omega_{pi} + \frac{\omega_{pi}v_i}{\eta_T} \left( s - \frac{1}{2} \right)^{1/2} \right],$$

(35)

where \(\omega_{pi} = \omega_i(1 + \eta_i), \omega_{Ti} = -ck_iT_i/eBL_n, \) \(s\) is the magnetic shear. For \(s > \frac{1}{2}\) the growth rate decreases with increasing \(s\) (note \(\omega_{Di}, \omega_{pi} < 0.\) On the other

![Fig. 2](image_url) Ion heat conductivity in nonlinear gyro-kinetic simulations with \(R/L_T = 5.3\) versus the ion-ion collision frequency [13].

![Fig. 3](image_url) The normalized growth rate as a function of \(s\) for the two cases of \(q = 1\) and \(4,\) when \(k_\rho T_i/\eta_i = 0.5, \eta_i = 4, \epsilon_n = 0.2, \) and \(T_i/T_e = 1.\) [22].
hand, for $s < \frac{1}{2}$ the growth rate has the form
\[
\gamma \simeq \sqrt{a(1 + i\epsilon)} \\
\sim \sqrt{b(1 + \frac{\epsilon}{2} - \frac{3}{8} \epsilon^2 + \cdots)},
\]
with $a = \omega_{DIW}'$ and $\epsilon = \frac{v}{2(1/2 - s)}B_{DIW}$. The growth rate becomes smaller with decreasing $s$ by the third expansion term. (The second expansion term $i\epsilon/2$ just changes the real frequency.) Thus we find that the growth rate of the toroidal ITG mode is maximized around $s \approx 1/2$, and becomes weaker with increasing $q$, as shown in Figure 3. The negative shear may give a substantial stabilizing effect when $q$ is small.

### 2.6 Internal Transport Barrier formation

The internal transport barrier (ITB) which makes an improved confinement has been observed in many tokamaks with a reversed-magnetic-shear (RS) configuration. ITB is beneficial for the steady-state operation of ITER and advanced tokamak reactors, and RS plasma is one of the attractive candidates for the advanced tokamak reactor because of its compatibility with the large bootstrap current produced by ITB. However, the physical mechanism of the ITB structure in the RS plasma has not been fully understood yet. The improved performance with ITB was surveyed in multi-machine comparisons [23]. On the basis of JT-60U experimental database of the box-type ITB, two scalings for the narrow ITB width [24] and the energy confinement inside ITB [25] have been developed. The former scaling shows that the ITB width is proportional to the ion poloidal gyro-radius at the ITB center. The latter scaling is equivalent to a condition for the core poloidal beta, $\beta_{p,core}$, i.e. $\epsilon_f/\beta_{p} \approx 0.25$, where $\epsilon_f$ is the inverse aspect ratio at the ITB foot. Data is classified into three groups: narrow ITB width $\Delta_{ITB} \propto \rho_{p,ITB}$ with $\rho_f/\rho_{p,ITB} \leq 1$ and large $\chi_{GB}$, wide ITB width $\Delta_{ITB}$ with $\rho_f/\rho_{p,ITB} \leq 1$ and small $\chi_{GB}$ ($\triangle$) and wide ITB width $\Delta_{ITB}$ with $\rho_f/\rho_{p,ITB} > 1$ ($\bigcirc$). Where $\Delta_{ITB} = \rho_f - \rho_{sh}$, $\rho_f$ is the ITB foot position on the normalized minor radius, and $\rho_{sh}$ is the ITB shoulder position inside which the $T_i$ profile becomes flat. $\rho_{p,ITB}$ is the ion poloidal gyro-radius, $\rho_{p,min}$ is a position of the minimum safety factor and $\chi_{GB}$ is defined as $\chi_{GB} \propto T_{i0}^{0.5}/(B_{p,f}^{2} \Delta_{ITB})$, where $T_{i0}$ is $T_i$ at $\rho \sim 0$ and $B_{p,f}$ is the poloidal magnetic field at the outer midplane ITB foot. These experimental data and the physical mechanism of strong ITB structures has been verified and studied by Hayashi et al. [26], through the modeling on the 1.5 dimensional time-dependent transport simulation. For the scaling for the narrow ITB width ($\Delta_{ITB} \propto \rho_{p,ITB}$), the following three physics are reported to be important: (1) the sharp reduction of the anomalous transport below the neoclassical level in the RS region, which results in $\rho_f/\rho_{p,min} \sim 1$, (2) the autonomous formation of pressure and current profiles through the neoclassical transport and the bootstrap current, which results in the strong RS configuration and (3) the large difference between the neoclassical transport and the autonomous transport in the normal-shear region, which affects the autonomous formation. The value of $\epsilon_f/\beta_{p,core}$ reaches to the saturation value, i.e. $\epsilon_f/\beta_{p,core} \approx 0.25$, when the box-type ITB is formed in the strong RS plasma with a large asymmetry of the poloidal magnetic field regardless of the details of the transport and the non-inductive current.

