

UNIVERSITÉ DE PROVENCE



PHYSIQUE DES INTERACTIONS IONIQUES ET MOLECULAIRES

(UNITÉ MIXTE DE RECHERCHE DU C.N.R.S. 6633)

Centre de St. Jérôme - Case 321 - F 13397 Marseille Cedex 20



THÈSE POUR OBTENIR LE GRADE DE DOCTEUR

DISCIPLINE: RAYONNEMENT ET PLASMAS

ECOLE DOCTORALE: PHYSIQUE ET SCIENCES DE LA MATIÈRE

ETUDE PAR SIMULATION NUMÉRIQUE DU TRANSPORT RADIAL DANS LE PLASMA DE BORD D'UN TOKAMAK

Satoru SUGITA

Thèse présentée et soutenue publiquement le 11 Janvier 2011
devant le jury composé de Messieurs:

S. BENKADDA	Directeur de Recherche CNRS, Université de Provence, Marseille, FRANCE	(Directeur de thèse)
S.-I. ITOH	Professeur à Kyushu University, Kasuga, JAPAN	(Co-Directeur de thèse)
K. ITOH	Professeur à NIFS, Toki, JAPAN	(Rapporteur)
I. CALDAS	Professeur à Universidade de São Paulo, São Paulo, BRAZIL	(Rapporteur)
M. YAGI	Professeur à Kyushu University, Kasuga, JAPAN	(Examineur)

Simulation Study on Radial Transport in Tokamak Scrape-off Layer

Satoru SUGITA^{*1}

updated December 1, 2010

^{*1} Interdisciplinary Graduate School of Engineering Sciences, Kyushu University

RÉSUMÉ

L'un des récents développements des observations expérimentales de plasma de bord (SoL, pour "Scrape-Off Layer") dans les dispositifs à plasma à confinement magnétique, est que le profil radial de densité s'étend loin de la surface du plasma. Le gradient de densité à proximité de la surface magnétique la plus externe est nettement différent du gradient dans la région loin de la surface du plasma. Un profil de densité plat dans cette dernière région, la SoL externe, a été observé, et est souvent appelé "2ème SoL". Une telle densité plasma étendue dans la 2ème SoL induit une interaction plus importante entre le plasma et la paroi, ce qui augmente la quantité d'impuretés et le taux de recyclage sur la paroi. Il est donc important d'explorer le mécanisme de formation du profil de densité dans la 2nde SoL. Dans des observations plus détaillées, il existe de nombreuses indications expérimentales que les filaments de plasma alignés avec le champ magnétique, appelés "blobs", sont générés dans la SoL et causent un transport radial convectif rapide. Sur la base de ces études expérimentales, il est maintenant accepté que ces filaments jouent un rôle important dans le transport dans le plasma de bord. Les mécanismes qui les génèrent et leur dynamique sont donc des sujets importants.

Dans cette thèse, les phénomènes fondamentaux du transport dans le plasma de bord sont étudiés par des simulations numériques en mettant l'accent sur le phénomène de filaments plasma. Dans un premier temps, les mécanismes de propagation de blobs uniques et isolés sont envisagés. Puis la génération de

blobs par la turbulence de bord est étudiée, et le transport turbulent est discuté en tant que phénomène collectif. Des particularités du transport turbulent, incluant les blobs auto-organisés dans la SoL, sont reliées à un transport de type Bohm (c'est à dire des perturbations avec des corrélations radiales longues, et un coefficient de transport effectif qui suit la dépendance Bohm). De plus, en prolongement de ce travail, un effort initial vers une transposition du transport non-local au plasma de bord est décrite.

Cette thèse est organisée de la manière suivante. Dans le premier chapitre, la motivation et les objectifs de la thèse sont présentés. Dans le second chapitre, l'état des connaissances sur le transport de blobs plasma dans la SoL sont résumés. Dans le chapitre 3, les équations du modèles utilisées dans les chapitres suivants sont expliquées. Les perturbations de type blob ont en réalité une structure tridimensionnelle. Cependant, on peut en constituer des équations 2D. En partant d'équations de type fluide, les équations 2D Hasegawa-Wakatani sont introduites pour décrire la dynamique des blobs dans la SoL à l'aide des concentrations locales de densité et de vortacité. Dans le chapitre 4, sur la base du modèle théorique introduit dans le chapitre précédent, la propagation de blobs uniques est étudiée en détail. Les conditions de la simulation sont expliquées, et ensuite les modèles sont comparés. L'attention centrée sur le modèle, que l'on appelle "modèle de conservation", la vitesse du blob est étudiée. L'évolution de sa taille et la relation avec la vitesse est obtenue. Dans le chapitre 5, le transport induit par la turbulence d'échange ("interchange") dans la SoL est étudiée en effectuant une simulation non linéaire à flux fixe des l'équations du modèle 2D Hasegawa-Wakatani, qui est dérivé dans le chapitre 3. L'évolution autoconsistante de perturbations de type blob est confirmée. La vitesse du blob généré est évaluée par les fonctions d'autocorrelation du flux turbulent de particules dans l'état quasistationnaire. Elle confirme l'estimation théorique, qui est construite par (i) une règle de sélection de

la taille des blobs et (ii) la formule de la vitesse d'un blob en fonction de sa taille. La vitesse radiale de blobs générés spontanément suit la relation $V_x \sim 0.66 (L_{\parallel}/R)^{1/5} (\rho_s/R)^{2/5} c_s$, o $L_{\parallel}$ indique la longueur de connexion de la ligne de champ magnétique dans la SoL, R le grand rayon du tokamak, ρ_s le rayon de Larmor acoustique ionique, et c_s la vitesse acoustique. La loi d'échelle du coefficient de transport est étudiée en détail, et sa dépendance de type Bohm est discutée. Il est observé que le coefficient de transport D_{eff} (qui est introduit comme la moyenne sur la distance radiale de propagation du blob) suit la relation $D_{eff} \propto (L_{\parallel}/R)^{2/5} \rho_s c_s$. Dans le chapitre 6, l'étude du transport turbulent 3D dans des simulations à flux fixe du plasma de coeur (c'est à dire avec un profil évoluant de manière consistante avec un flux de source donné) révèle le comportement du gradient de pression comme une fonction du flux total d'énergie imposé. Le rôle de la rotation du plasma dans ce processus est étudié en examinant les différents scénarios par rapport à la rotation du plasma (turbulence induite par la rotation, rotation supprimée artificiellement ou imposée). Le résumé et une discussion sont présents dans le chapitre 7.

Abstract

One of the recent important developments in experimental observations of Scrape-off Layer (SoL) plasmas in magnetic confinement device is that the radial density profile in SoL extends far away from the plasma surface. The density gradient near the outermost magnetic surface is distinctively different from the gradient in the region far from the plasma surface. The flat density profile in the latter region, i.e., the outer SoL region has been observed, which is often referred to as the “2nd SoL”. Such an extended density of plasmas in the 2nd SoL leads to the stronger plasma-wall interaction, so as to increase impurities and recycling rate on the first wall. Therefore, the mechanism of density profile formation in the 2nd SoL should be explored. In detailed observations, there are abundant experimental indications that the magnetic-field-aligned plasma filaments, referred to as “blobs”, are generated in the SoL, and they cause the rapid radial convective transport. Based on these experimental studies, it is now experimentally accepted that plasma filaments or blobs play important roles in the transport of SoL plasmas. The mechanisms that generate blobs and the dynamics of blobs are important issues.

In this thesis, putting an emphasis on the plasma blob phenomenon, we study fundamental processes of SoL transport using numerical simulation. At first, we investigate the propagation mechanisms of single and isolated blobs. Next, we study the generation of blobs from edge turbulence, and discuss the SoL turbulent transport as a collective phenomenon. Features of turbulent trans-

port, which includes the self-organized blobs in SoL, are identified as Bohm-like transport (i.e., the perturbation with long radial correlations and the effective transport coefficient that follows the dependence of Bohm-like transport). Additionally, as an advancement of study, we describe an initial effort to extend the view of nonlocal transport to edge plasmas.

The thesis is organized as follows: In Chapter 1, the motivation and the goal of the thesis is discussed. In Chapter 2, the background and the status of present research of the plasma blob transport in the SoL are reviewed. In Chapter 3, we explain the model equations which are used in the following chapters of the thesis. The blob-like perturbation has, in reality, three-dimensional structures; However one can constitute the 2D model equations. Starting from fluid-like equations, 2D Hasegawa-Wakatani type model equations are introduced to describe the dynamics of blobs in SoL with local concentrations of density and vorticity. In Chapter 4, based on the theoretical model introduced in the previous chapter, the single blob propagation is investigated. We study in detail the evolution of a single blob. We first explain the simulation conditions, and then models are compared. Focused on one model, which we call ‘conserved model’, we study the velocity of the blob. Evolution of the blob size and relation with the velocity is obtained. In Chapter 5, the transport driven by interchange turbulence in SoL is investigated by performing a flux-driven nonlinear simulation of a two-dimensional Hasegawa-Wakatani-type model equation, which is derived in Chapter 3. The selfconsistent evolution of blob-like perturbation is confirmed. The velocity of the generated blob is evaluated from autocorrelation function of turbulent particle flux in the quasi-steady state. It confirms the theoretical estimate, which is constructed by (i) selection rule of blob size and (ii) the velocity formula of a blob as a function of the blob size. The radial velocity of spontaneously generated blobs follows the relation $V_x \sim 0.66 (L_{\parallel}/R)^{1/5} (\rho_s/R)^{2/5} c_s$, where $L_{\parallel}$ indicates the connection

length of the magnetic field line in the SoL, R the tokamak major radius, ρ_s the ion sound Larmor radius, and c_s the ion sound speed. The scaling of the transport coefficient is studied in detail, and its Bohm-like dependence is discussed. Effective transport coefficient D_{eff} (which is introduced as an average over the radial distance of blob propagation) is found to follow the relation $D_{\text{eff}} \propto (L_{\parallel}/R)^{2/5} \rho_s c_s$. In Chapter 6, we report the 3D turbulent transport in a flux-driven simulation for edge plasmas (i.e. with selfconsistent profile evolution for given source of flux), revealing the behavior of the pressure gradient as a function of the imposed total energy flux. Role of plasma rotation in this process is studied by examining the different scenarios with respect to plasma rotation (turbulence driven rotation, artificially suppressed or imposed rotation). The summary and the discussion are given in Chapter 7.

Contents

RÉSUMÉ	i
Abstract	v
1 Introduction	1
1.1 Nuclear Fusion	1
1.2 ITER	3
1.3 Tokamak Concept and Scrape-off Layer	4
1.3.1 Tokamak	4
1.3.2 Divertor configuration	5
1.4 Subject of This Thesis	6
2 Reviews of Physics in Scrape-off Layer	9
2.1 Studies of Scrape-off Layer in Tokamaks	9
2.1.1 Core and scrape-off layer	9
2.1.2 Influence of SoL plasmas on fusion research	10
2.1.3 Characteristics of particle flux and distribution in SoL plasma	11
2.1.4 Note on the problem	13
2.2 SoL Phenomena and Plasma Blob Concept (Experiment) . . .	14
2.2.1 Observation of the second SoL	14
2.2.2 Blobs as an origin of enhanced transport in SoL	16

2.2.3	Details of a blob motion	17
2.2.4	Short summary of experimental observations	20
2.3	Theoretical Model of Blobs	20
2.3.1	Model and propagation mechanisms of blobs	20
2.3.2	Directions in the theory and simulation research	25
2.4	Impact on Physics of Multiple-scale Turbulence	26
2.4.1	Multiple-scale physics	26
2.4.2	Study of nonlinear process in laboratory plasmas	27
2.5	Problem Definition for This Thesis	27
3	Blob Transport Model	29
4	Single Blob Dynamics	39
4.1	Simulation Conditions	40
4.2	Model Comparisons	41
4.3	Single Blob Evolution and Propagation Velocity Analysis	44
4.3.1	Deformation of blobs	44
4.3.2	Velocity of blobs	45
4.3.3	Evolution of blob size	48
4.4	Summary of Chapter 4	51
5	Selfconsistent Blobby Transport	55
5.1	Model	56
5.2	Linear Stability of Interchange Mode in SoL Plasma	58
5.3	Formation of Mesoscale Strucutre	59
5.4	Turbulence-driven Radial Flux	65
5.5	Statistical Analysis of Turbulent Transport	68
5.6	Scaling of Transport Velocity	72
5.7	Scaling Relation of Turbulence-driven Flux	76

5.8	Summary of Chapter 5	81
6	Extension to Edge Turbulence: on Stiffness of Transport	83
6.1	Introduction	83
6.2	Illustration of the Problem	84
6.3	Model	84
6.3.1	Model equation	84
6.3.2	Geometry and simulation conditions	86
6.4	Simulation Results	87
6.4.1	Evolutions of pressure gradient and flow (case of spontaneous rotation)	87
6.4.2	Fitting to gradient-flux relation	89
6.4.3	Role of rotation in gradient-flux relation	91
6.5	Summary of Chapter 6	95
7	Summary	97
A	Dependence on Model and Numerical Scheme	101
A.1	Partially Linearized Model	101
A.1.1	Comparison between spectral and pseudo-spectral codes	101
A.1.2	Evolution of small blobs	104
A.1.3	Evolution of large blobs	106
A.2	Full Nonlinear Model	107
A.2.1	Evolution of small blobs	108
A.2.2	Evolution of large blobs	110
A.3	Detailed Discussion on Effects of Density Nonlinearity	113
B	Spectral and Pseudo-Spectral Method	117
C	Linear Stability Analysis of Interchange Mode in SoL	121

Acknowledgements	127
Bibliography	129

Chapter 1

Introduction

1.1 Nuclear Fusion

It is now widely considered that the average temperature of Earth's near-surface atmosphere and oceans have been increasing since the middle of the 20th century. The increase of the temperature is conjectured to be caused by increasing concentrations of greenhouse gases, which result from human activity such as the burning of fossil fuel and deforestation. Moreover, fossil fuel is limited. The principle of supply and demand suggests that the price of fossil fuel will keep rising. Thus, demands for exploring new energy source have become more and more pressing. Many candidates for the new energy sources have been investigated such as photovoltaic power, biomass fuel, wind power, geothermal power, etc. Nevertheless, the further efforts to realize another way for new energy are necessary in order to provide mankind sufficient amount of energy for globally increasing consumption in the future. Nuclear Fusion is expected as an alternative energy source that can sustain the global economic growth.

Nuclear Fusion is the process of energy source for the Sun and stars. In the core of the Sun, at temperatures of 10-15 million Kelvin, Hydrogen is converted to Helium by fusion to provide enough energy to keep the Sun burning.

A world-wide research programme is underway, aiming at fusion energy to

produce electricity on the Earth. When it succeeds, this will offer an enough alternative energy supply, with significant advantages in the environmental, supply and safety aspects, over present energy sources. To realize fusion on the Earth, different and more efficient fusion reactions than that is working in the Sun are chosen - those between the two heavy isotopes of Hydrogen: Deuterium (D) and Tritium (T). All isotopes of Hydrogen contain one proton and one electron. Protium, the common form of Hydrogen has no neutron, Deuterium has one neutron, and Tritium has two. If forced together, the Deuterium and Tritium nuclei fuse and then break apart to form a helium nucleus (two protons and two neutrons) and an uncharged neutron. The excess energy from this fusion reaction (which is released because the products of the reaction are bound together in a more stable way than the reactants) is mostly contained in the free neutron.

Fusion occurs at a sufficient rate only at very high temperatures and pressures - on earth, temperatures greater than 100 million Kelvin are required. At these extreme temperatures, the Deuterium - Tritium (D-T) gas mixture becomes a plasma. In a plasma, the atoms are disintegrated - electrons are stripped from the ions. For the positively charged ions to fuse, their temperature (or energy) must be sufficiently high to overcome their natural charge repulsion.

In order to utilize fusion energy, scientists and engineers have been learning how to control very high temperature plasmas. The plasmas at much lower temperature are now widely used in industry, especially for nano-technologies. However, the control of high temperature fusion plasmas has several major scientific and engineering problems - how to heat a plasma to in excess of 100 million Kelvin, how to confine such a plasma, and how to sustain it so that the fusion reaction can become established.

Table. 1.1 Main parameters of ITER [1].

Total fusion power	500 MW
Additional heating power	50 MW
Q	≥ 10
Average 14 MeV neutron wall loading	$\geq 0.5 \text{ MW/m}^2$
Plasma inductive burn time	300-500 sec
Plasma major radius R	6.2 m
Plasma minor radius a	2.0 m
Plasma current	15 MA
Toroidal field at 6.2 m radius	5.3 T

1.2 ITER

ITER is an international scientific project that aims at demonstrating that it is possible to produce commercial energy from magnetic confinement nuclear fusion. The ITER device is currently under construction as the world's largest and most advanced experimental nuclear fusion reactor, at Cadarache, France. The members of projects are Japan, China, the European Union, India, Korea, Russia and the United States. The ratio of fusion power to input power is abbreviated as Q . " $Q \geq 10$ " represents the scientific goal of the ITER project: to deliver ten times the power that it consumes. From 50 MW of input power, the ITER machine is designed to produce 500 MW of fusion power. Construction of the facility began in 2008 and first plasma is expected in 2019 in the present project. ITER's objectives also include: To ignite a burning, self-sustaining plasma and to verify the tritium breeding concept. Table 1.1 shows the main parameters of ITER and fig. 1.1 shows the schematic view of ITER device [1].

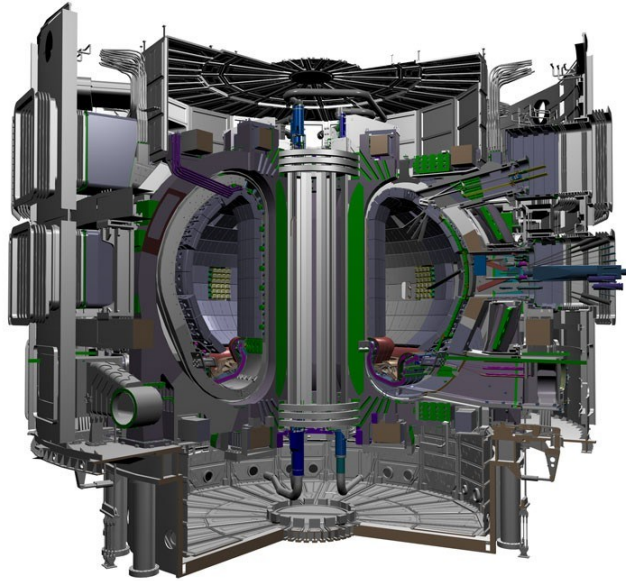


Fig. 1.1 Schematic view of ITER device [<http://www.iter.org>].

1.3 Tokamak Concept and Scrape-off Layer

1.3.1 Tokamak

Tokamak is one of the most-intensively-researched concepts for producing thermonuclear fusion power, and is a type of machine that uses a magnetic field to confine a plasma in the shape of a torus (fig. 1.2) [2]. In a tokamak the plasma is heated in a torus and kept away from the vessel walls by applied magnetic fields. The basic components of the tokamak's magnetic confinement system are (see figs. 1.2 and 1.3):

- The toroidal field - which produces a field around the torus. This is maintained by magnetic field coils surrounding the vacuum vessel. The toroidal field provides the primary mechanism of confinement of the plasma particles.

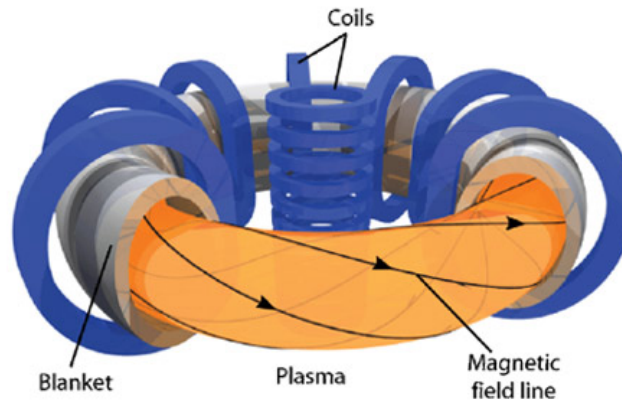


Fig. 1.2 In a tokamak the plasma is held in a torus-shaped vessel [<http://www.efda.org>].

- The poloidal field - which produces a field around the plasma cross section. It pinches the plasma away from the walls and maintains the plasma's shape and equilibrium. The poloidal field is induced both internally, by the current driven in the plasma (one of the plasma heating mechanisms), and externally, by coils that are positioned around the perimeter of the vessel.