### 3. Model

#### 3.1 Global simulation of ITG turbulence

In order to investigate the ITG turbulence and transport, we extend the gyro-fluid model equations proposed by Ottaviani et al. [5]. We consider a high temperature plasma with the major radius $R$, the minor radius $r$, immersed in a toroidal magnetic field $B_0$ and use the toroidal coordinate system $(r, \theta, \phi)$. Our model is composed of three conservation equations for the ion density, the parallel momentum and the energy as follows:

**Ion continuity equation**

\[
\frac{dW}{dt} + \nabla \cdot (\rho_{\epsilon} \mathbf{v} \nabla U) = \epsilon \mathbf{F} + \rho_s \frac{g}{2} \mu^{NC} \frac{\partial U_E}{\partial r} - \rho_s^2 \mu \nabla_\bot F
\]

**Ion parallel momentum equation**

\[
\frac{dV}{dt} = -\nabla_\| F + 4\mu \nabla_\| V - \mu^{NC} U_p + A \frac{1}{2} \frac{1}{r^2} \| \mathbf{k}_i \| V + \frac{2}{5} A \frac{1}{r^2} \| \mathbf{k}_i \| T
\]
Ion energy equation
\[
\frac{3}{2} \left( \frac{dF}{dt} + \kappa_T \frac{1}{r} \frac{\partial F}{\partial \Phi} \right) - \left( \frac{dn}{dt} + \kappa_n \frac{1}{r} \frac{\partial n}{\partial \Phi} \right) = -\frac{9}{5\sqrt{\pi}} \frac{1}{r} A \left| V_{||} \right| T + \frac{2}{5} A \nabla_{||} V + \chi_{n} V_{\perp} T
\]
where, \( W = n - \nabla_{\perp}^2 F \), \( F = \Phi + \frac{p}{\tau} \) is generalized potential, \( p = n + T \) is pressure. We assumed that density evolves according to the Boltzmann response, except the zonal flow component which satisfies
\[
\frac{\partial n}{\partial t} = \Phi - \bar{\Phi}
\]
where \( \bar{\Phi} \) represents the flux surface averaged potential, in the Fourier space it is equivalent to \( \Phi_{m=0, n=0} \), therefore, \( n_{m=0, n=0} = 0 \). This condition means that, there is no relaxation of ion density profile. In other words, particle flux to the \( r \) direction does not exist. Here \( U_p = V + \rho_s \frac{\partial F}{\partial \Phi} \) is the poloidal velocity, \( \kappa_T = -\frac{d}{dr} \ln T_{\tau}/\frac{dr}{dr} \) and \( \kappa_n = -\frac{d}{dr} \ln n_{\tau}/\frac{dr}{dr} \) indicate gradient scale lengths of ion temperature and density. \( \omega_d = 2 \cos \theta \frac{\partial}{\partial \Phi} + 2 \sin \theta \frac{\partial}{\partial \Phi} \) is the normalized \( \nabla \Phi \) (including contribution of toroidal curvature) drift frequency, which drives toroidal ITG mode. \( V_{\perp} = \frac{1}{2} \left( \frac{\partial}{\partial \Phi} + \frac{\partial}{\partial \Phi} \right) \) is the parallel derivative operator. \( \varepsilon = \rho_s / \rho \), \( \varepsilon_a = \rho_a / \rho \), the inverse aspect ratio, \( \varepsilon = r/R_0 \), the local inverse aspect ratio and \( \rho_s = \rho_s / a \) denotes the normalized ion gyro-radius. Note that molecular transport coefficients are introduced, i.e., thermal diffusion \( \chi \), neoclassical viscosity for poloidal motion \( \mu^{NC} \) and the ion viscosity \( \mu \).
Ion Landau damping is also taken into account as a collisionless dissipation \[29\]. This is corresponding to the 4th term of R.H.S. in Eq.(37) and the 1st term of R.H.S. in Eq.(38). This normalization is taken to be
\[
\frac{t}{t_B} \rightarrow t \quad , \quad \frac{r}{r_B} \rightarrow r \quad , \quad \frac{z}{z_B} \rightarrow z \quad , \\
\frac{\Phi}{\Phi_B} \rightarrow \Phi \quad , \quad \frac{V}{V_B} \rightarrow V \quad , \quad \frac{T}{T_B} \rightarrow T \quad , \\
\frac{\chi}{\rho_s c_s} \rightarrow \chi \quad , \quad \frac{\mu}{m_i n_0 \rho_s c_s} \rightarrow \mu \quad , \quad \frac{\mu^{NC} c_a^2}{\rho_s c_s} \rightarrow \mu^{NC}
\]
where,
\[
t_B = \frac{a^2}{\chi_B} \quad , \quad \chi_B = \frac{c_T e a}{e B}
\]
i.e., the characteristic time for the perpendicular dynamics is chosen to be the Bohm time \( t_B = \rho_s^{-1} a / c_s \). This model is a simplified version of so-called 3+1 gyrofluid model \[30\], which is applicable to the meso-scale dynamics. Meso-scale is the intermediate scale length between large scale and microscale. The main difference between reference \[5\] and our model is the dissipation form. The fluctuating ion poloidal flow is introduced according to the neoclassical MHD ordering \[6\]. Our model ensures the positive definiteness of dissipations, i.e.,
\[
\frac{dH}{dt} = -\frac{3}{2} \kappa_T \left( \frac{1}{\tau} \frac{\partial F}{\partial \Phi} \right) - \mu^{NC} (U_p^2) - \rho_s^2 \mu (\nabla_{\perp}^2 F)^2
\]
\[
-4\mu (\nabla_{\perp} V^2) - \frac{\chi}{\tau} (\nabla_{\perp} T^2)
\]
with
\[
H = \frac{1}{2} \rho_s^2 (\nabla_{\perp} F)^2 + \frac{1}{2} (V^2) + \frac{1}{2} (\tau) (\phi^2) + \frac{3}{4} (T^2)
\]
where \( \langle \cdot \rangle \) indicates the volume average.
Other parameters for the molecular viscosity and thermal diffusivity are chosen as \( \mu = \chi = 1.0 \times 10^{-3} \) and neoclassical viscosity is \( \mu^{NC} = 1.0 \). These values allow a stationary state of turbulence without a significant impact on the level of turbulence.