These two fields consist the helical magnetic field lines.

1.3.2 Divertor configuration

Particles are confined within the volume composed of closed field lines. The border of the confined region is known as the Last Closed Flux Surface (LCFS) or separatrix, and the term Scrape-off Layer (SoL) indicates a region outside this border as is illustrated in fig. 1.3. In the SoL, the magnetic field lines are open (i.e. the magnetic field lines intersect with material plates), and the plasmas (which are exhausted from the core region) move along the field line and are pumped out at the divertor plate, with Helium ashes and impurities. There are two ways by which the last closed field line can be delimited. In

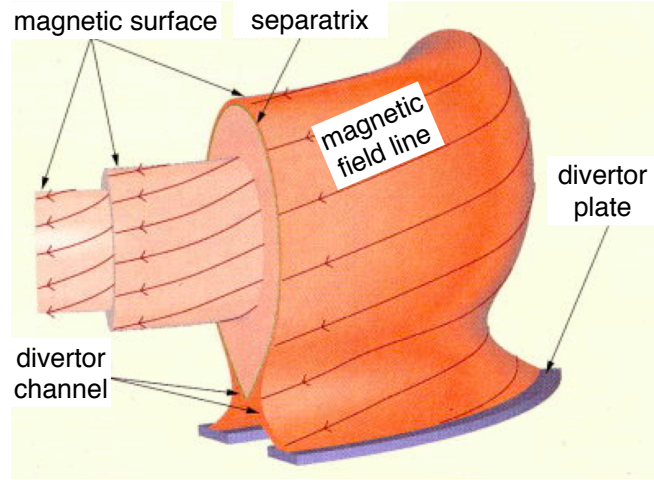


Fig. 1.3 Magnetic geometry of divertor configuration in tokamak device
[<http://www.lhd.nifs.ac.jp>].

the simplest option, the confined region is limited by a plate that is inserted to the plasma. This is called a limiter and essentially it was to protect the first wall from the hot core plasma. As a progress of the magnetic confinement techniques, a more sophisticated geometry was developed, using a separatrix of the magnetic field lines at the plasma edge, so that the field lines in the SoL are open and terminated at the plates at each end (see fig. 1.4). By guiding plasma flow along the magnetic field line to divertor, the motion of plasmas to the direction into the device wall is expected to be reduced. If the motion of plasmas to the device wall is *not* controlled, the strong plasma-wall interactions occur at the device wall. Then unfavorable effects (damage of the wall, impurity generation, etc.) are induced. The understanding of plasma dynamics in SoL is essential for magnetic confinement fusion research.

1.4 Subject of This Thesis

In this thesis, the physics of SoL plasma is studied, especially the radial turbulent transport is focused on. The thesis is organized as follows:

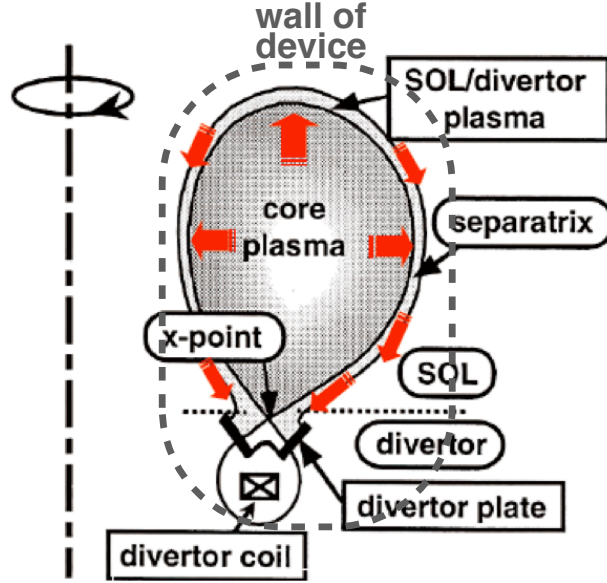


Fig. 1.4 Divertor Configuration in tokamak device.

In Chapter 2, the background and the present status of research of the plasma ‘blob’ [3] transport in the SoL are reviewed. In Chapter 3, we explain the model equations which are used in the following chapters of the thesis [4]. Starting from fluid-like equations, we introduce 2D model equations, which are modified from Hasegawa-Wakatani equations [5], in order to describe the dynamics of blobs (which is characterized by localized concentrations of density and vorticity). In Chapter 4, based on the theoretical model introduced in the previous chapter, the single blob propagation is investigated [6, 7]. We study in detail the evolution of a single blob. We first explain the simulation conditions, and then models (‘conserved’, ‘simplified’, ‘linearized’ models which are introduced in Chapter 3) are compared. Focused on one model, which we call ‘conserved model’, we study the velocity of the blob. Evolution of the blob size and relation with the velocity are obtained. In Chapter 5, the transport which is driven by interchange turbulence in SoL is investigated by performing a

flux-driven non-linear simulation of the two-dimensional Hasegawa-Wakatani-type model equation (derived in Chapter 3) [8,9]. The selfconsistent evolution of blob-like perturbation is simulated. The velocity of the generated blob is estimated from autocorrelation function of turbulent particle flux in the quasi-steady state. The scaling of the transport coefficient is studied in detail, and its Bohm-like dependence is discussed. In Chapter 6, we report the 3D turbulent transport in a flux-driven simulation (i.e. with selfconsistent profile evolution for given source of flux), revealing the behavior of the pressure gradient as a function of the imposed total energy flux. Role of plasma rotation in this process is studied by examining the different scenarios with respect to plasma rotation (turbulence driven rotation, artificially suppressed or imposed rotation). The summary and the discussion are given in Chapter 7.

Chapter 2

Reviews of Physics in Scrape-off Layer

2.1 Studies of Scrape-off Layer in Tokamaks

2.1.1 Core and scrape-off layer

The achievement of the plasma parameters in present-day tokamak plasmas, which has been the driving force to realize the ITER, was the outcome of research efforts that have focused at achieving high temperature and long confinement of core plasmas, as well as at controlling interactions of hot plasmas with wall of confining device. In the progress of the study of plasma turbulence, the turbulent convection due to fluctuating $\mathbf{E} \times \mathbf{B}$ velocity induces a major contribution to the observed anomalous radial transport. The importance of turbulent transport has also been recognized for the edge and SoL of all major fusion experiments, in limiter as well as in divertor configurations. The plasma parameters at the plasma edge and SoL have large impacts on the core confinement, so that the study of turbulent transport at the SoL gives enlightenments on the control of plasma-surface interaction as well as achieving the reactor-relevant plasmas in the core. One example is reproduced from ref. [10] in order to illustrate the impact of the edge region on the core plasma

(fig. 2.1).

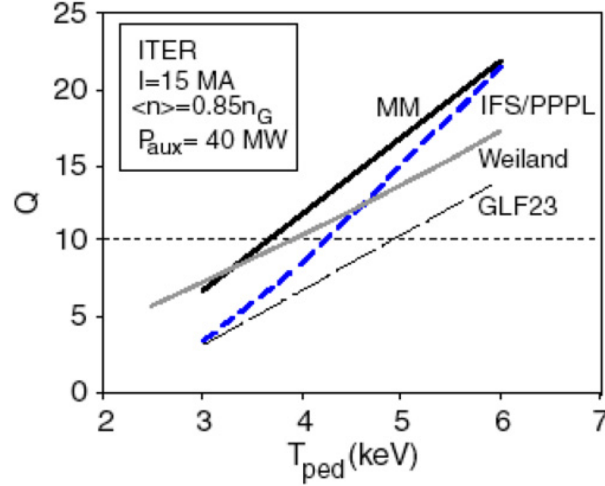


Fig. 2.1 Q versus T_{ped} (ion temperature at the pedestal top) predicted for ITER by the MM, IFS/PPPL, GLF23, and Weiland transport models. (Reproduced from [10]).

Here we first review the experimental observations of SoL transport, focusing on the picture of plasma blobs (Section 2.2). Then, theoretical efforts are illustrated (Section 2.3). Finally, the problems discussed in this thesis are described at the end of this chapter (Section 2.4). By investigating the dynamics of fluctuations in SoL plasmas, we hope to understand better the underlying plasma turbulence and to get insight into how to control the fluctuation-induced transport in SoL plasmas by experimental, theoretical, and numerical approaches.

2.1.2 Influence of SoL plasmas on fusion research

The physics of SoL plasmas is important for the development of fusion research. This is because the heat load on and damage of the first wall as well as divertor plates are one of the most crucial issues toward ITER. Moreover, the recycling of the impurity ions into the core plasma, across the SoL plasma region, has negative effect on the performance of the core plasmas. These key factors (e.g.,

heat load, generation and diffusion of impurities, etc.) are influenced primarily by the spatial distribution of SoL plasmas. Thus, the SoL physics and transport have attracted attentions in the history of fusion research. The recent trends in the research are reviewed in literatures, e.g., [11–21].

2.1.3 Characteristics of particle flux and distribution in SoL plasma

At first, the characteristic processes that control the radial density distribution in SoL are explained by considering the particle balance. Figure 2.2 shows the schematic diagram of the plasma flow from the confined plasma into the SoL. In

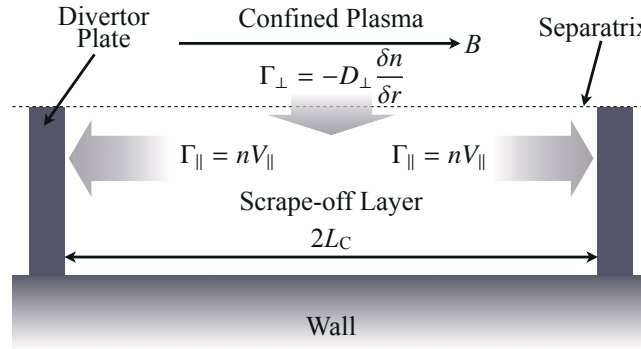


Fig. 2.2 Schematic diagram of the plasma flow from the confined plasma into the SoL by cross-field diffusion and the one along the field line in the SoL to the divertor target.

the classical picture of the SoL transport, the particle flux in the perpendicular direction (to the magnetic field line) is given by local transport process and is often assumed to be $D_{\perp} (dn/dr)$, and the rate of loss along the field line is evaluated by c_s/L_c , where $D_{\perp}$ is the cross-field diffusion coefficient, $c_s = \sqrt{T_e/m_i}$ is the ion acoustic speed and L_c is the connection length along the flux tube to the divertor plate. In the steady state, if there are no source or sink (such as ionization or recombination) in the scrape-off layer, and *if* the radial transport is expressed by a diffusion model, the loss along the field line

from any given flux tube is balanced by the net flow across the field line into the flux tube:

$$-\frac{d}{dr} \left[D_{\perp} \frac{dn}{dr} \right] = \frac{nc_s}{L_c}. \quad (2.1)$$

For the simplicity of argument, we consider the case that $D_{\perp}$ and c_s are independent of radius. In this case, eq. (2.1) is easily integrated to obtain

$$n(r) = n(a) \exp \left[-\frac{r-a}{\lambda_n} \right] \quad (2.2)$$

with

$$\lambda_n = \left(\frac{D_{\perp} L_c}{c_s} \right)^{1/2}, \quad (2.3)$$

where a is the plasma minor radius at the separatrix, r is minor radius, and λ_n is the electron density e -folding length. Equation (2.3) shows that the radial density profile in the SoL plasma decays exponentially.

Similarly, a temperature profile is evaluated using the electron heat balance as [22]:

$$T_e(r) = T_e(a) \exp \left[-\frac{r-a}{\lambda_T} \right] \quad (2.4)$$

with

$$1 + \frac{\lambda_n}{\lambda_T} = \frac{\delta}{5/2 + \chi_{\perp} \lambda_n / D_{\perp} \lambda_{T_e}}, \quad (2.5)$$

where λ_T is the electron temperature e -folding length, and $\chi_{\perp}$ is the cross-field thermal diffusivity. Equations (2.2) and (2.4) give a radial distribution of n_e and T_e .

Equations (2.3) and (2.5) give the e -folding lengths in the scrape-off layer in terms of the cross-field transport coefficients. Experimentally, the measured values of λ_n and λ_T are used to deduce $D_{\perp}$ and $\chi_{\perp}$. Typical values for λ_n and λ_T are ~ 10 mm. If there are no sources in the scrape-off layer, this measurement, together with the assumption that $T_e = T_i$, gives the estimation

of c_s and $D_\perp$. This estimation gives, typically, $D_\perp \sim 1\text{m}^2\text{sec}^{-1}$. Experimental results are often interpreted as that $D_\perp$ is inversely proportional to density. However, because the changes of density and temperature in the boundary region are correlated, this behaviour can also be interpreted as a temperature dependence of $D_\perp$. The decisive conclusion of the parameter dependence has not yet been obtained. If either particle or energy sources are present in the SoL, then the calculation of transport coefficients is much more difficult.

The parameters in the SoL plasma is connected to those in core plasmas. A particle confinement time of the core plasma, τ_p , can be defined as

$$\Gamma_n = \frac{\bar{n}V}{\tau_p}, \quad (2.6)$$

where Γ_n is the total particle outflux across the plasma surface, $\bar{n}$ is the average plasma density and V is the plasma volume. Using eq. (2.2) the total flux flowing into the SoL region is expressed as

$$\Gamma_L = 4\pi a \int_a^\infty n(a) \exp\left[-\frac{r-a}{\lambda_n}\right] c_s dr. \quad (2.7)$$

By carrying out the integration in eq. (2.7) and equating Γ_n to Γ_L , the edge density is given by

$$n(a) = \frac{\bar{n}V}{\tau_p} \frac{1}{4\pi a \lambda_n c_s}. \quad (2.8)$$

Equation (2.8) shows the relation between $n(a)$, $\bar{n}$, τ_p , and λ_n .

2.1.4 Note on the problem

The consideration of the particle transport in Subsection 2.1.3 illustrates the important role of radial transport of particle in the SoL plasma. It controls the density-fall-off length, which has strong effects on the plasma-wall interaction at the device wall. The interaction at the wall dictates the generation of impurities, so that the quantity of impurities is also affected. The relation eq. (2.8)

implies the connection with the core confinement. Thus, the clarification of the cross-field transport has primary importance for fusion research. In the Subsection 2.1.3, the conventional model based on diffusive transport process is explained. In the next sections, we review the experimental observations and then theoretical studies.

2.2 SoL Phenomena and Plasma Blob Concept (Experiment)

2.2.1 Observation of the second SoL

Motivated by the recognition of important role of SoL plasmas for fusion research, detailed experimental observations of SoL plasmas have been accumulated. One of the recent important developments is that the measurement of radial density profile in SoL has reported “the second SoL”. Experiments have shown that the density gradient near the outermost magnetic surface is distinctively different from the gradient in the region far from the plasma surface. The flat density profile in the outer SoL region has been observed (see fig. 2.3), which is often referred to the 2nd SoL. The typical example for the 2nd SoL is quoted from JT-60 experiments [23] in fig. 2.3.

The 2nd SoL is observed not only in the Large tokamak (JT60 in fig. 2.3), but also in middle size tokamaks (Alcator C-mod [24] in fig. 2.4). Similar observations have been reported in many tokamaks.

This observation of 2nd SoL means that the density profile in SoL cannot be described by the simple exponential distribution and that the transport process is *not* governed by a simple diffusive process with one diffusion coefficient. Such an extended density of plasmas in the 2nd SoL (figs. 2.3 and 2.4) leads to the increase of impurities and recycling rate on the first wall. Therefore, the mechanism of density profile formation in the 2nd SoL should be explored.

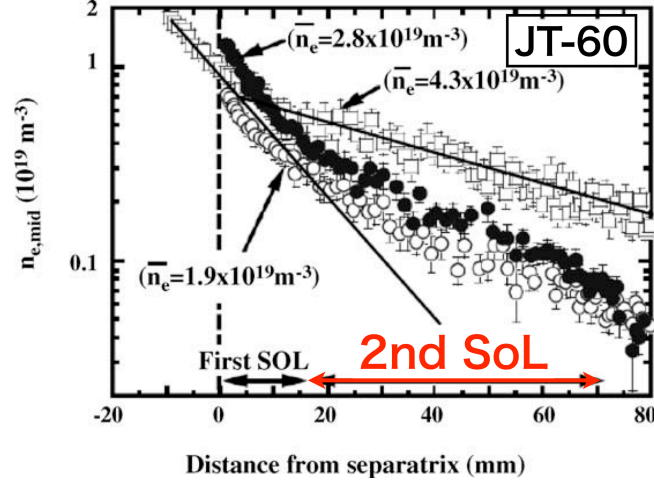


Fig. 2.3 Electron density as a function of distance from separatrix for two average bulk plasma densities. Vertical dashed line shows location of separatrix. Change of the decay length of density is clearly seen and two SoLs (first and second) are seen [23].

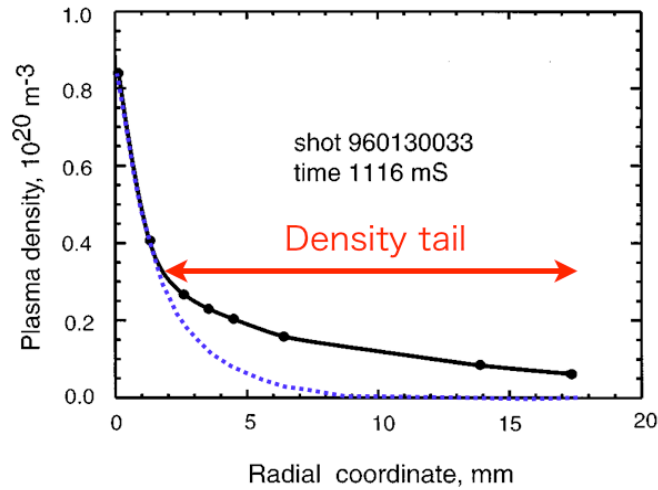


Fig. 2.4 Plasma density profile in SoL of Alcator C-mod tokamak [24]. Black points indicate experimental observations and solid line indicates fitting of them. Dotted line is the exponentially decaying curve, which implies collisional diffusive process.

2.2.2 Blobs as an origin of enhanced transport in SoL

In studying the origin of the 2nd SoL that has a long fall-off length, detailed observations have been made experimentally. There are abundant experimental indications that (1) the magnetic-field-aligned plasma filaments, referred to as “blobs”, are generated in the SoL, and that (2) they cause the rapid radial convective transport. Figure 2.5 is a characteristic example of filament structures observed on MAST spherical tokamak [25]. The evidence of plasma blob

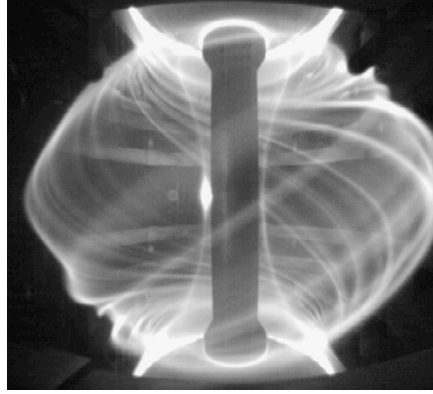


Fig. 2.5 Filament structure in MAST tokamak by high-speed camera imaging [25]

transport in the SoL is given in pictures taken by high-speed camera imaging (fig. 2.5). In recent years, the high-speed cameras have become common and image quality is improved significantly. Currently, a rate of more than 10^5 frames per second is possible. Figure 2.6 shows an image taken at the spherical tokamak NSTX with a high-speed camera [26]. The images were produced by filtered D_α light with the use of a local gas puff in the camera field of view to increase the local neutral density and thus the brightness (and contrast) of the fluctuation images (Gas Puff Imaging Method). We can see the filament structure near the edge.

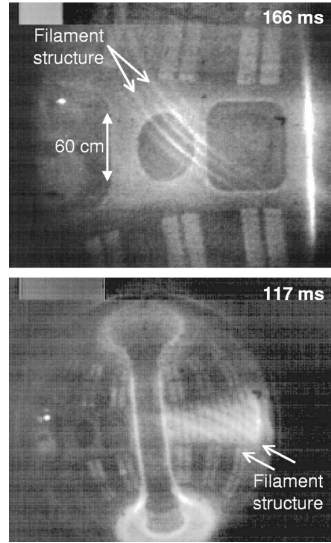


Fig. 2.6 Filament structure of plasma blob observed by gas puff imaging on NSTX [26].

2.2.3 Details of a blob motion

Figures 2.5 and 2.6 show the coexistence of many filaments. Picking up one filament, a single, magnetic-field-aligned-stripe structure is illustrated in fig. 2.7. The dynamics of an individual filament is now studied in detail.



Fig. 2.7 The predicted structure of an ELM in the MAST tokamak plasma geometry, based on the nonlinear ballooning mode theory [25]

The two-dimensional measurements of filaments are studied in many toka-

maks. Rapid propagation of plasma density clumps, which is associated with individual filament, are observed (figs. 2.8 and 2.9). Figure 2.8 shows consec-

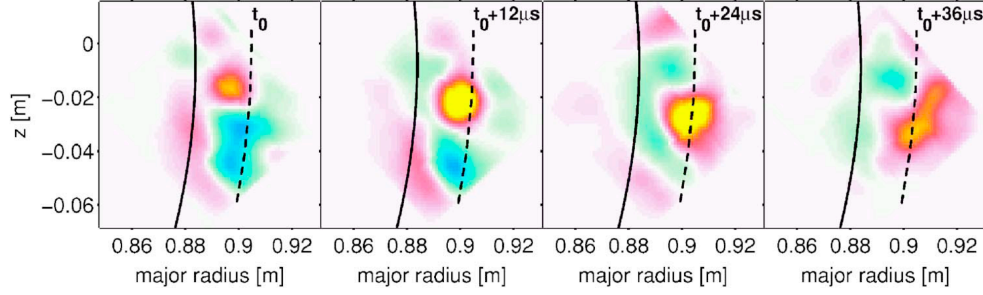


Fig. 2.8 The evolution of a long-living blob in Alcator C-mod [27].

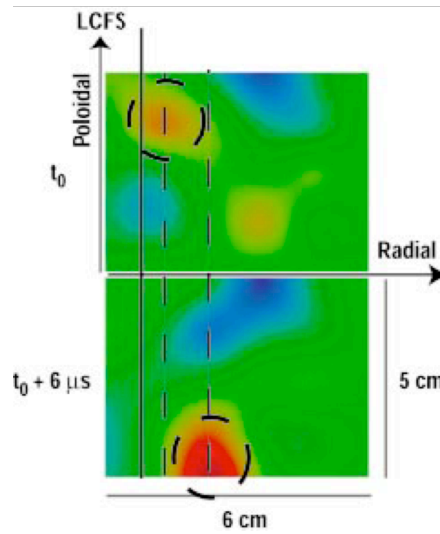


Fig. 2.9 Two frames from BES (Beam Emission Spectroscopy showing 2D density plots. Red indicates high density and blue low density [28].)

utive frames from an ultrafast camera of the D_α emission intensity fluctuation in a poloidal cross section in Alcator C-mod tokamak [27]. Red and yellow colors correspond to positive fluctuation amplitudes (i.e. local higher density), green and blue to negative. The separatrix is indicated by the black solid line. The picture shows the rapid convective propagation of the plasma blob in the SoL. The observed pair of the clump and hole of density moves radially

about one cm in 10 microseconds. That is, the radial velocity is of the order of $10^3[m/sec]$. This velocity is much faster than that associated with simple diffusion process. Figure 2.9 shows the experimental results in DIII-D, visualized by the BES (Beam Emission Spectroscopy) diagnostics [28]. The radial velocity of this density clump has similar value in comparison with the case in fig. 2.8. The density burst in plasma, marked with a dotted circle and shown in both frames, undergoes both poloidal and radial motions, maintaining the clump structure.

Filament structures (plasma blobs) are observed as an isolated perturbation of macroscopic density. They are formed at the plasma surface as is illustrated in fig. 2.10. (Figure 2.10 is observed in QUEST [29].) The motion of a solitary

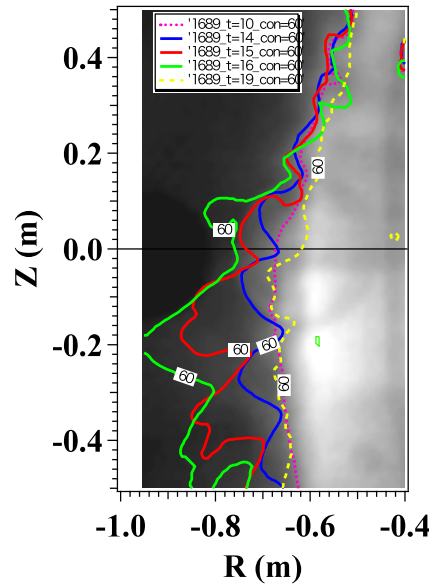


Fig. 2.10 Blob generation in slab geometry in the QUEST tokamak [29]. Contour on the poloidal cross section is shown for various time slices. Region of $R > -0.6$ indicates the high-density plasma.

structure is observed as a pulsive variation in a time series of density at one location. The sawtooth-like intermittent evolution of pulses was reported from the one-point measurement of SoL density on JT-60U (fig. 2.11) [30]. Each pulse

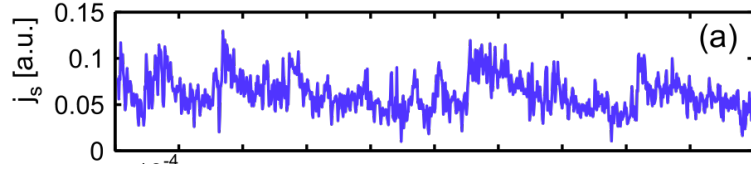


Fig. 2.11 Time evolution of ion saturation current near the low-field side SoL of JT-60U [30].

is considered to correspond to the passage of a blob (filamentary structure). The asymmetry in the fast rise and slow decay of a pulse indicates the asymmetry in spatial structure of a blob.

2.2.4 Short summary of experimental observations

As mentioned in the above discussions, plasma blobs with filamentary shape are observed in a large number of tokamak devices [23,24,26–28,30–40]. Recent detailed experimental studies in the laboratory plasmas and small devices have approached to clarify the plasma blob characteristics [25, 41–56]. Based on these experimental studies, it is now widely established that plasma filaments or blobs play important roles in the transport of SoL plasmas. The mechanisms that generate blobs and the dynamics of blobs are important issues.

2.3 Theoretical Model of Blobs

2.3.1 Model and propagation mechanisms of blobs

As a basis of the analysis in this thesis, we here make a brief overview of the theoretical model of blobs. The theoretical model for the plasma blob propagation was recently proposed by Krasheninnikov [3]. The basic mechanism behind the rapid convection appears to be the following: In general, any macroscopic clump of particles in a toroidal plasma would tend to be polar-

ized due to species-dependent processes such as the ∇B and curvature drifts (fig. 2.12). In the core plasma with its closed flux surfaces, this charge separation is short-circuited by free-streaming electrons. In the SoL, however, a finite potential can be supported on the open fields due to sheath formation and its resistivity at the end points (fig. 2.13). A poloidal electric field is thus formed, coupled with the toroidal magnetic field, and then leads to a rapid $\mathbf{E} \times \mathbf{B}$ drift of the blob. Figures 2.12 to 2.14 show conceptual diagrams of blob propagation mechanisms.

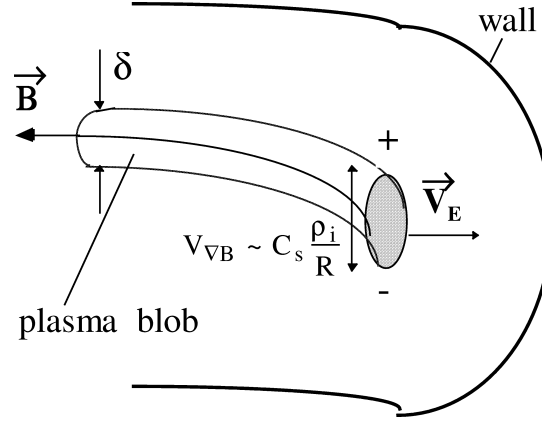


Fig. 2.12 The grad- B plasma polarization and associated $\mathbf{E} \times \mathbf{B}$ drift result in outward motion of plasma blob in tokamak far scrape-off layer [3].

We next introduce a heuristic argument of the blob motion. For simplicity, we assume that the SoL ion and electron temperature, T , is constant, and ignore the flow and the resistivity along the magnetic field lines. The electric current is described in a quasi-neutrality condition by the conservation of charge,

$$\nabla \cdot \mathbf{j}_{\perp} + \nabla \cdot \mathbf{j}_{\parallel} = 0, \quad (2.9)$$

where $\mathbf{j}_{\perp} = (\mathbf{B} \times \nabla p) / B^2$ and $\mathbf{j}_{\parallel}$ are the current densities in the directions which are perpendicular and parallel to the magnetic field line, respectively, $p = nT$ is the plasma pressure and n is the plasma density. In eq. (2.9), $\nabla \cdot \mathbf{j}_{\perp}$

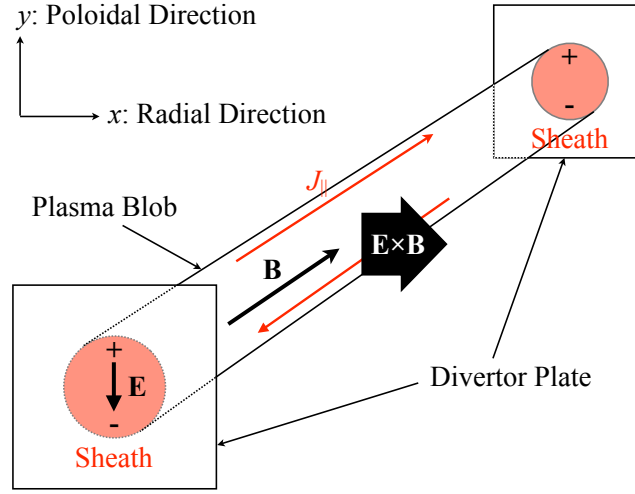


Fig. 2.13 Theoretical model and geometry of blob.

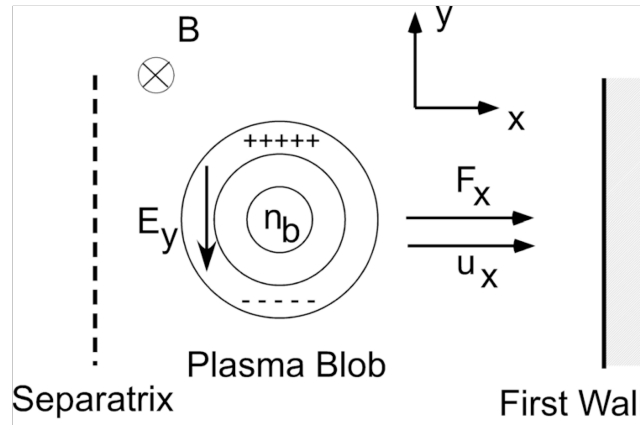


Fig. 2.14 Plasma blob 2D propagation model [12].

can be calculated as follows:

$$\begin{aligned}
 \nabla \cdot \mathbf{j}_{\perp} &= \nabla \cdot \left(\frac{\mathbf{B} \times \nabla p}{B^2} \right) \\
 &= \nabla \left(\frac{1}{B^2} \right) \cdot (\mathbf{B} \times \nabla p) + \left(\frac{1}{B^2} \right) \mathbf{B} \cdot (\nabla \times \nabla p) \\
 &\quad - \left(\frac{1}{B^2} \right) \nabla p \cdot (\nabla \times \mathbf{B}),
 \end{aligned} \tag{2.10}$$

where the second term becomes zero. We assume $\nabla p \cdot (\nabla \times \mathbf{B}) \approx 0$ (i.e. the

current is nearly perpendicular to the pressure gradient), then the third term can be neglected. Hence, we have

$$\nabla \cdot \mathbf{j}_\perp = \nabla \cdot \left(\frac{\mathbf{B} \times \nabla p}{B^2} \right) \cong -\frac{T}{B^2} \nabla \ln B \cdot (\mathbf{B} \times \nabla n). \quad (2.11)$$

Integrating eq. (2.9) along the magnetic field line gives

$$\int dl \frac{T}{B^2} \nabla \ln B \cdot (\mathbf{B} \times \nabla n) = \int dl \nabla \cdot \mathbf{j}_\parallel = 2j_\parallel|_{\text{target}}. \quad (2.12)$$

Here, $j_\parallel|_{\text{target}}$ is the current density at the end of the blob, which is connected to the sheath. The coefficient of 2 means the blob is connected to the sheath at each end and we assume that the current density is equal at each end.

We can estimate $j_\parallel|_{\text{target}}$ by introducing the sheath resistivity R_{sheath} , which is approximated by

$$R_{\text{sheath}} = \frac{V_f}{I_{\text{sat}}} = \frac{(1/2) \ln [4\pi m_e/m_i] T}{en_T c_s} \simeq \frac{T}{en_T c_s}, \quad (2.13)$$

where V_f is the floating potential, I_{sat} is the ion saturation current density in the sheath. n_T is the plasma density close to the wall, $c_s = \sqrt{T/m_i}$ is the ion acoustic speed. Thus, $j_\parallel|_{\text{target}}$ accompanied with the electrostatic potential at the blob end, ϕ , which is caused by the ∇B drift, is estimated from eq. (2.13). Replacing V_f and I_{sat} with ϕ and $j_\parallel|_{\text{target}}$, respectively in eq. (2.13), one has

$$j_\parallel|_{\text{target}} \approx en_T c_s \left(\frac{e\phi}{T} \right). \quad (2.14)$$

Eliminating $j_\parallel|_{\text{target}}$ from eqs. (2.12) and (2.14), we get

$$\frac{e\phi}{T} = \frac{\rho_i}{2n_T B} \int dl \nabla \ln B \cdot (\mathbf{B} \times \nabla n), \quad (2.15)$$

where $\rho_i = m_i v_\perp / (eB)$ is the ion sound Larmor radius and $v_\perp = \sqrt{T/m_i}$ is the ion thermal velocity. The magnetic curvature is evaluated by $\nabla \ln B = \hat{\mathbf{x}}/R$, where R is the major radius, x and y are the coordinates along radial

and poloidal directions, respectively, and $\hat{\mathbf{x}}$ is the unit vector in the radial direction (fig. 2.14). For the transparency of arguments we assume that the blob density is uniform along the magnetic field lines, and that the current density perpendicular to the field lines (diamagnetic current) is uniform. Considering a plasma blob of density n_b with parallel length L_b , we obtain the evaluation of the R.H.S. of eq. (2.15), which gives

$$\frac{e\phi}{T} = \frac{L_b \rho_i}{2R n_T} \frac{\partial n_b}{\partial y}. \quad (2.16)$$

Approximating $\partial n_b / \partial y \approx n_b / \delta$, $E_y \approx \phi / \delta$, where δ is the poloidal scale length of the blob, we have

$$E_y \simeq c_s B \left(\frac{\rho_i}{\delta} \right)^2 \frac{L_b}{R} \frac{n_b}{n_T}. \quad (2.17)$$

Thus we get a rough estimation of the blob radial propagation velocity $V_x = E_y / B$ as

$$V_x = c_s \left(\frac{\rho_i}{\delta} \right)^2 \frac{L_b}{R} \frac{n_b}{n_T}. \quad (2.18)$$

This equation indicates that a smaller blob has larger convective velocity. The stability of blobs should be taken into account in blob transport, since instability can break blobs into pieces so that blob has a larger velocity than the original one. As a result, the first-wall might suffer from more damages by larger blobs.

Next, we discuss the life time of a blob. The life time gives an estimate of the radial distance, Δ_x , that the blob can travel during its life time, $\tau_b \approx L_b / c_s$, before terminating due to parallel plasma transport. Taking an estimate, $L_b \sim qR$, where q is the safety factor, from eq. (2.18) we find

$$\Delta_x \approx V_x \tau_b \approx \left(\frac{\rho_i}{\delta} \right)^2 \frac{L_b^2}{R} \frac{n_b}{n_T}. \quad (2.19)$$

Description of isolated blobs makes sense only when $\Delta_x > \delta$, which gives

$$\delta < \delta_{\max} = R \left\{ \left(\frac{L_b \rho_i}{R^2} \right)^2 \frac{n_b}{n_T} \right\}^{1/3}. \quad (2.20)$$

Using number of plasma particles in a blob $N_b \approx n_b L_b \delta^2$, another expression is obtained as

$$N_b \approx n_b L_b \delta^2 < n_b R^3 \left\{ q^7 \left(\frac{\rho_i}{R} \right)^4 \right\}^{1/3} \left(\frac{n_b}{n_T} \right)^{1/3}. \quad (2.21)$$

For a large tokamak device, typical parameters are given by $q \sim 3$, $R \sim 1.5\text{m}$, $n_b \sim 10^{19}\text{m}^{-3}$, $c_s \sim 3 \times 10^4\text{m/sec}$, $\rho_i/R \sim 2 \times 10^{-4}$, and $n_b/n_T \sim 10$. For the blob of the size $\delta \sim 0.01\text{m}$ from eqs. (2.18)-(2.21), we find $V_x \sim 10^3\text{m/sec}$, $\Delta_x \sim 0.15\text{m}$, $\delta_{\max} \sim 0.03\text{m}$, and $N_b \lesssim 10^{16}$ from eqs. (2.18)-(2.21). These values are in the range of experimental reports.

We here note that the picture of ‘blob transport’ is complementary to the diffusive transport which is explained in the Subsection 2.1.3. The blob transport process is that, if the density perturbation of an isolated flux tube becomes steep enough, (say the radius of a filament δ of the order $(\rho_i/R)^{2/3}$ smaller than R as is in the illustrative example eq. (2.20)) the filament moves at large velocity V_x to a distance Δ_x , before the loss along the field line eliminates the high density in the blob. If the initial density perturbation is broad, the solitary feature of blob is lost rapidly, and diffusive transport takes place. Thus we would expect that blob-like process is more dominant in the case where large amplitude and localized perturbation is generated in SoL.

2.3.2 Directions in the theory and simulation research

As is illustrated in the Subsection 2.3.1, the model of ‘blob’ is relevant in understanding the dynamics of filamentary structures and associated transport phenomena. The key in the model is that, for the parameters of present day

tokamak experiments, one filament keeps an identity as a filament during the life time (Δ_x/V_x). Thus, in the study of a filamentary blob, a two-dimensional model (where inhomogeneity along the field line and the loss at the end are neglected) can be employed for the analytic transparency of arguments, as long as the motion within the radial distance of Δ_x is analyzed.

There are two lines of researches that investigate the blob dynamics and transport. One thread of ideas is, based on this concept of blob, the dynamics of a single blob in the SoL (focusing at the radial propagation and at deformation of blob shape) is investigated by many researchers. The outcome of researches, which have used analytical approach and numerical simulations, are reported in literature [3, 6, 7, 57–79].

The other thread of studies is to identify the generation mechanism of blobs and to analyze the selfconsistent evolution of blobs. The primary modes, which induce blobs in SoL, are considered such as curvature driven interchange or sheath interchange instabilities, the Kelvin-Helmholtz instability, rotational instability and so on. Several types of geometries (2D slab, or 3D torus) and model equations (Hasegawa-Wakatani-type resistive interchange, 3-field reduced MHD, or drift-Alfven ITG, etc.) are used to investigate the SoL transport, which is dominated by blob dynamics [8, 80–107]. Although instability mechanism for primary modes is different, the blob transport was found universal: The blob is the secondary mode and the convective transport is driven by the charge polarization.

2.4 Impact on Physics of Multiple-scale Turbulence

2.4.1 Multiple-scale physics

The study of blob transport in SoL plasmas has strong mutual interaction with the progress of modern plasma physics. One of the most remarkable progress

in the area of physics of magnetized plasmas is that the physics picture for the multiple scale turbulence has been established [108]. That is, turbulence and turbulent transport are governed not only by the local microscopic fluctuations, but also by mesoscale and macro scale perturbations. The new views of turbulence have been developed [108] theoretically [109,110] and experimentally [111,112]. Theories have shown that the fluctuations with long radial correlations can induce impacts on transient transport [108]. In the sense of the physics approach, the study of blobs belongs to the family of the multi-scale dynamics. The progress of the understanding of blobs will, no doubt, enrich the modern physics of plasmas.

2.4.2 Study of nonlinear process in laboratory plasmas

Detailed measurements of laboratory plasmas have given understanding of fundamental nonlinear process and dynamics (in which the perturbation scale length can be comparable to the gradient scale length) [113–116]. In these experiments, the correlation length, which is close to the scale length of gradients, is observed. In addition, the nonlinear interactions between fluctuation modes are directly measured. The elementary process that may induce the blob structure from fluctuations is being studied in the laboratory plasmas.