### 3.2 Source term

We introduce a heat source term into the model equation in this study, i.e., inhomogeneous term of ion thermal conduction equation. The source is added in the ion energy equation of Eq.(38). In the Fourier space, it is included in the form of \( T_{m=0, n=0} \). Thereby we can investigate the interaction between the heat transport and the ITG turbulence itself during the heating process.
We consider two types of the radial profile for the heat source. The first type of source \( S_1 \) is given by,
\[
S_1(r) = -8.0 \times 10^{-3} \left[ \frac{2r^2 - 1}{(1 - r^2)^2} \right]
\]
In the case without a convective nonlinearity, this source gives the condition that the ion temperature profile asymptotically reaches to \( T_\infty(r) = \frac{1 - r^2}{(1 - r^2)^2} \), as \( r \rightarrow \infty \), shown in Figure 5. Here, we set \( r_\text{r} = 0.6 \).
This source profile is negative, namely, it assumes the existence of heat sink in the outside of the plasma, as is shown in Figure 5.
We also investigate a Gaussian shape of source \( S_2 \), which is positive all over the plasma column, i.e.,
\[
S_2(r) = \frac{8.0 \times 10^{-3}}{(1 - r^2)^2} \exp \left[ -\frac{r^2}{0.273316} \right]
\]
The parameters of \( S_1 \) and \( S_2 \) are determined such that, both source terms have the same amplitude at the center, and supply the same amount of energy per unit time. Here, the density profile is taken to be the same for both cases, that is,
\[
n_0(r) = \frac{1}{1 - r^2}
\]
shown in bottom of Figure 5. The source term \( S_2 \) gives gentle temperature profile asymptotically at \( r \rightarrow \infty \) as shown in Figure 5.
3.3 Reversed shear profile