2.5 Problem Definition for This Thesis

The above discussion clarifies the following: major advances in experimental measurements and theoretical modelling have revealed that, in SoL plasmas, the transport across the magnetic field lines can be largely dominated by the radial motion of filamentary structures, which are localized in the plane perpendicular to the magnetic field and are elongated along the field lines. These filaments contain excess particles and heat, compared to background plasmas, and are thus referred to as plasma blobs. The rapid radial motion of blob-

like structures in the SoL causes a large cross-field transport of particles and heat, and influences strongly the properties of the SoL and operation of divertor. In particular, the blob-like transport leads to a broadening of the radial plasma profile (the second SoL) and thereby erosion of the main chamber walls. The transport mediated by such blob structures is likely related to a number of outstanding problems for the steady operation of a fusion experiment, including particle recycling, material migration, plasma rotation and the density disruption limit.

Thus, the dynamics of blobs and the associated radial transport are the subject of this thesis. Particularly, the blob velocity and the evolution of size during fast radial propagation are important issues. Characteristics of the transport are also crucial. For instance, transport velocity, transport coefficient, radial profile associated with the transport are the key elements. These issues are addressed in Chapters 3, 4, and 5. In addition, this thesis describes an initial effort to extend the view of nonlocal transport to edge plasmas in Chapter 6. By this effort, unification for the multiple-scale picture of turbulence will be advanced [108,117].

Chapter 3

Blob Transport Model

In this chapter we explain the derivation of the model equations which are used in the following chapters of the thesis. The essence of the process considered was explained in Chapter 2, figs. 2.12 to 2.14. The blob-like perturbation has, in reality, three-dimensional structures; However, one can constitute 2D model equations. As is explained in Subsection 2.3.1, the loss along the field line may be neglected during the period that the blob moves the distance Δ_x (eq. (2.19)). Thus simple 2D model has relevance to real 3D system in order to understand the physics of blob within the period, during which a blob moves a distance Δ_x . By starting from fluid-like equations, 2D Hasegawa-Wakatani type model equations are introduced to describe the dynamics of blobs in SoL with local concentrations of density and vorticity. Derivation of the model equations was developed by D'Ippolito, *et al.* [58], and we follow the procedure to derive the model equations. Then three different models are explained after.

The two-fluid model consists of the vorticity equation and the continuity equation:

$$\frac{d}{dt} \nabla_{\perp}^2 \Phi = \frac{B^2}{m_i n} \nabla_{\parallel} J_{\parallel} + \frac{2B}{m_i n} (\mathbf{b} \times \boldsymbol{\kappa}) \cdot \nabla p \quad (3.1)$$

$$\frac{\partial n}{\partial t} + \nabla \cdot (n \mathbf{v}_e) = 0. \quad (3.2)$$

where Φ is the electrostatic potential, $df/dt = \partial f/\partial t + [\Phi, f]$, $[\Phi, f] = \mathbf{b} \cdot (\nabla_{\perp} \Phi \times \nabla_{\perp} f)$ is the Poisson bracket which indicates the convective derivative, $\mathbf{B} = B\mathbf{b}$ is the magnetic field, $\kappa = \mathbf{b} \cdot \nabla \mathbf{b}$ is the magnetic curvature, $J_{\parallel}$ is the current density parallel to the ambient magnetic field, $\mathbf{v}_e$ is the electron flow velocity, $n = n_e = n_i$ is the particle density, $p = n(T_e + T_i)$ is the total pressure, and T_e and T_i are the electron and ion temperature. Here we assume $T_e \gg T_i$ and T_e is constant along the magnetic field lines. The Boussinesq approximation is used to derive the vorticity equation.

From eq. (3.2), we obtain:

$$\left(\frac{\partial}{\partial t} + \mathbf{v}_e \cdot \nabla \right) n = -n (\nabla_{\parallel} v_{e\parallel} + \nabla_{\perp} \cdot \mathbf{v}_{e\perp}) \quad (3.3)$$

where, the perpendicular component of the electron drift velocity is given from equation of motion of electrons as

$$\mathbf{v}_{e\perp} = \frac{\mathbf{B}}{B^2} \times \mathbf{W}, \quad (3.4)$$

where

$$\mathbf{W} = \nabla \Phi - \frac{\nabla p}{ne} = \nabla \Phi - \frac{T_e}{e} \nabla \ln n, \quad (3.5)$$

and the first and second terms in the R.H.S of eq. (3.5) result in the $\mathbf{E} \times \mathbf{B}$ and diamagnetic drifts, respectively. The second term of R.H.S. of eq. (3.3), is evaluated by use of eq. (3.5) as

$$\begin{aligned} \nabla_{\perp} \cdot \mathbf{v}_{e\perp} &= -\frac{\nabla_{\perp} B^2 \times \mathbf{B}}{B^4} \cdot \mathbf{W} - \frac{\mathbf{B}}{B^2} \cdot \nabla_{\perp} \times \mathbf{W} + \frac{\mathbf{W}}{B^2} \cdot \nabla_{\perp} \times \mathbf{B} \\ &\simeq -\frac{\nabla B^2 \times \mathbf{B}}{B^4} \cdot \mathbf{W}, \end{aligned} \quad (3.6)$$

for the lowest order. (The second and third terms in the R.H.S. of eq. (3.6) represent the inhomogeneity in the direction of $\mathbf{b}$ and the current, respectively. Both remain small corrections.) Then eq. (3.2) reduces to

$$\frac{dn}{dt} = -n \nabla_{\parallel} v_{e\parallel} + n \frac{2}{B} \kappa \times \mathbf{b} \cdot \left(\nabla \Phi - \frac{T_e}{e} \nabla \ln n \right), \quad (3.7)$$

where the relation $\boldsymbol{\kappa} \simeq (\nabla_{\perp} B) / B$ is used.

Now, we calculate the current density at the sheath. Electron flux is given by

$$\Gamma_e = n \int_{\sqrt{2e\Phi/m_e}}^{\infty} dv_{\parallel} v_{\parallel} f(v_{\parallel}), \quad (3.8)$$

where we assume that $f(v_{\parallel})$ is a Maxwellian distribution:

$$f(v_{\parallel}) = \left(\frac{m_e}{2\pi T_e} \right)^{3/2} \exp \left[-\frac{m_e v_{\parallel}^2}{2T_e} \right]. \quad (3.9)$$

Then the current density at the sheath, $J_{\parallel \text{sh}} = enc_s - e\Gamma_e$, is given by

$$J_{\parallel \text{sh}} = en \left(c_s - \frac{v_{\text{the}}}{\sqrt{2\pi}} \exp \left[-\frac{e\Phi}{T_e} \right] \right), \quad (3.10)$$

where $c_s = (T_e/m_i)^{1/2}$ is the ion acoustic speed and $v_{\text{the}} = (T_e/m_e)^{1/2}$ is the electron thermal speed.

Equations (3.1) and (3.7) describe 3D dynamics. The variation of perturbations along the field line is considered weak (because we are interested in the period in which parallel loss remains small), so that they can be simplified by integrating along the magnetic field lines from 0 to $L_{\parallel}$ in the limit where all quantities are assumed to be constant along the magnetic field line except the current. ($L_{\parallel}$ indicates the connection length of the magnetic field line in the SoL.) We note here that the ion dynamics is not solved here; The ion current into the sheath is evaluated by enc_s in eq. (3.10), where the density n is taken constant along the field line in the 2D model. One has a system of equations which are reduced to 2D:

$$L_{\parallel} \frac{d}{dt} \nabla_{\perp}^2 \Phi = \frac{B^2}{m_i n} \{ J_{\parallel}(L_{\parallel}) - J_{\parallel}(0) \} + L_{\parallel} \frac{2B}{m_i n} (\mathbf{b} \times \boldsymbol{\kappa}) \cdot \nabla p \quad (3.11)$$

$$L_{\parallel} \frac{dn}{dt} = -n \{ v_{e\parallel}(L_{\parallel}) - v_{e\parallel}(0) \} + L_{\parallel} n \frac{2}{B} \boldsymbol{\kappa} \times \mathbf{b} \cdot \left(\nabla \Phi - \frac{T_e}{e} \nabla \ln n \right). \quad (3.12)$$

Next, we relate $v_{e\parallel}$ to Φ and n . Assuming that both ends of each field line terminate at grounded conducting plates as is illustrated in fig. 3.1, we have

$$J_{\parallel}(L_{\parallel}) = 2J_{\parallel\text{sh}} = en \left(c_s - \frac{v_{\text{the}}}{\sqrt{2\pi}} \exp \left[-\frac{e\Phi}{T_e} \right] \right) \quad (3.13)$$

$$J_{\parallel}(0) = 0 \quad (3.14)$$

$$v_{e\parallel}(L_{\parallel}) = \left(\frac{v_{\text{the}}}{\sqrt{2\pi}} \right) \exp \left[-\frac{e\Phi}{T_e} \right] \quad (3.15)$$

$$v_{e\parallel}(0) = 0. \quad (3.16)$$

With the help of eqs. (3.13) to (3.16), eqs. (3.11) and (3.12) are reduced to a set of equations for Φ and n . We have

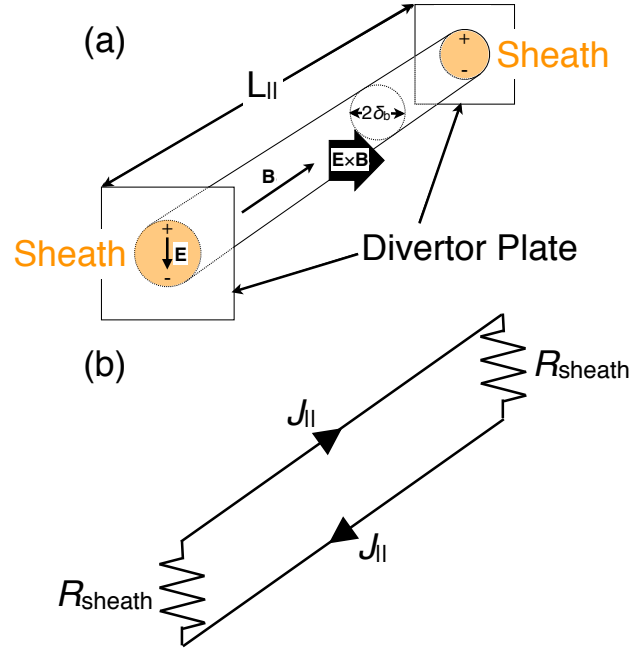


Fig. 3.1 (a) Analytical model of plasma blob propagation mechanisms (reproduction of fig. 2.13) and (b) equivalent circuit.

$$\frac{d}{dt} \nabla_{\perp}^2 \Phi = \frac{1}{L_{\parallel}} \frac{2eB^2}{m_i} c_s \left(1 - \nu \exp \left[-\frac{e\Phi}{T_e} \right] \right) + \frac{2B}{m_i n} (\mathbf{b} \times \boldsymbol{\kappa}) \cdot \nabla p \quad (3.17)$$

$$\frac{dn}{dt} = -n \frac{1}{L_{\parallel}} c_s \nu \exp \left[-\frac{e\Phi}{T_e} \right] + n \frac{2}{B} \boldsymbol{\kappa} \times \mathbf{b} \cdot \left(\nabla \Phi - \frac{T_e}{e} \nabla \ln n \right), \quad (3.18)$$

where $\nu = (m_i/2\pi m_e)^{1/2}$. Introducing the normalization of times to Ω_i^{-1} , length scales to ρ_s , and other quantities to reference values as follows (e.g., separatrix values n_s and T_{es}):

$$\begin{aligned}\Omega_i t &\rightarrow t, & \rho_s \nabla &\rightarrow \nabla, & \rho_s \boldsymbol{\kappa} &\rightarrow \boldsymbol{\kappa}, \\ e\Phi/T_{es} &\rightarrow \Phi, & n/n_s &\rightarrow n, & T_e/T_{es} &\rightarrow T, \\ p/(n_s T_{es}) &\rightarrow p,\end{aligned}$$

where $\Omega_i = eB/m_i$ is the ion cyclotron frequency and $\rho_s = c_s/\Omega_i$ is the gyro-radius. Using the relation

$$\Phi = \Phi_B + \varphi,$$

where the Bohm potential is defined by $\Phi_B \equiv T \ln \nu$, we obtain

$$\frac{d}{dt} \nabla_{\perp}^2 \varphi = \alpha \sqrt{T} \left(1 - \exp \left[-\frac{\varphi}{T} \right] \right) + \frac{2}{n} (\mathbf{b} \times \boldsymbol{\kappa}) \cdot \nabla p \quad (3.19)$$

$$\frac{dn}{dt} = -\alpha n \exp \left[-\frac{\varphi}{T} \right] + 2n \boldsymbol{\kappa} \times \mathbf{b} \cdot (\nabla \varphi - \nabla \ln n), \quad (3.20)$$

where the parameter

$$\alpha = 2 \frac{\rho_s}{L_{\parallel}}$$

measures the conductivity at the plate that governs the net parallel current into divertor plates, and $df/dt = \partial f/\partial t + [\varphi, f]$.

Since the plasma filaments are assumed to be localized perpendicular to the magnetic field, we can use a slab model with orthogonal coordinates $\{x, y, z\}$ such that the x coordinate is taken in the direction of the major radius and the z coordinate is taken along the magnetic field line (fig. 2.14). The dimensionless curvature can be written as

$$\boldsymbol{\kappa} = -\frac{\rho_s}{R} \hat{\mathbf{x}},$$

where R is the major radius, $\hat{\mathbf{x}}$ is the unit vector in the x -direction, so that the second term of the R.H.S. of eq. (3.19) takes a form

$$2 (\mathbf{b} \times \boldsymbol{\kappa}) \cdot \nabla = -\beta \frac{\partial}{\partial y}, \quad (3.21)$$

where

$$\beta = 2 \frac{\rho_s}{R}$$

is the parameter measuring the strength of the curvature drift. Assuming that $|\varphi| \ll 1$ holds, and expanding exponential functions in (3.19) and (3.20) with respect to φ , we obtain

$$\frac{d}{dt} \nabla_{\perp}^2 \varphi = \alpha \varphi - \frac{\beta}{n} \frac{\partial n}{\partial y} + \mu \nabla_{\perp}^4 \varphi \quad (3.22)$$

$$\frac{dn}{dt} = -\alpha n + \beta \left(n \frac{\partial \varphi}{\partial y} - \frac{\partial n}{\partial y} \right) + D \nabla_{\perp}^2 n, \quad (3.23)$$

where $\nabla_{\perp}^2 = \partial^2 / \partial x^2 + \partial^2 / \partial y^2$, and the ion viscosity μ and the diffusion coefficient D are introduced here. We call eqs. (3.22) and (3.23) Hasegawa-Wakatani type equation. The main nonlinearities in the L.H.S. of eqs. (3.22) and (3.23) are the same as those in Hasegawa-Wakatani equations. If one considers the limit of β -terms in R.H.S. to zero, and takes into account adiabaticity terms in Hasegawa-Wakatani equation as substitutes of α terms, this set of equations reduces to Hasegawa-Wakatani equation [5].

We here describe the conservation property of eqs. (3.22) and (3.23). Multiplying $1/n$ to each side of (3.23) to set $N = \log n$, then multiply φ to eq. (3.22) and N to eq. (3.23), we get

$$\varphi \frac{d}{dt} \nabla_{\perp}^2 \varphi = \alpha \varphi^2 - \beta \varphi \frac{\partial N}{\partial y} + \mu \varphi \nabla_{\perp}^4 \varphi \quad (3.24)$$

$$N \frac{dN}{dt} = -\alpha N + \beta N \frac{\partial \varphi}{\partial y} - \beta N \frac{\partial N}{\partial y} + DN (\nabla_{\perp} N \cdot \nabla_{\perp} N) + DN \nabla_{\perp}^2 N, \quad (3.25)$$

where we use the following relation:

$$\begin{aligned}
\nabla_{\perp} N &= \frac{1}{n} \nabla_{\perp} n \\
n \nabla_{\perp} N &= \nabla_{\perp} n \\
\nabla_{\perp} (n \nabla_{\perp} N) &= \nabla_{\perp}^2 n \\
\nabla_{\perp} n \cdot \nabla_{\perp} N + n \nabla_{\perp}^2 N &= \nabla_{\perp}^2 n \\
\frac{1}{n} \nabla_{\perp} n \cdot \nabla_{\perp} N + \nabla_{\perp}^2 N &= \frac{1}{n} \nabla_{\perp}^2 n \\
\nabla_{\perp} N \cdot \nabla_{\perp} N + \nabla_{\perp}^2 N &= \frac{1}{n} \nabla_{\perp}^2 n.
\end{aligned}$$

Subtracting eq. (3.24) from eq. (3.25),

$$\begin{aligned}
N \frac{dN}{dt} - \varphi \frac{d}{dt} \nabla_{\perp}^2 \varphi &= -\alpha N - \alpha \varphi^2 + \beta \left(N \frac{\partial \varphi}{\partial y} - N \frac{\partial N}{\partial y} + \varphi \frac{\partial N}{\partial y} \right) \\
&\quad + DN (\nabla_{\perp} N \cdot \nabla_{\perp} N) + DN \nabla_{\perp}^2 N - \mu \varphi \nabla_{\perp}^4 \varphi,
\end{aligned} \tag{3.26}$$

and gathering β terms using y -derivative, one has

$$\begin{aligned}
N \frac{dN}{dt} - \varphi \frac{d}{dt} \nabla_{\perp}^2 \varphi &= -\alpha N - \alpha \varphi^2 + \beta \left(\frac{\partial (\varphi N)}{\partial y} - \frac{1}{2} \frac{\partial N^2}{\partial y} \right) \\
&\quad + DN (\nabla_{\perp} N \cdot \nabla_{\perp} N) + DN \nabla_{\perp}^2 N - \mu \varphi \nabla_{\perp}^4 \varphi.
\end{aligned} \tag{3.27}$$

Finally, taking the volume average of eq. (3.27), one obtains the relation

$$\begin{aligned}
\frac{d}{dt} \langle H' \rangle &= -\alpha \langle N \rangle - \alpha \langle \varphi^2 \rangle \\
&\quad + D \langle N (\nabla_{\perp} N \cdot \nabla_{\perp} N) \rangle + D \langle N \nabla_{\perp}^2 N \rangle - \mu \langle \varphi \nabla_{\perp}^4 \varphi \rangle,
\end{aligned} \tag{3.28}$$

where

$$H' = \frac{1}{2} |\nabla_{\perp}^2 \varphi|^2 + \frac{1}{2} N^2$$

is the Hamiltonian, and the volume average is defined by the integral as

$$\langle \quad \rangle \equiv \frac{1}{L_x L_y} \int_0^{L_y} \int_0^{L_x} dx dy.$$

Equation (3.28) can be rewritten as

$$\frac{d}{dt} \langle H' \rangle = -\alpha \langle N \rangle - \alpha \langle \varphi^2 \rangle - D \langle (1 - N) |\nabla_{\perp} N|^2 \rangle - \mu \langle \varphi \nabla_{\perp}^4 \varphi \rangle. \quad (3.29)$$

This model conserves the quadratic moment $\langle H' \rangle$, which may be referred as ‘energy’, in the limit of $\alpha \rightarrow 0$, $\mu \rightarrow 0$ and $D \rightarrow 0$. This energy conserving model is abbreviated as “C-model”. We here note the role of transport coefficients α , D and μ for the conservation of the fluctuation energy $\langle H' \rangle$. The viscosity μ causes the decay of $\langle H' \rangle$. The particle diffusivity D (which is originated from collisions, etc.) tends to reduce $\langle H' \rangle$ so long as $N < 1$ holds. In the case of small amplitude fluctuation, N is smaller than unity, and the particle diffusivity reduces $\langle H' \rangle$. The role of α works in two ways. It decreases $\langle H' \rangle$ for finite $|\varphi|$. This is natural, because the finite value of α (i.e. the finite conductivity at end plate) tends to suppress the potential difference. $\alpha \langle N \rangle$ term reduces $\langle H' \rangle$ in time if $\langle N \rangle > 0$. If $\langle N \rangle$ is negative, $\langle H' \rangle$ can increase.

We compare this model with two other models. In the second model, ‘linearized model’, the nonlinear terms in the R.H.S. of eqs. (3.22) and (3.23) are linearized, and are approximated by $(\beta/n) (\partial n / \partial y) \rightarrow \beta \partial n / \partial y$ in eq. (3.22) and $\beta n \partial \varphi / \partial y \rightarrow \beta \partial \varphi / \partial y$ in eq. (3.23). then eqs. (3.22) and (3.23) are modified as follows:

$$\frac{d}{dt} \nabla_{\perp}^2 \varphi = \alpha \varphi - \beta \frac{\partial n}{\partial y} + \mu \nabla_{\perp}^4 \varphi. \quad (3.30)$$

$$\frac{dn}{dt} = -\alpha n + \beta \left(\frac{\partial \varphi}{\partial y} - \frac{\partial n}{\partial y} \right) + D \nabla_{\perp}^2 n. \quad (3.31)$$

This model is abbreviated as “L-model”. In this case, the Hamiltonian reduces to

$$H = \frac{1}{2} |\nabla_{\perp}^2 \varphi|^2 + \frac{n^2}{2}.$$

In the third model, the nonlinear term in the R.H.S. of eq. (3.22) is kept, but the β term on the R.H.S. of eq. (3.23) is ignored, i.e., eqs. (3.22) and (3.23) are

reduced to

$$\frac{d}{dt} \nabla_{\perp}^2 \varphi = \alpha \varphi - \frac{\beta}{n} \frac{\partial n}{\partial y} + \mu \nabla_{\perp}^4 \varphi, \quad (3.32)$$

$$\frac{dn}{dt} = -\alpha n + D \nabla_{\perp}^2 n. \quad (3.33)$$

This ‘simplified model’ is the same as that used by Aydemir [67] and is abbreviated as “S-model”. However, the energy conservation does not hold in this model.

Considering the important role of nonlinear terms and the construction of conservation relation, we conclude here that eqs. (3.22) and (3.23) are more relevant model in order to compare with experimental observations. Before going to the step of solving eqs. (3.22) and (3.23), we explain roles of some of nonlinear terms in the Appendix A. Simplified model in [67] is also referred to.

In summary, a set of basic equations (3.22) and (3.23) is derived in this chapter for the study of blob transport problems. The model equation is characterized by dimensionless parameters α and β , $\alpha = \rho_s / L_{\parallel}$ and $\beta = \rho_s / R$, where α represents the impact of finite conductivity at the sheath and β denotes the bad magnetic curvature. In addition, particle diffusivity and viscosity by collisional process, D and μ , are introduced. The model system is controlled by these four parameters. The results of analysis are explained in Chapters 4 and 5. In Chapter 4, an evolution of a single blob (figs. 3.2 and 3.3) is studied. In Chap 5, the selfconsistent evolution of blob in the flux-driven nonlinear simulation is studied in the geometry like fig. 3.4.

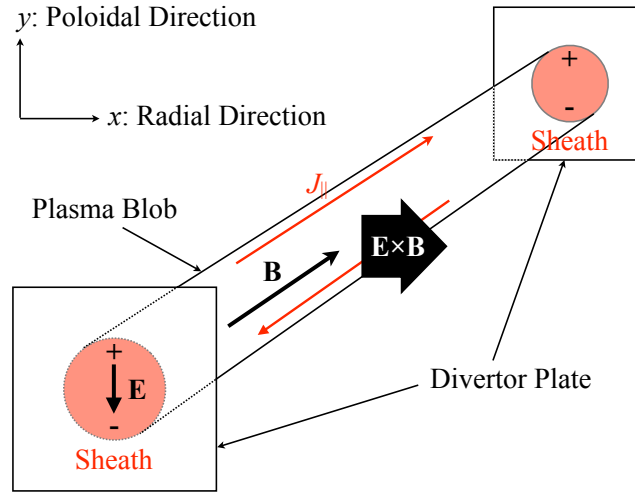


Fig. 3.2 Theoretical model and geometry of single blob. Reproduction of fig. 2.13.

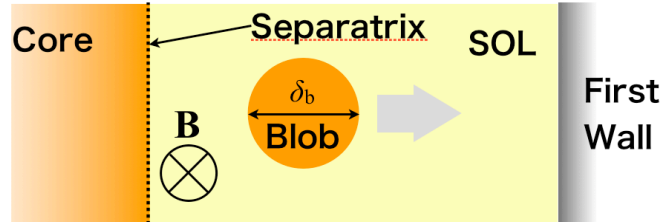


Fig. 3.3 2-dimensional simulation field for single isolated blob simulation.

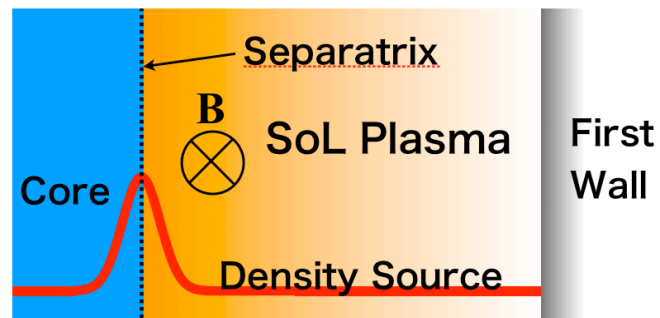


Fig. 3.4 2-dimensional simulation field for flux-driven-turbulence simulation.

Chapter 4

Single Blob Dynamics

In the previous chapter, theoretical picture and equations in 2D model are explained. The concept of ‘blob’ is introduced as an isolated density clump in a magnetic-field-aligned flux tube in the SoL, connected to the divertor plates at each end (fig. 4.1). Owing to its rapid velocity in the radial direction, the

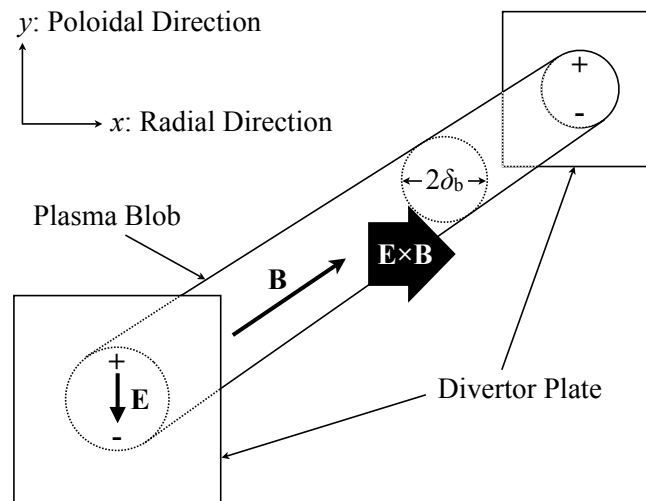


Fig. 4.1 Theoretical model and geometry of blob. Reproduction of fig. 2.13.

approximation to treat it in two-dimensional model is relevant to the real three-dimensional system for the basic understanding of transport in SoL plasma. With this theoretical model, the single blob propagation is investigated. In this chapter we study in detail the evolution of a single blob [6, 7].

We first explain the simulation conditions, and then models (C, L, S in Chapter 3) are compared. Focusing on the conservation model, we study the velocity of the blob. Evolution of the blob size and relation with the velocity are obtained.

4.1 Simulation Conditions

The simulations are performed with the initial value problem. Initial condition for the blob density distribution is taken to be a 2D Gaussian:

$$n_b \exp \left[\left(-\frac{x - 2\delta_b}{\delta_b} \right)^2 - \left(\frac{y - L_y/2}{\delta_b} \right)^2 \right],$$

which exists in a ambient plasma with constant density. Here, δ_b indicates the initial blob size, and L_y is the system size in the y direction (fig. 4.2). The

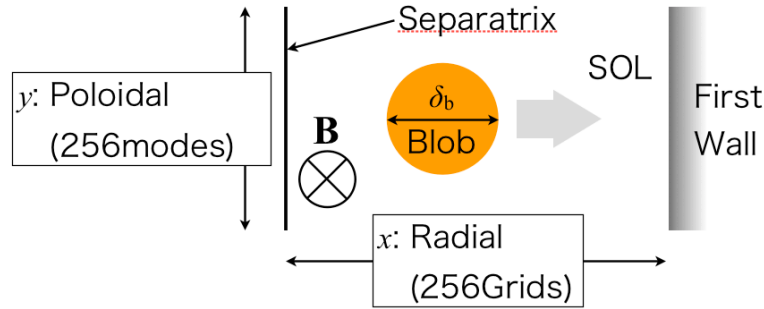


Fig. 4.2 Simulation field, initial conditions, and geometry.

simulation parameters are set as $\alpha = 3 \times 10^{-5}$, $\beta = 6 \times 10^{-4}$, $n_b/n_0 = 10$, and $\mu = D = 2 \times 10^{-3}$ [67], where n_0 is the initial background density. (Note that all quantities are normalized as explained in Chapter 3.)

Discretization for numerical simulation is done in Fourier expansion in y -direction, and the Fourier mode number is abbreviated as m . The following

boundary conditions are used for simulations. In the vorticity equation

$$\begin{aligned} \text{for } m = 0 \text{ mode:} \quad & \nabla_{\perp}^2 \varphi' \big|_{x=0} = 0, & \nabla_{\perp}^2 \varphi \big|_{x=L_x} = 0. \\ \text{for } m \neq 0 \text{ mode:} \quad & \nabla_{\perp}^2 \varphi \big|_{x=0} = 0, & \nabla_{\perp}^2 \varphi \big|_{x=L_x} = 0. \end{aligned}$$

where the prime indicates the derivative in the x -direction and L_x indicates the system size in the x -direction. Similarly,

$$\begin{aligned} \text{for } m = 0 \text{ mode:} \quad & \varphi' \big|_{x=0} = 0, & \varphi \big|_{x=L_x} = 0. \\ \text{for } m \neq 0 \text{ mode:} \quad & \varphi \big|_{x=0} = 0, & \varphi \big|_{x=L_x} = 0. \end{aligned}$$

are imposed. In the continuity equation,

$$\begin{aligned} \text{for } m = 0 \text{ mode:} \quad & n \big|_{x=0} = n_0, & n' \big|_{x=L_x} = 0. \\ \text{for } m \neq 0 \text{ mode:} \quad & n \big|_{x=0} = 0, & n' \big|_{x=L_x} = 0. \end{aligned}$$

are imposed. The periodic boundary conditions are used in the y -direction. 256 radial grids and 256 Fourier modes are taken for the discretization. Simulation field size depends on the initial blob size δ_b .

4.2 Model Comparisons

Before going into intensive simulations, we compare Conserved (C), Linearized (L), and Simplified (S) models. Figure 4.3 shows the time evolution of the radial position of the peak of a single blob, its velocity, and the peak value of the poloidally-averaged density in the case with $\delta_b = 5$ for three models. Initial conditions are identical for all three cases. In the early stage of propagation, the velocity of the ‘peak’ increases ($0 < t \lesssim 100$ for the C-model and the S-model, and $0 < t \lesssim 200$ for the L-model). In this phase, the peak position of the density shifts forward in the radial direction from the initial Gaussian shape, keeping the foot position. This phase is defined as ‘initial phase’. In the initial phase, the density also increases. This indicates that the blob density

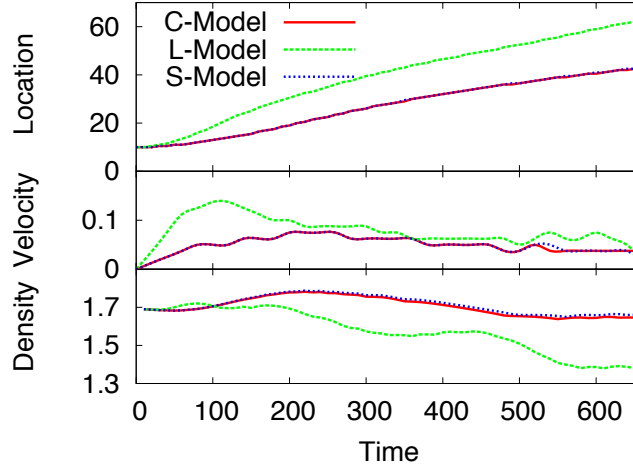


Fig. 4.3 A time evolution of the peak location (x -coordinate), the peak velocity, and the peak value of poloidally averaged density using three models. The length and time are normalized in ρ_s and Ω_i .

concentrates at the peak position. After the initial phase, the blob spreads out and the propagation velocity approaches to the stationary value ($t \gtrsim 200$ for every model). This phase is defined as ‘stationary phase’. The S-model agrees well with the C-model, while the L-model is quite different from the C-model. The density’s nonlinearity in eq. (3.22) has a more dominant effect than the one in eq. (3.23) in the phase of nonlinear evolution. Hence, we focus on the C-model and the S-model in the following.

Figure 4.4 shows the 2D contour plots of density for a blob (with initial parameter of $\delta_b = 5$) in the C-model and the S-model, respectively. The difference between the initial phase and the stationary phase for both cases is identified. According to the initial condition, the blob density is given by the Gaussian shape ($t = 0$). In the initial phase, it concentrates at the front end of the blob ($t \simeq 600$). Then, it is broken into pieces by the Kelvin-Helmholtz (K-H) instability ($t \simeq 1900$). The difference between the C-model and the S-model is not remarkable in the contour plots up to $t \sim 1200$; however, the difference

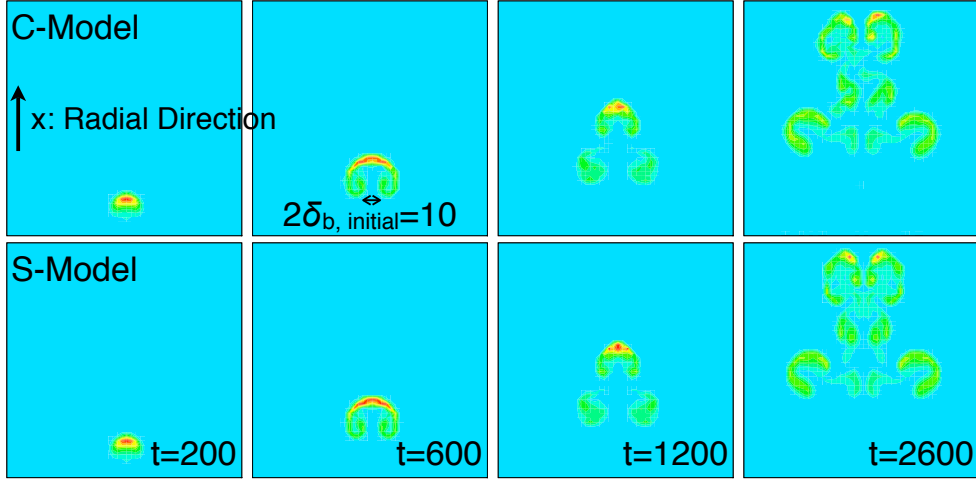


Fig. 4.4 Density contour plots of a blob with $\delta_b = 5$ propagation time evolution. The C-model is used in upper plots and the S-model in lower plots. System size is 128×128 . The length and time are normalized in ρ_s and Ω_i . Blobs break and spread with the Kelvin-Helmholtz instability.

becomes noticeable in the time evolution of the kinetic energy. Figure 4.5 shows the time evolution of the kinetic energy $\langle |\nabla_\perp \varphi|^2 \rangle$, for a blob with $\delta_b = 5$ in the C-model and the S-model, respectively. In the linearly growing phase ($t < 500$),

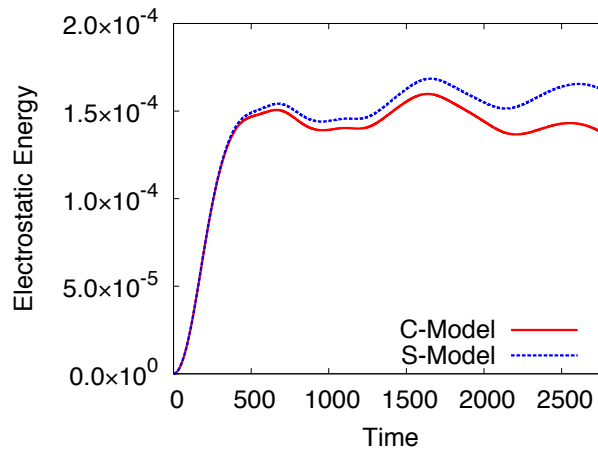


Fig. 4.5 The time evolution of the electrostatic energy $|\nabla_\perp \varphi|^2$ for blob with $\delta_b = 5$ propagation in the C-model and S-model.

the energy in one model agrees with the other. In the saturation phase, the difference becomes larger. The β term in the eq. (3.23) has small effect in the linearly growing phase. In the nonlinear phase, the fine structure appears, so that a steep gradient like $\partial/\partial y \sim 1$ can occur. Thus, the terms which include the y -dependence have this effect on the nonlinear evolution. In this phase, the symmetry in the y -direction is broken by the density's nonlinearity. The numerical issues concerning on models C and L are discussed in Appendix A.

4.3 Single Blob Evolution and Propagation Velocity Analysis

4.3.1 Deformation of blobs

In the following, we focus on the analysis by the C-model. During the nonlinear evolution, the blobs are broken due to the instabilities. As has been pointed out in the reference [67], the Kelvin-Helmholtz (K-H) instability appears for a small size blob (due to the velocity shear accompanying its rapid convection); on the other hand, the interchange instability appears for a large blob. The critical blob size, which categorizes the type of instability, is given as

$$\delta_{\text{crit}} = \left(\frac{L_{\parallel}}{2\rho_s} \right)^{1/5} \left(\frac{L_{\parallel}}{R} \right)^{1/5} \rho_s. \quad (4.1)$$

The K-H instability appears for $\delta_b < \delta_{\text{crit}}$, while interchange mode grows for $\delta_b > \delta_{\text{crit}}$. For our parameters explained in the beginning of section 4.1, the dimensionless characteristic length is estimated as $\delta_{\text{crit}} \sim 15$ (in the unit of ρ_s). Figure 4.6 shows the density contour plots for various values of δ_b , and indicates the characteristics of instabilities for various blob sizes. For the blob with $\delta_b = 10$, the K-H instability appears, and the evolution is similar as the one in fig. 4.4. The larger blob with $\delta_b = 25$ propagates in the radial direction with keeping its round structure. For much larger blobs with $\delta_b = 40$ and

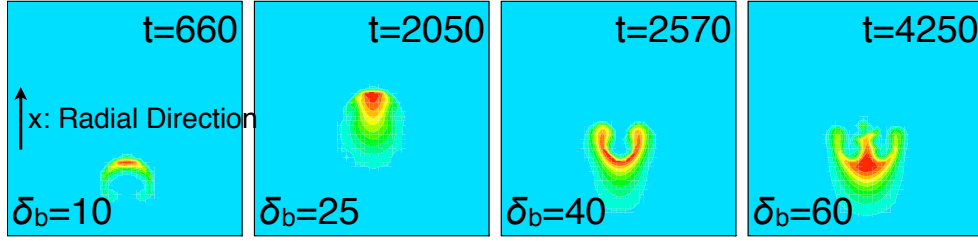


Fig. 4.6 Snap shots of density contour of blobs with $\delta_b = 10, 25, 40$, and 60 in the C-model. System sizes are 192×192 , 256×256 , 384×384 , and 512×512 from the left picture. Blobs break and spread with the Kelvin-Helmholtz and interchange instability or propagate in the radial direction with keeping its coherent structure.

$\delta_b = 60$, the Rayleigh-Taylor (R-T) fingers develop, which indicate the growth of the interchange instability. The number of the R-T fingers depends on the unstable wave number and the blob size. (The wavelength of the interchange instability, which determines initial number of R-T fingers, is explained in the next section.)

4.3.2 Velocity of blobs

We study here how these modifications of blobs shape influence the velocity of blobs in the stationary phase. In the stationary phase, small blobs break up and spread due to K-H instability. The velocity of the peak of $n(x, y)$ almost saturates in the stationary phase. Figure 4.7 shows the plot of the initial blob size and the velocity of the peak in the stationary phase. The solid line represents the results calculated by the theoretical blob velocity

$$V_x = \frac{(2\beta/\alpha)}{\delta_b^2}. \quad (4.2)$$

This formula was derived by assuming a traveling waveform [67], in which, the traveling wave form solution $n(x, y, t) = n(x - V_x t, y)$ is assumed to solve the continuity equation (3.23) analytically. Equation (4.2) does not have the

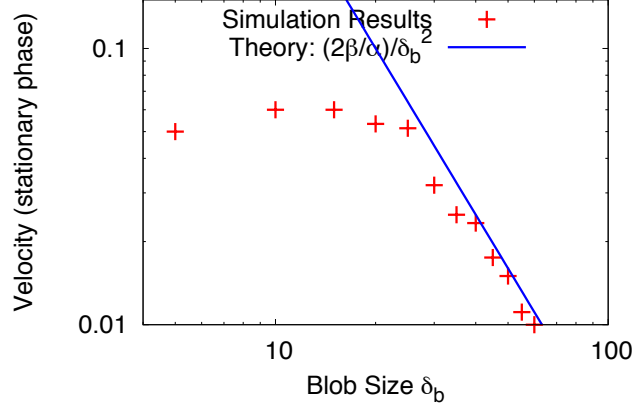


Fig. 4.7 The velocity of the peak of blob in their stationary phase v.s. initial blob size using the C-model. A log scale is used on the vertical and horizontal axes.

dependence on the blob density, while eq. (2.18) contains the blob density n_b . This is because the evaluation $n_b/n_T \rightarrow 1$ is employed in 2D model. During the blob moves at large velocity V_x a radial distance Δ_x (eq. (2.19)) before the loss along the field line eliminates the blob, it is valid to assume that quantities except the current is uniform along the field line (recall the procedure in Chapter 3). Thus, $n_b/n_T \rightarrow 1$ in eq. (2.18) can be employed. Equations (2.18) and (4.2) are consistent during the blob life time Δ_x/V_x .