For the choice of safety factor $q$ profile, we employ reversed magnetic shear profile, proposed by Garbet et al. [28].

$$q(r) = q_{\text{min}} + C_2 (r^2 - r_{\text{min}})^2 + C_3 (r^2 - r_{\text{min}}^2)^3$$

where, $r_{\text{min}} = 0.6$, $q_{\text{min}} = 1.35$, $C_2 = 4.66$, $C_3 = -0.987$. The magnetic shear is null at $r = 0.6$. The profile is shown in Figure 6.

3.4 Numerical settings

For the numerical studies, we use the cylindrical co-ordinate $(r, \theta, z)$ and discretize all variables in the $r$ direction and expand in Fourier series, $\exp(i(m\theta - nz))$ in $\theta$ and $z$ directions, respectively. As the Fourier mode, resonant modes and their sideband modes are selected. Left hand side of Figure 7 shows Fourier spectrum in wave number $(m,n)$ space which we select by blue and red points. The distribution of rational surfaces of them is shown in right hand side of Figure 7. We put damping zones in $r < 0.1$ and $r > 0.9$. In addition, the Fourier components colored by red in left hand side of Figure 7 are disposed at every prediction steps in ion continuity equation and ion parallel momentum equation. The number of Fourier modes which are retained in the computation is 2179, and the number of cell for $r$ direction is 256. Our model equations are schematically written as

$$\frac{\partial f_{m,n}(r)}{\partial t} = \mathcal{L}(f_{m,n}(r)) + \sum_{m' = m, n' = n} \mathcal{N}(f_{m',n';r}, f_{m',n';r}(r)) + \sum_{m' = m} \mathcal{B}(f_{m+m'}(r))$$

where, $\mathcal{L}$ represents the linear operator, $\mathcal{N}$ is nonlinear operator, and $\mathcal{B}$ is toroidal coupling operator. We solve the linear part by the implicit Cranc-Nicolson scheme using the LU decomposition. Nonlinear and toroidal coupling parts are integrated with the predictor-corrector scheme, i.e.,

the predictor step:

$$f^{t+\Delta t} = \left( I - \frac{\Delta t \mathcal{L}}{2} \right) \left( I + \frac{\Delta t \mathcal{L}}{2} \right) f^t + \left( I - \frac{\Delta t \mathcal{L}}{2} \right)^{-1} \Delta t \mathcal{N}(f^t, f^t) + \mathcal{B}(f^t)$$

the corrector step:

$$f^{t+\Delta t} = \left( I - \frac{\Delta t \mathcal{L}}{2} \right) \left( I + \Delta t \mathcal{L} \right) f^t + \left( I - \frac{\Delta t \mathcal{L}}{2} \right)^{-1} \Delta t \mathcal{N}(f^{t+\Delta t}, f^{t+\Delta t}) + \mathcal{B}(f^{t+\Delta t})$$

4. Simulation Results

4.1 Time evolution of fluctuations

At first, we investigate the case with the source term $S_1$. Figure 8 shows the temporal evolution of fluctuat-
ing internal energies of 10 modes, which have the largest amplitude at $t = 128$. The system starts from $T_1(r) = 0$, and gradually evolves by the drive of source term. In this simulation $T_{(3,0)}$ component dominates the system. In the early stage of evolution, the temperature gradient is small enough everywhere so that the system is stable. In this stage, fluctuations are suppressed to a negligible level. When the temperature gradient exceeds the threshold value of ITG instability, the energy of the most unstable mode of $(29,20)$ and $(30,20)$ start to grow and then energies of other modes such as $(1,0)$ and $(3,2)$ successively increase. These modes reach to saturated turbulent state by the nonlinear interaction in a short period. A very complicated behavior of modes is observed, at $t \approx 60$. With a short time delay, other modes such as $(16,11)$ and $(26,17)$ start to grow at $t \approx 75$.

Figure 9 shows the power spectrum of the internal energy in the linearly growing phase at $t = 55.5$. At $t = 55.5$, the amplitude of random fluctuations, which are shown as a mass of cross symbols at bottom right ($m > 50$) in Figure 9, is relatively comparable to that of the linearly growing modes, although peaks of unstable modes are seen at $m \sim 30$ and $m \sim 68$. These modes start to grow at $t = 49.5$.