In the regime where the blob is subject to the K-H instability, the simulation results are quite different from the theoretical prediction eq. (4.2). Blobs spread out into a broader area, and the assumption of the traveling waveform is violated, then the velocity saturates at a value lower than the predicted one. On the other hand, the theoretical estimation of blob propagation velocity in the stationary phase agrees well with the simulation results for larger blobs, since the spreading is weak in the interchange regime. It is captured that propagation velocity is proportional to δ_b^{-2} in this regime, $\delta_b > \delta_{\text{crit}}$. (The critical size δ_{crit} depends on simulation conditions such as curvature and sheath resistivity

and so on, as is explained in the beginning of Subsection 4.3.1.)

So far, the peak location is used to estimate the blob velocity. The velocity of peak is suitable to capture the blob characteristics in the interchange regime. However, it does not contain the information of the blob behavior as a whole. In order to take this into account, the center-of-mass motion is introduced. The coordinate of the center-of-mass is defined as

$$x_{\text{CM}}(t) = \frac{\int x \cdot n_{\text{blob}}(x) dx}{\int n_{\text{blob}}(x) dx}, \quad (4.3)$$

and the radial velocity of the center of mass is defined as

$$v_{\text{CM}} = \frac{\partial x_{\text{CM}}}{\partial t}, \quad (4.4)$$

where $n_{\text{blob}}(x)$ is the radial profile of blob density which is poloidally averaged, and is obtained from the difference between the instantaneous total density profile and the initial background density profile. Figure 4.8 shows the time evolution of v_{CM} of the various blob sizes. It is found that the blob with $\delta_b = 5$

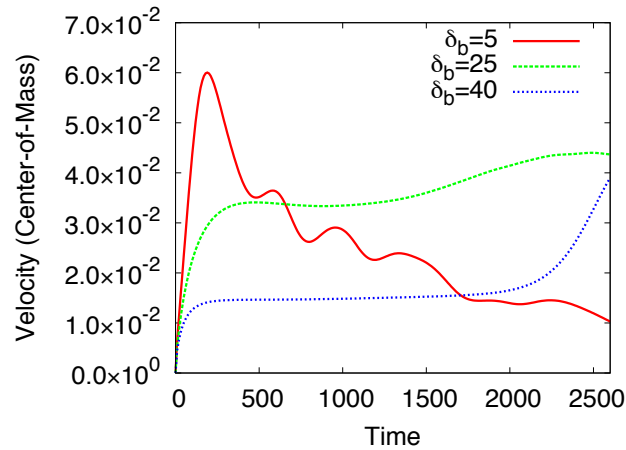


Fig. 4.8 Time evolution of the propagation velocity of blobs with $\delta_b = 5$, 25, and 40 estimated by center-of-mass.

changes the propagation velocity dynamically. The velocity increases initially.

In this phase, the blob density concentrates to the front end as is given in fig. 4.4. The local minimum of the curve in the phase of gradual reduction ($t \gtrsim 500$) corresponds to the vortices generation and the upslope is related to the detachment of the vortices (see fig. 4.4). In the evolution of K-H instability, the blob breaks and spreads out. Hence, the velocity gradually decreases with the time evolution. The blob with $\delta_b = 25$ is less affected by instabilities, and the velocity increases gradually. For the blob with $\delta_b = 40$, the velocity is almost stationary till the growth of the interchange instability takes place. Once it appears, the velocity grows up with R-T fingers. From fig. 4.8 it is seen that the decay time of v_{CM} is roughly estimated as $t \sim 1500$ for the blob with $\delta_b = 5$ which is affected by K-H instability. During the decay time, the small blob propagates, up to the half scale of system size.

The dependence of the blob velocity on the blob size, which is illustrated in fig. 4.7, is compared with experimental observation by the TORPEX group in EPFL/CRPP (Switzerland) [118]. In fig. 4.9, the contour map shows their experimental observation. The color indicates the probability of observations: that is, data in the red domain are frequently observed, while the observation is very rare in the region shown by blue. The red points show our simulation result in fig. 4.7, which were chosen as representative theoretical-simulative work. The velocity and size are normalized by using TORPEX device parameters. (See [118] for the details of normalization.) We can see that the simulation result gives good estimation for the experimental observation.

4.3.3 Evolution of blob size

Next, the variance is introduced to estimate the spread of the blob quantitatively. The variance is defined as

$$\sigma^2(t) = \frac{\int (x - x_{CM})^2 n_{\text{blob}}(x) dx}{\int n_{\text{blob}}(x) dx} \quad (4.5)$$

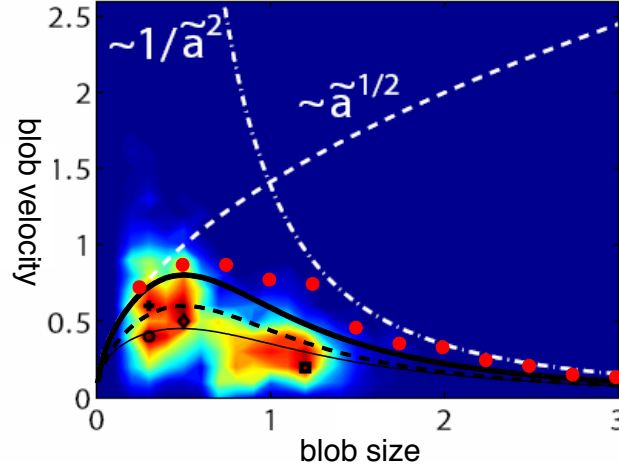


Fig. 4.9 The dependence of blob velocity on blob size, which was observed in TORPEX toroidal device is presented. Comparison between experimental observation (contour map) and 2D simulation (red points) is shown. Data of fig. 4.7 are reproduced and are shown by red points. The normalized blob size is denoted by $\tilde{a}$, definition of which is given in [118]. (Reproduced from [118])

using the center-of-mass coordinate x_{CM} . Comparing the variance definition for the Gaussian distribution function and the blob initial density profile, it is shown that $\sqrt{2}\sigma$ corresponds to the effective blob size in the initial state $r_b(t)$, where $r_b|_{t=0} = \delta_b$. Figure 4.10 shows the time trace of v_{CM} on the variance $\sqrt{2}\sigma$ for the blob with initial size $\delta_b = 5$. The velocity of the peak v_{peak} for δ_b and the theoretical value eq. (4.2) are also plotted. It is seen that according to the temporal evolution, the variance is getting large and the velocity approached to the theoretical line although the assumption of the traveling waveform is already not valid in this phase.

At last, in order to analyze the case where multiple blobs exist (as is the case of next chapter), the blob size ensemble average of physical quantity Q is

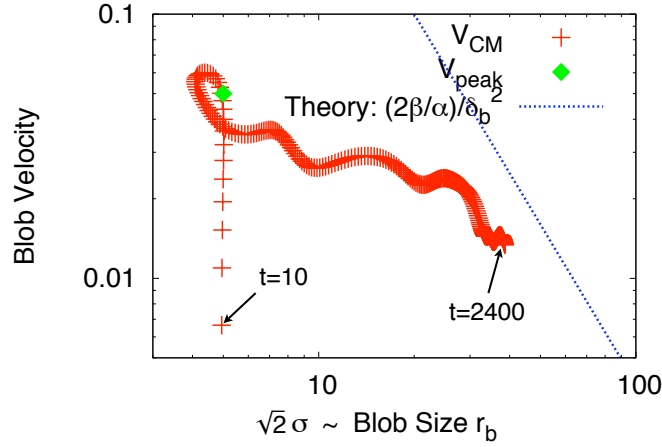


Fig. 4.10 Similar as fig. 4.6, for a blob with $\delta_b = 5$ the peak velocity and the theoretical prediction are plotted. Red crosses indicate the time trace of v_{CM} versus $\sqrt{2}\sigma$.

introduced by

$$\langle Q \rangle = \frac{\int_1^{\delta_{\max}} d\delta_b \delta_b f(\delta_b) Q}{\int_1^{\delta_{\max}} d\delta_b \delta_b f(\delta_b)}, \quad (4.6)$$

according to the literature [58], where $f(\delta_b)$ is a size distribution of blobs and the lower bound of integration is cut off at the gyro-radius $\delta_b = 1$. The upper cut off is determined by the requirement that the blob moves fast enough before decaying to contribute to transport. As is the discussion in Subsection 2.3.1, $\delta_{\max}$ is discussed by the consideration that the distance, which blob propagates with the speed eq. (4.2) in the life time $L_{\parallel}/c_s$, is equal to $\delta_{\max}$. That is,

$$\delta_{\max} = \left(\frac{\beta}{\alpha^2} \right)^{1/3} = \left(\frac{L_{\parallel}}{2\rho_s} \right)^{1/3} \left(\frac{L_{\parallel}}{R} \right)^{1/3}. \quad (4.7)$$

Then, the outward particle flux Γ is estimated as

$$\Gamma(t) = \left\langle \frac{N_b(\delta_b)}{S(r_b(\delta_b, t))} v_{CM}(\delta_b, t) \right\rangle, \quad (4.8)$$

where N_b is the total number of particles in the blob and $S(r_b(\delta_b, t))$ is the effective area of the spreading blob. The quantity N_b is calculated from the

initial blob density profile as

$$N_b(\delta_b) = \int_x \int_y n_b \exp \left[- \left(\frac{X}{\delta_b} \right)^2 - \left(\frac{Y}{\delta_b} \right)^2 \right] dX dY \quad (4.9)$$

and the effective area S is approximated as

$$S(r_b) \approx \pi (2r_b)^2. \quad (4.10)$$

Figure 4.10 shows that v_{CM} asymptotically approaches to the theoretical value of eq. (4.2). Using the asymptotic relation

$$r_b \sim \frac{2\beta/\alpha}{v_{CM}}, \quad (4.11)$$

we obtain

$$\Gamma \sim \frac{\alpha}{4\pi\beta} \langle N_b(\delta_b) v_{CM}^2(\delta_b) \rangle. \quad (4.12)$$

This formula might be applicable to estimate the flux which is averaged over blob size ensemble.

Systematic study of the evolution of blob size is performed for the parameter of $\mu = D = 2.0 \times 10^{-2}$. The choice of this set of dissipation parameters is motivated by the selfconsistent simulation in Chapter 5. Figure 4.11 illustrates the long time evolution of blobs for the case of $\delta_{init} = 5, 15$, and 25 . The time evolution of the size, defined by the variance, is summarized in fig. 4.12. For blobs, which are smaller than δ_{crit} in the initial phase, approaches to a common value δ_{crit} .

4.4 Summary of Chapter 4

We investigated the motion of a single blob by solving eqs. (3.22) and (3.23), i.e., the Hasegawa-Wakatani-type model equation as an initial value problem. It was found that the density's nonlinearity in the vorticity equation has a dominant effect on the blob propagation and that the nonlinearity in the continuity

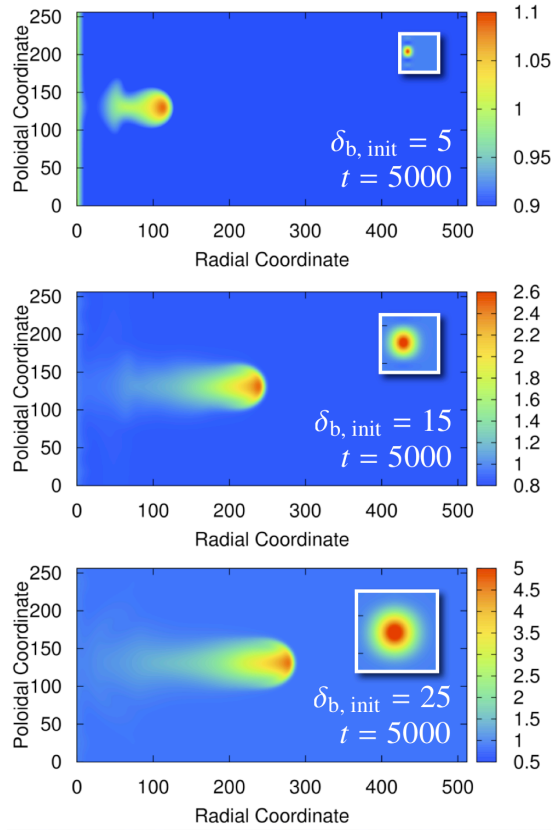


Fig. 4.11 Long time evolution of blobs whose initial sizes are different. $\mu = D = 2.0 \times 10^{-2}$ in these cases.

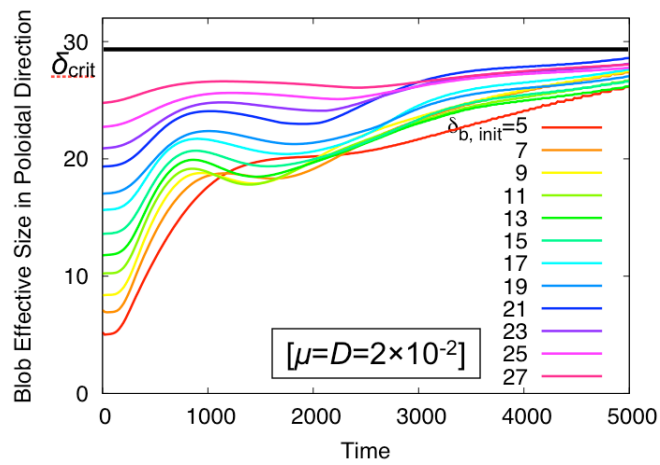


Fig. 4.12 Time evolution of effective blob size in the y -direction for various initial sizes. $\mu = D = 2.0 \times 10^{-2}$ in these cases.

equation causes the asymmetry in the poloidal direction. The dependence of the propagation velocity on the blob size was investigated. When the initial blob size δ_b is large, the propagation velocity scales as δ_b^{-2} . In contrast, when the initial blob size is small, the blob is broken into the pieces due to the Kelvin-Helmholtz instability, and the size approaches to δ_{crit} . As a result of it, the velocity of small blob is limited to a velocity which is of the order of V_x at $\delta = \delta_{\text{crit}}$.

Chapter 5

Selfconsistent Blobby Transport

In this chapter, the transport driven by interchange turbulence in SoL plasmas is investigated by performing a flux-driven nonlinear simulation of a two-dimensional Hasegawa-Wakatani-type model equations ((3.22) and (3.23)) [8,9]. The selfconsistent evolution that generates blob-like perturbation is confirmed. The velocity of the generated blob is evaluated from the slope of the autocorrelation function in the quasi-steady state. It is found that the velocity agrees with that estimated by the velocity model for a single blob whose size corresponds to that of a typical mesoscale structure. In Section 5.1, the model for SoL interchange turbulence is explained. In Section 5.2 a linear analysis is performed. In Section 5.3, numerical simulation results are discussed. The formation of a mesoscale structure in the initial nonlinear stage is explained. In Section 5.4, the transport is evaluated in the quasi-steady state before the mesoscale structure breaks up owing to the boundary effect. In Section 5.5, the ballistic propagation velocity is characterized using the velocity model for a single blob whose size corresponds to that of a typical mesoscale structure. In Section 5.6, the scaling of the transport coefficient is discussed. The effective diffusion coefficient is found to have ‘Bohm-like’ diffusion.

5.1 Model

The model geometry, which we analyze here, is illustrated in fig. 5.1. Two dimensional dynamics is solved in radial and poloidal planes (i.e., x, y -coordinates). The system is periodic in the poloidal (y) direction. The plasma is bounded at $x = 0$ (core-SoL interface) and at $x = L_x$ (wall). The source is prescribed near $x \simeq 0$ (boundary), and the flux-driven direct nonlinear simulation is done by solving eqs. (3.22) and (3.23).

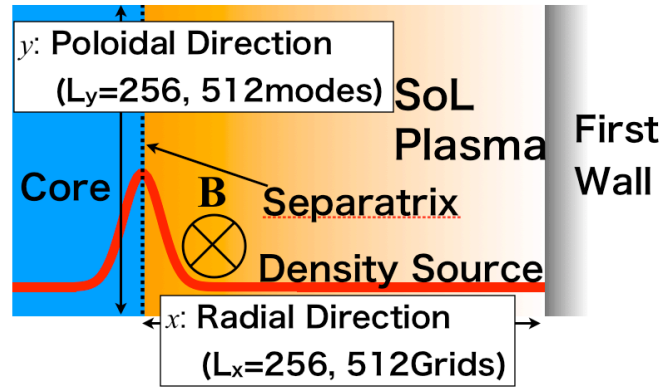


Fig. 5.1 Simulation field and geometry

The following boundary conditions are imposed in the x -direction. Discretization for numerical simulation is done in Fourier expansion in the y -direction, and the Fourier mode number is abbreviated as m . In the vorticity equation [eq. (3.22)]:

$$\begin{array}{lll}
 m = 0 \text{ mode:} & \varphi'|_{x=0} = 0, & \varphi|_{x=L_x} = 0, \\
 & \nabla_{\perp}^2 \varphi'|_{x=0} = 0, & \nabla_{\perp}^2 \varphi|_{x=L_x} = 0, \\
 m \neq 0 \text{ mode:} & \varphi|_{x=0} = 0, & \varphi|_{x=L_x} = 0, \\
 & \nabla_{\perp}^2 \varphi|_{x=0} = 0, & \nabla_{\perp}^2 \varphi|_{x=L_x} = 0.
 \end{array}$$

Here the prime indicates the derivative with respect to x and L_x indicates the system size in the x -direction. These boundary conditions for the vortic-

ity equation are same as the conditions used in the isolated blob propagation simulation in Chapter 4. Likewise, in the continuity equation [eq. (3.23)], the conditions

$$\begin{array}{lll} m = 0 \text{ mode:} & n'|_{x=0} = 0, & n|_{x=L_x} = 0, \\ m \neq 0 \text{ mode:} & n|_{x=0} = 0, & n|_{x=L_x} = 0 \end{array}$$

are imposed. A periodic boundary condition is used in the y -direction such that $\varphi|_{y=0} = \varphi|_{y=L_y}$ and $n|_{y=0} = n|_{y=L_y}$, where L_y indicates the system size in the y -direction.

The initial condition is chosen that $n(x, y) = 1$ at $t = 0$, and that $\varphi(x, y)$ is given by a small random noise.

A Gaussian type density source term is introduced at the $x = 0$ boundary, which is given by

$$S(x) = S_0 \exp \left[\frac{-x^2}{\lambda^2} \right],$$

where $S_0 = 5 \times 10^{-3}$ and $\lambda = 8$.

The simulation box size is set up as $L_x = L_y = 256$ (in units of ρ_s). For the dissipation parameters, $\mu = D = 2 \times 10^{-2}$ are assumed as default values. For the characteristic parameters α and β , three cases are investigated:

$$\begin{array}{llll} \text{a reference case with} & \alpha = 3 \times 10^{-5} & \text{and} & \beta = 6 \times 10^{-4}, \\ \text{a low-}\alpha \text{ case with} & \alpha = 3 \times 10^{-6} & \text{and} & \beta = 6 \times 10^{-4}, \\ \text{a high-}\beta \text{ case with} & \alpha = 3 \times 10^{-5} & \text{and} & \beta = 6 \times 10^{-3}. \end{array}$$

These values are obtained by taking example of parameters $c_s = 6 \times 10^4$ [m/s], $\Omega_i = 10^8$ [s⁻¹], $L_{\parallel} = 40$ [m], and $R = 2$ [m] for the reference case [67].

The numerical method is based on the finite difference in the x -direction and Fourier expansion in the y -direction. Additional explanation of the numerical scheme is given in the Appendix B.

5.2 Linear Stability of Interchange Mode in SoL Plasma

The interchange instability in SoL plasma plays an important role in generating plasma blobs at the separatrix and in breaking up large blobs during their propagation [67]. The interchange instability is responsible for SoL turbulence and transport. Here, we investigate the linear stability of the interchange mode in the SoL. In linearizing eqs. (3.22) and (3.23), the nonlinear density terms are reduced to $(\beta/n)(\partial n/\partial y) \rightarrow \beta \partial n/\partial y$ in eq. (3.22) and $\beta n \partial \varphi/\partial y \rightarrow \beta \partial \varphi/\partial y$ in eq. (3.23). We assume that the mean density is given by

$$n_0(x) = \tilde{n}_0 \exp(-x/\lambda),$$

where λ is the radial width of the SoL density, and we introduce perturbations of the form

$$\tilde{n}(x) \exp[i(k_y y - \omega t)].$$

The following dispersion relation is obtained:

$$\begin{aligned} (i\omega k_\perp^2 - \alpha - \mu k_\perp^4) (-i\omega + \alpha + ik_y \beta + Dk_\perp^2) \\ + ik_y \beta \left(-\frac{ik_y}{\lambda} + ik_y \beta \right) = 0, \end{aligned} \quad (5.1)$$

where $k_\perp^2 = k_x^2 + k_y^2$, and k_x and k_y are the wave numbers in the x - and y -directions, respectively. Solving the dispersion relation for ω and taking the imaginary part, the linear growth rate of the interchange instability in SoL plasma is obtained. Figures 5.2 to 5.4 shows the linear growth rate of the interchange mode in SoL plasma as a function of k_y . It is found that a larger value of α stabilizes the long-wavelength modes and a larger value of β destabilizes modes. The wave number that gives the maximum growth rate weakly depends on β , but is affected much by α , D and μ . A larger value of α indicates the higher sheath conductivity at the divertor plates. If α is large, φ is kept small

in order to sustain the cancellation current for the polarization. This means that a larger value of α is effective for the stabilization of the interchange instability. The stabilizing effect of α is also explained in Chapter 3. On the other hand, β is the curvature parameter and it drives the interchange instability. The radial width of the SoL density $\lambda = 32$ is assumed in the linear analysis in figs. 5.2 - 5.4, while the characteristic scale length in the nonlinear simulations in the following sections is found $\lambda \sim 8$. The parameter dependency discussed in this section is relevant in understanding nonlinear simulations. Detailed discussion on the linear stability analysis of the interchange mode in SoL plasma is discussed in Appendix C.

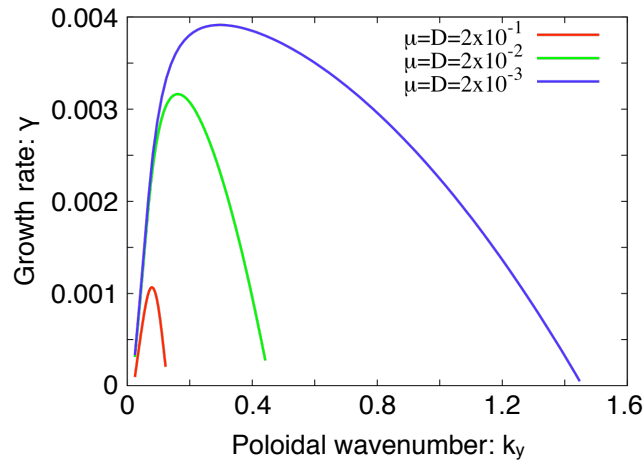


Fig. 5.2 Linear growth rate of SoL interchange mode as a function of the poloidal wave number k_y . Dissipation parameter dependence is shown. $\lambda = 32$, $\alpha = 3 \times 10^{-5}$, $\beta = 6 \times 10^{-4}$ are assumed.

5.3 Formation of Mesoscale Strucutre

First, we will discuss the spatial structure in the initial stage of the reference case. Figure 5.5 illustrates the evolution of n and φ in a linear phase. The perturbation which has the largest linear growth rate starts to grow. The wave

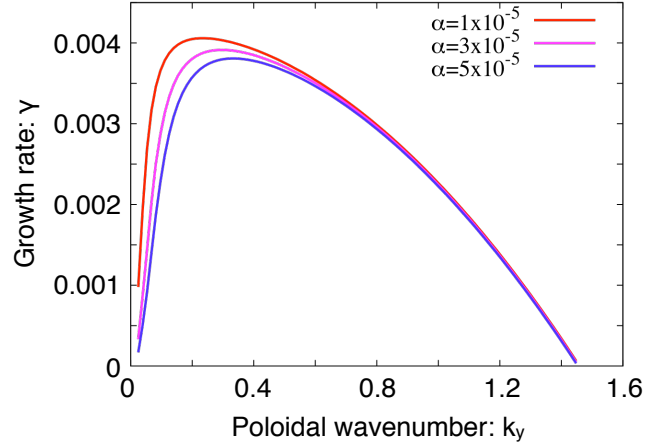


Fig. 5.3 Linear growth rate of SoL interchange mode as a function of the poloidal wave number k_y . α -dependence is shown. $\lambda = 32$, $\beta = 6 \times 10^{-4}$, $D = \mu = 2 \times 10^{-3}$ are assumed.

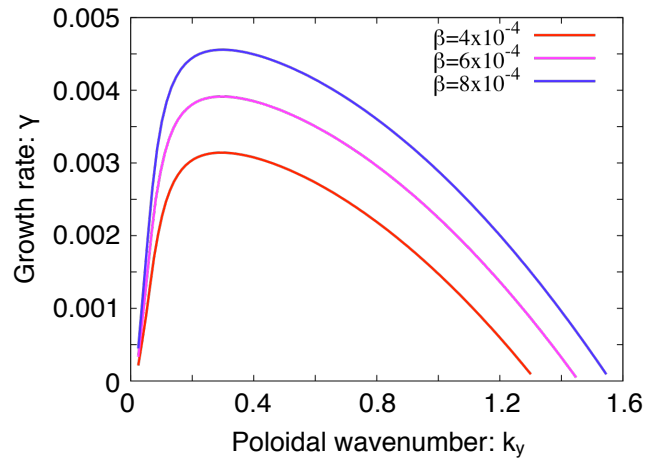


Fig. 5.4 Linear growth rate of SoL interchange mode as a function of the poloidal wave number k_y . β -dependence is shown. $\lambda = 32$, $\alpha = 3 \times 10^{-5}$, $D = \mu = 2 \times 10^{-3}$ are assumed.

length of the excited mode is given by the most unstable mode. The nonlinear

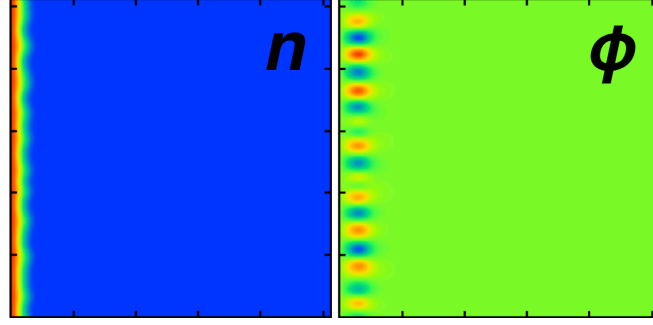


Fig. 5.5 Snap shot of density contour. Time is $t = 1400$. Horizontal axis is radial direction x and vertical axis is poloidal direction y . Length of each axes is $256[\rho_s]$ and time is in units of Ω_i^{-1} .

evolution is described in fig. 5.6. Figure 5.6 shows contour plots of the particle density n , electrostatic potential φ , and vorticity $\nabla_{\perp}^2 \varphi$ fields at times $t = 1800$, 2000, 3000, and 5500. Ten Rayleigh-Taylor (R-T) fingers develop at $t \sim 1800$. The number of fingers corresponds to the poloidal mode number that gives the maximum linear growth rate. For the simulation parameters, the poloidal mode number $m = 10$ corresponds to $k_y \sim 0.12$. Fingers detach from the separatrix ($x = 0$ boundary) because of the velocity shear due to the gradient of the electric field. As a result, blobs are formed at $t \sim 2000$.

During the rapid propagation in the radial direction, the shapes of the blobs are deformed. The instability that deforms blobs in this simulation is concluded to be the Kelvin-Helmholtz (K-H) instability owing to the following reasons. According to the study of the evolution of single blobs, the dominant mechanism for deformation is the K-H instability for small blobs and the interchange instability for large blobs. These types of instability can be distinguished by observing the deformation pattern. For the former (K-H), the deformation pattern is stirring vortices over the whole blob. For the latter (interchange), the

deformation takes place at the front side of propagation (i.e., the bad curvature side) of the blob. An additional criterion determining the type of stability is the size of the blob. The K-H instability occurs if the size of the blob is smaller than δ_{crit} , and the interchange instability occurs if the size is larger than δ_{crit} . δ_{crit} was evaluated in ref. [67] as $\delta_{\text{crit}} \sim 2(\beta/\alpha^2)^{1/5}$, and δ_{crit} is ~ 29 for the present cases. We compare the simulation result with the two rules that distinguish two types of instabilities. First is the shape of deformation. The deformation of each blob extends over the whole blob as is observed in fig. 5.6, i.e., K-H vortices are observed in the simulation. In addition, the size of the blobs observed in the initial phase of fig. 5.6 is roughly estimated to be ~ 10 to 20 (in the regime of $t \sim 2000$). This is smaller than δ_{crit} . Thus, the instability that breaks blobs in this simulation is concluded to be the K-H instability.

We next study the evolution in longer time. An inverse cascade causes a mesoscale structure to appear at $t \sim 3000$. Radially elongated mesoscale convective cells are formed at $t \sim 5500$. Charge separation occurs inside convective cells, which drives the radial $\mathbf{E} \times \mathbf{B}$ drift. It was pointed out in ref. [67] that a single blob radially propagates while keeping its coherent structure; this is because, in the study of single-blob propagation, results are constrained by the initial blob size, which is given artificially as an initial condition. While the dynamics of the mesoscale structure in this simulation has some common features with that of single blobs, there is a distinct difference in this simulation. For the mesoscale structure in the present simulation, which is considered as a group of blobs, blobs are generated in the initial phase of the instability-driven fluctuations, and are subject to collapse to form an autonomous structure during their radial propagation. The characteristic size in the poloidal direction is determined by the self-organization mechanisms. The selection rule is discussed in a later section of this chapter.

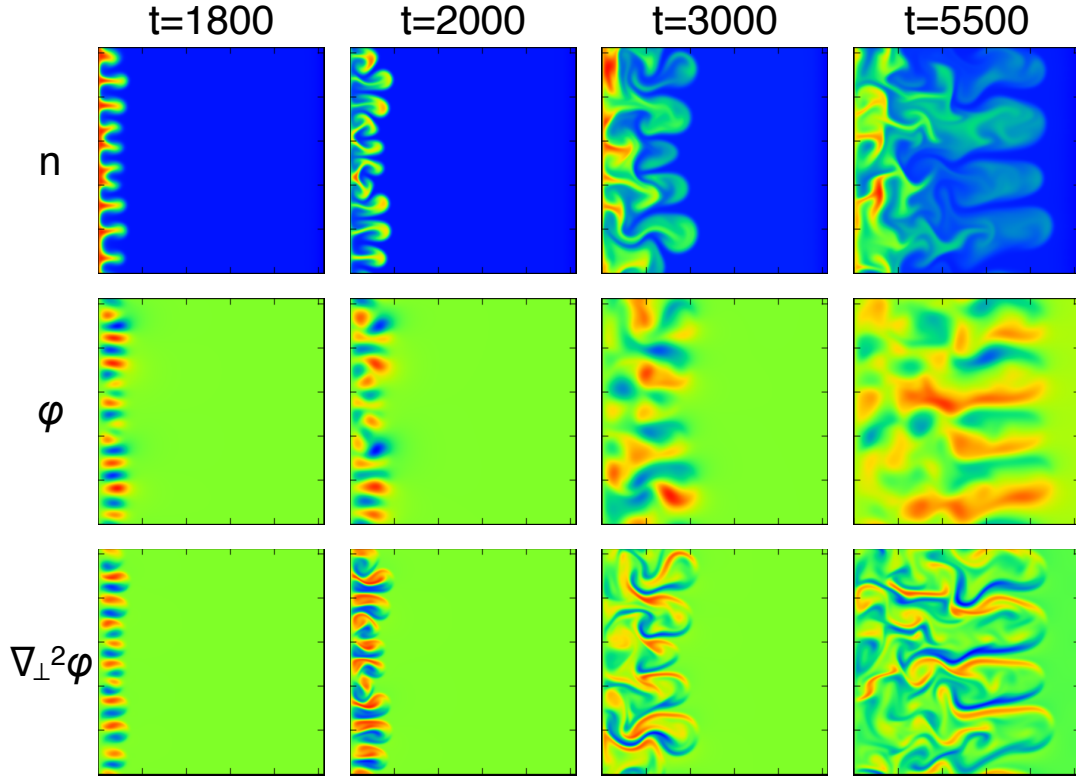


Fig. 5.6 From top to bottom: contour plots of particle density n , electrostatic potential φ and vorticity $\nabla_{\perp}^2 \varphi$. From left to right: time slices at $t = 1800, 2000, 3000$, and 5500 . The horizontal axis is the radial direction x and the vertical axis is the poloidal direction y . The length of each axis is $256[\rho_s]$ and time is in units of Ω_i^{-1} .

The temporal change of the characteristic scale length of the generated blobs can be clearly illustrated by observing the time evolution of Fourier components (in y -direction). Figure 5.7 shows the time evolution of the power spectra of the electrostatic energy defined by

$$E_{\varphi}(m) = \frac{\int_0^{L_x} |\nabla \varphi(x, m)|^2 dx}{L_x},$$

where $\varphi(x, m)$ is a coefficient of Fourier decomposition in y -direction, and m is the mode number. In the initial stage (phase A, $t \sim 1800$), blobs are formed at the separatrix and the poloidal mode number $m = 9$ is dominant. After the

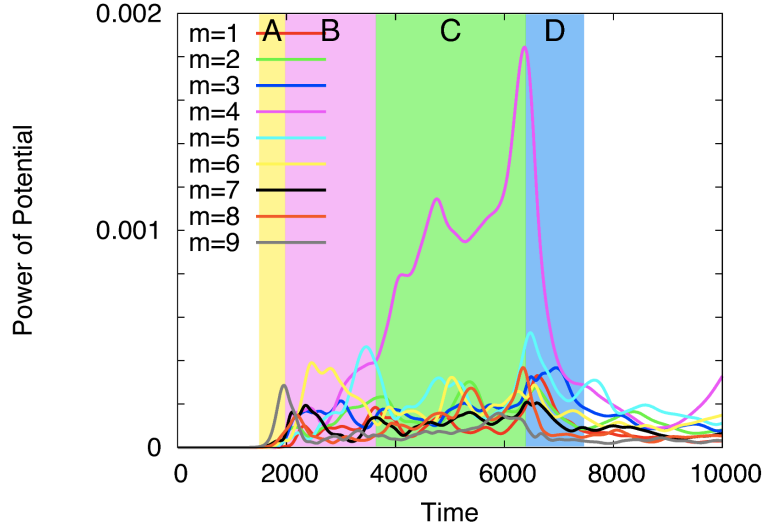


Fig. 5.7 Time evolution of power spectra of electrostatic potential energy. Each phase has a characteristic behavior. Phase A: blob formation, phase B: inverse cascading, phase C: mesoscale convective cell formation, phase D: break-up of mesoscale structure due to the boundary effect.

linear growth phase, the energy shifts to lower modes such as $m = 6$ and $m = 5$ (phase B, $t \sim 2000$). The inverse cascade produces a mesoscale structure with $m = 4$ (phase C, $t \sim 3000$). The mesoscale convective cells are broken into pieces by the boundary effect when they reach the wall and the $m = 4$ mode loses energy (phase D, $t \sim 5000$).

We here note that these results from simulation are not governed by the system size. The dependence of the mesoscale structure on the box size is also examined. Doubling the box size L_y , we see that eight mesoscale convective cells develop and the mode number $m = 8$ becomes prominent. That is, scale length that corresponds to $m = 4$ in original simulation is recovered.

The two-dimensional analysis of the spatial structure is performed. Figure 5.8 shows contour plots of the power spectra of the electrostatic potential $\varphi(x, m)$. The time slices are the same as those in fig. 5.6. The dominant modes are

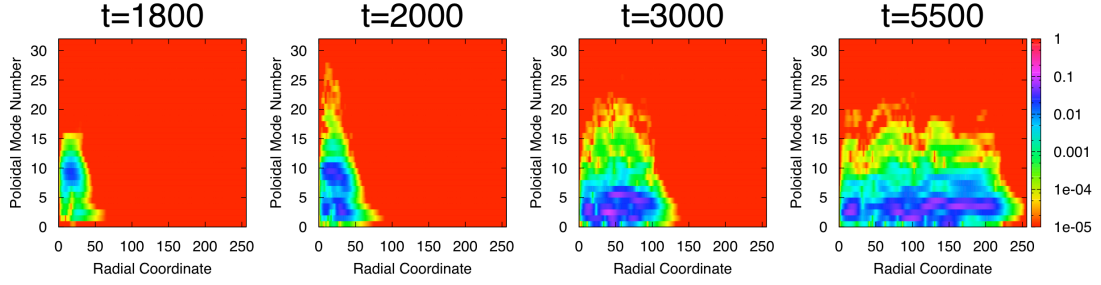


Fig. 5.8 Time slices of power spectra of electrostatic potential. The color scale is logarithmic. The horizontal axis is the radial coordinate x and the vertical axis is the Fourier mode m in the poloidal direction.

clearly observed for each time slice. In the initial phase, the linear growth of the interchange instability is dominant near the separatrix. Poloidal mode number is characterized by the linear growth rate. In the subsequent phase from $t = 2000$ to $t = 3000$, the formation of mesoscale structures is observed. Finally, these structures spread out radially and streamerlike formations appear at $t = 5500$. The scale length in the y -direction becomes much longer compared to the initial phase.

5.4 Turbulence-driven Radial Flux

We now evaluate the particle flux which is driven by turbulence. The turbulence-driven flux, Γ_t , is given by

$$\Gamma_t(x) = \frac{1}{L_y} \int_0^{L_y} n(x, y) (-\nabla_y \varphi(x, y)) dy, \quad (5.2)$$

while the flux which is caused by the molecular diffusion term $D\nabla_\perp^2 n$ in eq. (3.23) Γ_d is given by

$$\Gamma_d(x) = -D \frac{\partial}{\partial x} \frac{1}{L_y} \int_0^{L_y} n(x, y) dy. \quad (5.3)$$

The average over y is taken to study the mean density and mean flux. Figure 5.9 shows the time evolution of the poloidally averaged turbulence-driven and diffusive particle fluxes. The contour map in the top figure shows the time evolution of the Γ_t distribution in the radial direction. The light green stripes that extend from the bottom at $x = 0$ to the upper right direction indicate ballistic transport events. Intermittent transport also appears near the separatrix. The bottom three figures in fig. 5.9 show the time evolutions of Γ_t (turbulence-driven, red) and Γ_d (diffusive, blue) at the radial positions $x = 8$, $x = 64$, and $x = 128$. For turbulence-driven flux Γ_t near the separatrix (at $x = 8$), periodic bursts occur. However, the periodicity becomes less apparent at locations far from the separatrix. The turbulence-driven flux Γ_t is much larger than the diffusive flux Γ_d everywhere after the first avalanche takes place at $x = 8$. The SoL transport is dominated by the turbulence-driven transport.

We have investigated the effect of the parameter D (in eq. (3.23)), i.e., the collisional diffusion coefficient, on the turbulent transport associated with the mesoscale structures. We have confirmed that the SoL turbulent transport via blob formation is weakly dependent on D , except when D is exceptionally large. It would be advantageous to identify the order of magnitude of D , by comparing it with the quasi-linear estimation obtained by Kadomtsev [119],

$$D_{QL} \sim \gamma_{\max}/2k_y^2,$$

where $\gamma_{\max}$ is the maximum linear growth rate (for the instability occurring in the region of the source) and k_y is the wave number of the most unstable mode. Figure 5.10 shows the effect of the parameter D on the effective transport coefficient for blob transport, D_{eff} , and a comparison of D_{QL} with D . (The wave numbers are discrete because the system size is finite, thus, D_{QL} is discretized here.)