Figure 10 shows $q$ profile and $\hat{k}$ profile at $t = 49.5$ with the location of the resonant surfaces of $(30,20)$, $(29,20)$ and $(68,50)$ modes. The modes such as $(30,20)$ and $(29,20)$ are destabilized at $r \approx 0.45$. Figure 11 shows the eigen mode of $(30,20)$ at $t = 51.0$. Both the modes of $(30,20)$ and $(29,20)$ have relatively large growth rate, because of large value of $\hat{k}$ in this region. They grow faster, as is seen in Figure 8. On the other hand, the mode such as $(68,50)$ is destabilized at $r \approx 0.6$. The resonant surface of $(68,50)$ locates near $r = 0.6$, where magnetic shear vanishes, so this mode is easily destabilized and also grows faster similar to $(30,20)$ mode. However, its energy saturation level is lower than the other modes (see Figure 20), so that it is not plotted in Figure 8.

4.2 ITB formation and relaxation phenomena

During the temperature evolution, the generation of internal transport barrier (ITB) is observed. Figure 12 shows temporal evolution of ion temperature profile, $T_i(r, \theta = 0, \phi = 0)$. After the ITG turbulence saturates at $t \sim 60.5$, thermal conductivity $\chi$ decreases and temperature gradient increases in $0.53 \leq r \leq 0.66$. The steep gradient of temperature around $r \approx 0.6$ exists until $t \sim 96$ and collapses. As the results, the temperature increases remarkably at $r \sim 0.6$, and steep gradient moves to around $r \sim 0.35$. These regions with steep gradient appear at the positions of rational surfaces shown in Figure 12.

In detail, this collapse phenomenon can be divided into two phases. The first phase is associated with the growing of $(4,3)$ mode, as is observed at $90 \leq t \leq 96$ in Figure 8. The second phase is associated with the damping of $(4,3)$ mode and with the growing of $(7,4)$ mode at $96 \leq t \leq 104$. In this study, the safety factor is fixed and the minimum $q$ value is $q_{min} = 1.35$. So the
Fig. 12 The growth and relaxation of radial profile of the ion temperature and the positions of resonant surfaces of the modes of (7,4), (11,7), (68,50) and (70,50).

Fig. 13 χ profile t = 85.0 ~ 96.0.

Fig. 14 k profile t = 85.0 ~ 96.0.

Fig. 15 χ profile t = 96.0 ~ 96.7.

Fig. 16 k profile t = 96.0 ~ 96.7.

Figure 13 and 14 show the temporal evolution of the radial profile of the thermal conductivity and temperature gradient k at 85 < t < 96, respectively. The regions painted by gray color in Figures 13-16 indicate regions where the transport barrier is produced. They are also indicated in Figure 12.

Until t ~ 91, thick strong transport barrier is observed in Figure 13. The large value of k is sustained around r ~ 0.6 during this phase. Then, the heat flux rapidly increases around r = 0.5 at t = 92 ~ 96. Correspondingly, the collapse of thick transport barrier occurs, and the value of k decreases in this region. However, in Figure 13, the thin transport barrier still remains at r = 0.7 at t = 96. On the other hand, the value of k increases at r ~ 0.35. During this phase, i.e., t = 92 ~ 96, (4,3) mode abruptly grows. The (4,3) mode is off-resonant.

Next, the thin weak transport barrier at r = 0.7 decays eventually followed by saturation of the (4,3) mode and growing of the (7,4) mode. During the decaying stage of the weak transport barrier, nonlinear relaxation phenomena, such as an avalanche, are observed in Figure 16. The position of the maximum value of k moves outward from t = 96.2 to t = 96.6, with reducing its absolute value, as shown by × symbols in Figure 16. The avalanche propagates ballistically and its velocity is estimated as \( v_{av} = 0.128 \) at r ~ 0.7. The inward avalanche also occurs at r ~ 0.35, and its velocity is \( v_{av} \approx -0.1 \).

During the avalanche, the energy stored inside of ITB is released. After the avalanche arrives at the edge region, the value of k starts to increase at the inside. We observe that the next peak is growing on the position of r ~ 0.65 at t = 96.7 indicated by circles in Figure 16. Namely, the weak transport barrier appears again. Similarly, next avalanche propagates inward and outward.
Figures 17-18 show the contour plots of generalized potential $F_i$ (including poloidal flow) and parallel ion velocity $V_{||}$ at the time before ITB collapses ($t=85.0$) and after that ($t=100$), respectively. At $t = 85$, strong flow is observed as a red belt in relief. This implies the strong shear flow is generated at just outside of $q_{min}$ location. The flow shear in $V_{||}$ and $F_i$ suppresses turbulence, and quenches anomalous transport. After the ITB collapses, such structures vanish, and both $V_{||}$ and $F_i$ profiles entirely change. Obviously the ITB formation is attributed by the suppression of turbulence by shear flow. Once, the zonal flow is damped by neoclassical viscosity, turbulence convection destroys the ITB.

After the collapse event, steep gradient still remains around $r \approx 0.35$. The thicker strong transport barrier like observed in the first phase never appears again. Around $r \approx 0.65$, gradient increases slightly with thinner transport barrier, and decays in a short time, which repeats as a limit cycle. In this way, the system evolves in a quasi-periodic manner associated with an intermittent energy transport. Finally, the system approaches to marginal state where the energy supplied from source and the intermittent heat transfer are balanced.

4.3 Source term effect

The temporal evolution of ion temperature with different source is investigated. Figure 19 shows the temporal evolution of fluctuating internal energy in the case with $S_2$. 10 Fourier modes are plotted, which have the largest amplitudes at $t = 140$. In spite of the same input energy per unit time, destabilized modes appear in different time. It is found that the system is stable until $t \sim 100$. At first, the modes of $(33,21), (28,19)$ and neighbor of these modes are excited at $r \sim 0.4$. These modes saturate with lower amplitudes than the modes of $(m,n) \sim (30,20)$ in the case with $S_1$. So that, ITG turbulence is relatively weaker. It is found that, no modes
are excited around $r \sim 0.6$. Namely, the value of $\hat{k}$ is not large enough to drive instability at $r \sim 0.6$.

In both cases, we commonly observe that the mode of $(1,0)$, i.e., Geodesic Acoustic Mode (GAM) grows associates with $(m,n) \sim (30,20)$ mode. Later, $(4,3)$ mode grows and saturates, which gives rise to the global profile changes. After $(11,7)$ mode reaches to almost same amplitude of $(7,4)$ mode, off-resonant $(4,3)$ mode starts to grow. Figures 20 and 21 show the extended view of Figure 8 and 19 in the saturation phase.

As the results, $(11,7)$ mode looses its fluctuation energy. This implies that $(4,3)$ mode is generated by $(7,4)$ and $(11,7)$ modes via three waves coupling. Collapse of ITB occurs associates with the saturation of $(4,3)$ mode ($t \sim 126$). The collapse mechanism is common, in spite of different source profile, although $(4,3)$, $(7,4)$ and $(11,7)$ modes in the case with $S_2$ saturate more rapidly compared with the case with $S_1$.

Figures 22-25 show the temporal evolutions of temperature profiles at $t = 48, 80, 100, 120, 140$. For comparison, the cases without profile relaxation are also shown in these figures. In this case, temperature profile evolves according to the diffusion equation without the additional anomalous transport. Until $t = 100$, profiles of two cases are quite different. However, at last, after the ITB collapse, the profiles of the two cases are gradually become similar.

Figures 26-29 show the temporal evolution of the radial profile of the heat conductivity $\chi$ and the temperature gradient $\Delta r$ during the collapse event in the case with $S_2$. In this case, collapse event occurs at $t \sim 126$. The thick transport barrier is observed at $r \sim 0.6$ in the Figure 26, similar to case of the $S_1$ until $t \sim 120$. 

![Fig. 22 t = 80 : ITB is generated in case $S_1$.](image)

![Fig. 23 t = 100 : ITB of case $S_1$ is decayed.](image)

![Fig. 24 t = 120 : ITB is generated in case $S_2$.](image)

![Fig. 25 t = 140 : ITB of case $S_2$ is decayed.](image)
The collapse process is also similar, however, the average value of the $\chi$ is smaller than the case of $S_1$. The value of the $\kappa$ in the ITB region is small. The avalanche with collapse occurs also smaller in scale, as shown in Figures 28. The value of $\kappa$ at $r \sim 0.3$ is not so much changed in this case.

After the decay of (4,3) mode, energy of all modes attain the quasi-steady state and limit cycle appears in the heat flux evolution. Figures 30 and 31 show the temporal evolution of the heat flux measured at $r = 0.5$ and $r = 0.7$ in the case with $S_1$ and $S_2$, respectively. In this phase, the intermittent energy transport is accompanied with the avalanche. Before the collapse of ITB, we observe the precursor, then after the collapse, amplitude and the frequency of intermittent burst become larger. The peak values of the burst observed at $r = 0.5$ and $r = 0.7$ locations in the case with $S_1$ are larger than those with $S_2$. So are the averaged value of heat flux.

The heat source affects on the energy transport in the marginal state. However, the radial temperature profile of each case in the marginal state becomes similar. At first, the heat source effect appears as the difference of temperature profile and affects the ITG turbulence via the $\kappa$. In the case with $S_1$, continuously excited ITG turbulence is stronger, because of the larger supply of the driving term $\kappa$ to the region $r \sim 0.45$ than the case with $S_2$. The larger heat flux at $r = 0.5$ in Figures 30 seems to indicate the continuous existence of the stronger turbulence. The differences of the energy saturation level of ITG modes between the two cases, shown in Figures 20 and 21 imply that the balance point between the suppression of the ITG turbulence by zonal flow and the generation of zonal flow by turbulence is different in two cases. The difference of the property of intermittent energy transport between the cases with $S_1$ and $S_2$ is seemed to be the results of the different quasi-steady state arranged by the turbulence and the zonal flow. To investigate the relation between the intermittent transport and zonal flow interaction quantitatively, it is necessary to evaluate the profile of $E \times B$ flow shearing rate. It is left for a future work. The plasma adjusts the energy transfer through the multi-scale interaction between the turbulence and the shear flow for the different heat source $S_1$ and $S_2$, and self-organizes toward to the similar temperature profile.

5. Summary and discussion

We have developed global ITG simulation code and investigated the multi-scale interaction between transport and turbulence in the reversed shear configuration. The numerical simulations have been carried out under the finite heat source, starting from the initial temperature profile which is stable against all ITG modes. The multi-scale interaction between the ITG turbulence of short time scale and the profile modification of long time scale during the plasma heating has been clearly shown. The different form of source term leads to the different temperature profile, however, similar modes are excited according to the free energy profile developments.

In the temporal evolution of fluctuation, the time scale in which linear modes are excited depends on the source profile. This can be naturally understood as the result of steepness of temperature gradient provided by source term $S_1$ in comparison with $S_2$. The $\kappa$ profile is also different from each other. In the case with $S_1$, (30,20) mode at $r \sim 0.4$ and (68,50) mode at $r \sim 0.6$ are destabilized at almost the same time $t \approx 50$. However in the case with $S_2$, driving term $\kappa$ has not reached to threshold value at $r \sim 0.6$ when the (m,n) $\sim$ (30,20) becomes unstable at $t \approx 100$.

In the case with $S_1$, relatively stronger ITB formation is observed at $r \leq 0.6$, than the case with $S_2$. The thick ITB is generated at $r \sim 0.6$ in both cases. The (7,4) and (11,7) mode contribute to generate (4,3) mode via beat-wave interaction in both cases. In both cases, the growth and decay of (4,3) mode, which is off-resonant mode, around $r \sim 0.6$ are accompanied with the ITB collapse. The weaker ITB is observed at $r \sim 0.7$ in both cases after the collapse. It grows slightly and decays in a short time repeatedly with avalanche. The quasi-periodic manner associates with intermittent energy transport.

The properties of the intermittent transport in two cases were different from each other. However, the radial...
profiles of two cases were going to become similar at the marginal state through the multi-scale interaction. In other words, the self-organization was observed, and the plasma selected spontaneously its global structure, regardless of the different form of the heat source.

To investigate the source term effect in detail, it is necessary to study the dynamics at the $q$-minimum region and interaction between the off resonant modes. In addition, the model should be extended at least two points of view: i) the ballooning mode stability should be treated accurately, and ii) consistent treatment of the $\rho$, with the heating simulation should be done. The second part is more challenging issue to understand the interaction between transport and turbulence. Those are left as a future work.

References