We introduce the effective transport coefficient D_{eff} in nonlinear simulations using the ratio of the turbulence-driven transport flux Γ_t to the diffusive flux

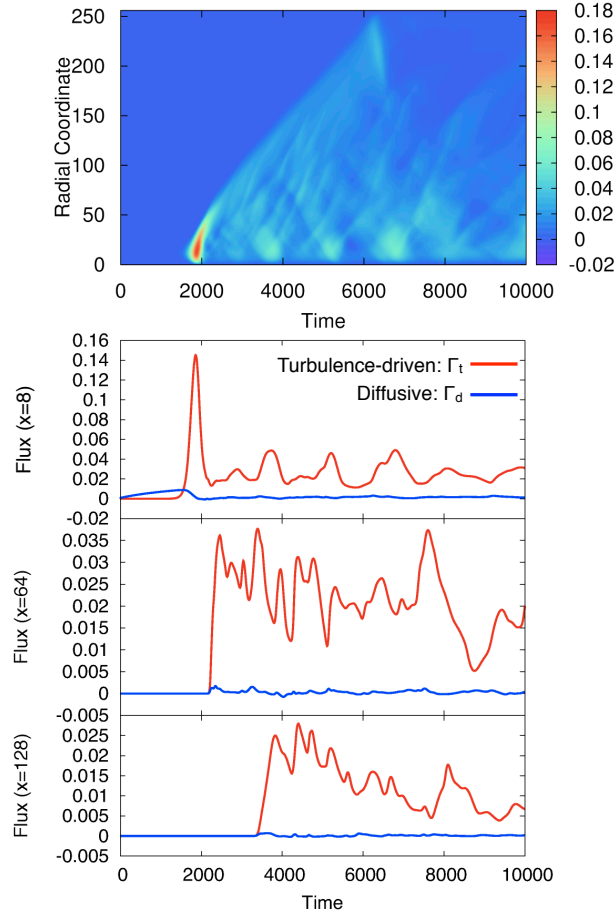


Fig. 5.9 Upper contour plot: time evolution of poloidally averaged turbulence-driven flux. The lower 3 plots show the time evolution of poloidally averaged turbulence-driven and diffusive fluxes at $x = 8$, $x = 64$, and $x = 128$.

Γ_d as

$$D_{\text{eff}} = (\Gamma_t / \Gamma_d) D. \quad (5.4)$$

The nonlinear simulation is performed for cases from $D = 2 \times 10^{-3}$ to 2×10^{-1} . The time average is taken from $t = 10000$ to 30000 to study the quasi-steady state. In the space, Γ_t and Γ_d are averaged over the y direction and from $x = 64$ to 192 to calculate D_{eff} . We here stress that D_{eff} (i.e., Γ_t) is evaluated by the value that is averaged over the finite length in x . Therefore, Γ_t and D_{eff} are no longer ‘local’ value. This is the reason for the term ‘effective diffusion’ is used.

This coefficient D_{eff} is *not* a local coefficient that connects local flux and local gradient.

In fig. 5.10, we plot D (denoted by D_{coll}), D_{QL} , and D_{eff} . (We see that the value of $D \sim 0.1$ is impractically large as it is on the order of the Bohm diffusion coefficient.) Upon choosing a realistic value for the collisional diffusion coefficient ($D < 0.1$), we observed that the turbulence-driven flux dominates the transport. Moreover, for small values of D , in which mesoscale fluctuations develop, the result is independent of D . The default parameter for D is assumed to be $D = 2 \times 10^{-2}$. Comparing the case for $D = 2 \times 10^{-2}$ with the case for the smallest value of D ($D = 2 \times 10^{-3}$), the characteristic size of the mesoscale structure is not different, the peak position of the k_y spectra is not affected, and the k_y spectra show the same tendency. (Of course, the mode number of perturbation in the initial stage ($t \lesssim 1000$) is different between these two cases reflecting the change in linear stability. However, the evolution in nonlinear phase is affected little by D in these simulations.) Thus, the result of turbulent transport in the reference case is considered to be almost independent of the choice of the parameter D . Under the situation that blob-like perturbation dominates the radial transport, the mean density profile deviates from the simple exponential form. Figure 5.11 shows the poloidally averaged density $n(x)$ at $t = 9820$. The dotted line is an exponential curve fitted to the data. We confirm that the deviation becomes larger for larger value of x . The density distribution that resembles the ‘2nd SoL’ is formed in the selfconsistent simulation.

5.5 Statistical Analysis of Turbulent Transport

The statistical method used for the analysis of the SoL plasma turbulence is discussed in this section. Three cases, the reference case ($\alpha = 3 \times 10^{-5}$, $\beta = 6 \times 10^{-4}$), low- α case ($\alpha = 3 \times 10^{-6}$, $\beta = 6 \times 10^{-4}$), and high- β case

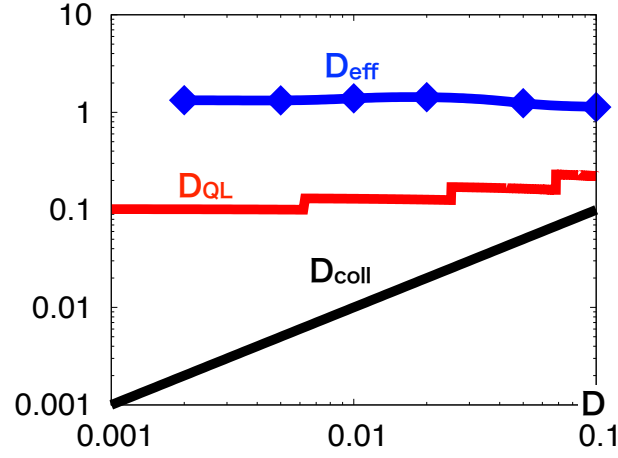


Fig. 5.10 D dependences of (collisional) diffusion coefficient D_{coll} , quasi-linear estimation of turbulent diffusion coefficient D_{QL} , and effective diffusion coefficient D_{eff} obtained by nonlinear simulation.

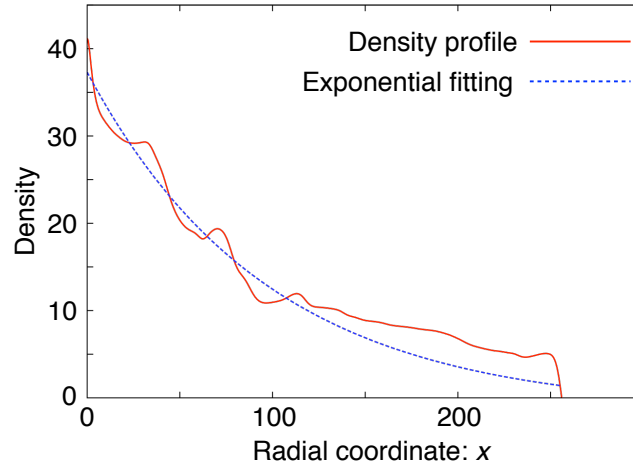


Fig. 5.11 Radial density profile of $m = 0$ mode at saturated phase and exponential fitting curve.

($\alpha = 3 \times 10^{-5}$, $\beta = 6 \times 10^{-3}$), are investigated. For the low- α case, modes with smaller poloidal mode numbers are unstable. The spatial scale length increases with the decrease of α . For the high- β case, the curvature is effective and the plasma tends to be polarized. As is shown by the linear growth rate

in figs. 5.2, 5.3, and 5.4, the interchange mode is strongly unstable and the fluctuation amplitude becomes larger. The perturbation includes modes with a large poloidal mode number. It was found that for all cases, the system reaches the quasi-steady state at $t > 10000$. Figure 5.12 shows power spectra of the electrostatic potential for three cases; the long-time evolutions up to $t = 30000$ are plotted. In each case, mesoscale modes are significantly strong in the quasi-steady state. The selection rule for the characteristic scale length of the mesoscale structure is discussed in Section 5.5.

Figure 5.13 shows the density and electrostatic potential contour plots for three cases at $t = 20000$. There is a tendency towards a finer structure when α increases and when β increases. In the low α case, it is seen that the mesoscale structure radially propagates while keeping its isolated structure. Charge separation occurs in the blob that produces $\mathbf{E} \times \mathbf{B}$ drift.

We also test the possible influence of the system size on the evolution of blob transport. The simulation result is compared with that of a simulation using the system size doubled. Figure 5.14 shows the k_y spectra of electrostatic potential in the quasi-steady state. The simulation parameters for the reference case are used. The spectra are averaged over time from $t = 15000$ to 30000. In fig. 5.14, for the case that the system size is doubled, the powers for unit modes are doubled to normalize them to those for a unit wave number. It is confirmed that the main conclusion is not dependent on the system size. In both cases, it is found that the mesoscale structure is important.

The phase relation between the turbulence-driven flux at two different radial positions allows us to estimate the spatiotemporal characteristics of the SoL radial turbulent transport. The autocorrelation function is defined by

$$C(\Delta x, \Delta t) = \frac{\langle \Gamma_t(x_0 + \Delta x, t_0 + \Delta t) \Gamma_t(x_0, t_0) \rangle}{\sqrt{\langle \Gamma_t^2(x_0 + \Delta x, t_0 + \Delta t) \rangle \langle \Gamma_t^2(x_0, t_0) \rangle}}, \quad (5.5)$$

where (x_0, t_0) indicates the reference position for the autocorrelation function;

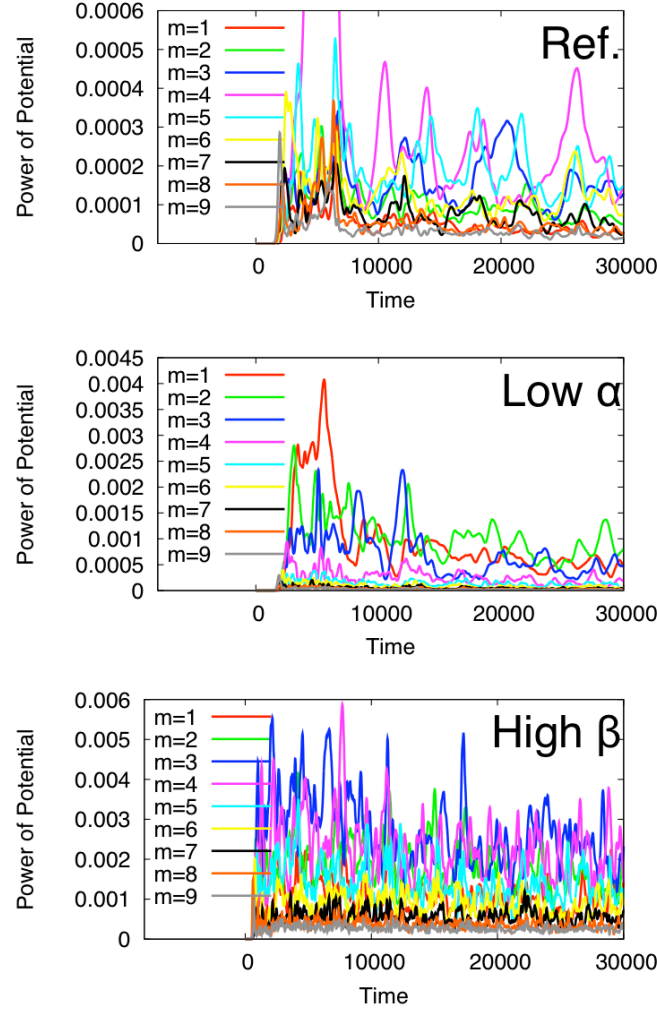


Fig. 5.12 Time evolution of power spectra of electrostatic energy for each Fourier mode. The results for three cases, reference, low- α , and high- β are shown.

$x_0 = 128$ and $t_0 = 10000$ are used. $\langle \rangle$ indicates the temporal average defined as

$$\langle f(t) \rangle = T^{-1} \int_t^{t+T} f(s) ds$$

with $T = 10000$. Figure 5.15 shows the contour plots of the autocorrelation function for three cases. It is seen that there is a strong spatial correlation between the separatrix at $x = 0$ and the wall at $x = 256$. The red stripes indicate strong ballistic transport with a scale length almost equivalent to the

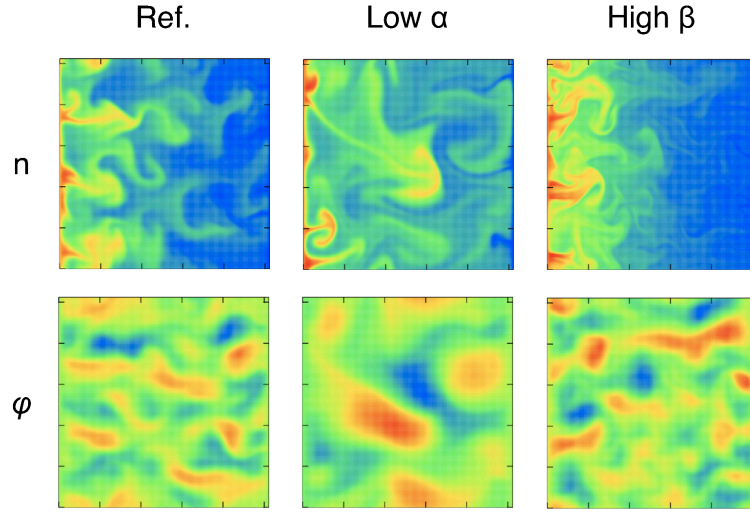


Fig. 5.13 Upper three contour plots: spatial profiles of particle density; lower three plots: electrostatic potential at $t = 20000$. Left to right: reference case, low- α case, and high- β case. The horizontal axis is the radial direction x and the vertical axis is the poloidal direction y . The length of each axis is $256[\rho_s]$ and time is in units of Ω_i^{-1} .

system size. The slope of the red stripes corresponds to the effective velocity of blobs that contributes SoL transport. A larger value of α with a larger value of β causes smaller reduction and larger enhancement of transport velocity, respectively. Parameters α and β are important factors in determining the size of the mesoscale structure and affect the transport velocity. Details of parameter dependence are discussed in the following section.

5.6 Scaling of Transport Velocity

In this section, we discuss how the transport velocity is determined on the basis of a dimensional argument. According to the study in Chapter 4, a single blob is deformed by the K-H instability when its size is small, and is modified by the interchange instability when it is large. The critical size that categorizes

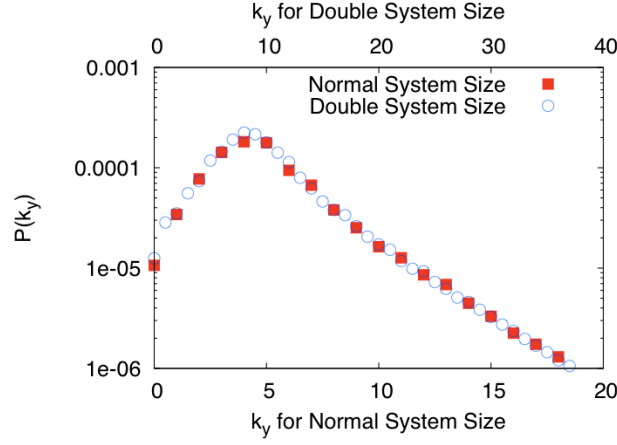


Fig. 5.14 Time-averaged k_y spectra of electrostatic potential perturbation $P(k_y)$ for simulations with normal system size and double system size. $P(k_y)$ for the double system size is doubled for normalization.

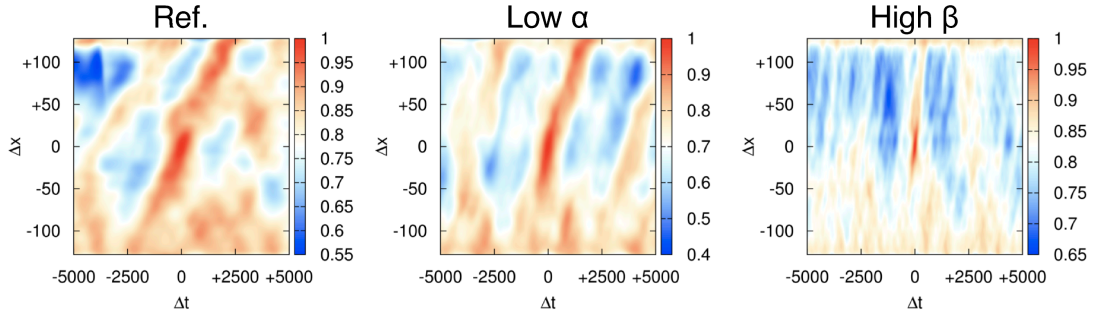


Fig. 5.15 Autocorrelation function of space and time for poloidally averaged turbulence-driven flux, $C(\Delta x, \Delta t)$, for reference, low- α , and high- β cases.

the type of instability [67] is given in eq. (4.1) by

$$\delta_{\text{crit}} \sim 2 \left(\frac{\beta}{\alpha^2} \right)^{1/5} = 2 \left(\frac{L_{\parallel}}{2\rho_s} \right)^{1/5} \left(\frac{L_{\parallel}}{R} \right)^{1/5}. \quad (5.6)$$

Above this size, inertia does not play an important role in the coherent blob dynamics; thus, the K-H instability is suppressed and the interchange instability becomes unstable. A blob of size $\sim \delta_{\text{crit}}$ can remain intact while radially

propagating.

The evolution of the size of a single blob was studied in fig. 4.12 of Subsection 4.3.3. The argument is reproduced here. The simulation is performed with the initial condition that the blob density distribution is prescribed by a two-dimensional Gaussian.

The effective poloidal size of the deformed blob is evaluated by the variance in the y -direction. The variance is defined by

$$\sigma^2(t) = \frac{\int (y - y_{\text{CM}}(t))^2 n_{\text{blob}}(y, t) dy}{\int n_{\text{blob}}(y, t) dy}, \quad (5.7)$$

where the radially averaged center-of-mass coordinate y_{CM} is defined as

$$y_{\text{CM}} = \frac{\int y \cdot n_{\text{blob}}(y, t) dy}{\int n_{\text{blob}}(y, t) dy}. \quad (5.8)$$

The radially averaged blob density profile $n_{\text{blob}}(y, t)$ is obtained by subtracting the constant background density profile from the poloidally averaged total density profile. The effective blob size is $\sqrt{2}\sigma$ in the y -direction. We recall the argument in Chapter 4. Figure 4.12 shows the time evolution of the effective blob size in the y -direction for various initial sizes. δ_{crit} is given by ~ 29 for the simulation parameters in the reference case ($\alpha = 3 \times 10^{-5}$, $\beta = 6 \times 10^{-4}$). The effective size approaches δ_{crit} asymptotically regardless of the initial blob size, where δ_b is smaller than δ_{crit} in the initial condition. We may adopt δ_{crit} as the blob size in the y -direction during the saturation phase. We here stress that the selfconsistent simulation shows that the fully-developed blobs have the size which is close to δ_{crit} . We therefore conclude that δ_{crit} might be applicable to estimate the effective blob size in the y -direction even for the multi blob situation shown in fig. 5.6.

These considerations give an analytic estimation for the velocity of generated blobs. The propagation velocity of a single blob can be estimated analytically.

This characteristic velocity in the SoL [67] is given by

$$V_x = \frac{2\beta/\alpha}{\delta_y^2}, \quad (5.9)$$

where V_x is the blob radial propagation velocity. Substituting $\delta_y \sim \delta_{\text{crit}}$ into eq. (5.9), we obtain the blob velocity as

$$V_x \sim 0.5\beta^{3/5}\alpha^{-1/5}. \quad (5.10)$$

In a dimensional form, it is rewritten as

$$V_x \propto \left(\frac{\rho_s}{R}\right)^{2/5} \left(\frac{L_{\parallel}}{R}\right)^{1/5} c_s. \quad (5.11)$$

Equation (5.11) indicates that how the magnetic confinement works in SoL [108]. It vanishes in the limit of $\rho_s/R \rightarrow 0$, but the speed of reduction is very slow. For instance, the diamagnetic drift velocity in the core plasma is given as $V_{\text{dia}} \propto (\rho_s/R) c_s$. It vanishes linearly with respect to ρ_s/R as $\rho_s/R \rightarrow 0$.

Here, we discuss the condition under which the formula of velocity, $V_x \sim (\rho_s/R)^{2/5} (L_{\parallel}/R)^{1/5} c_s$, is valid. To derive the 2D model equation (eqs. (3.22) and (3.23)), we integrated the 3D model along the magnetic field line, assuming that physical quantities except the current are uniform in the parallel direction. This approximation is valid as long as the loss of ions along the field line remains small. The parallel transport imposes the upper bound of the blob life time. The blob can propagate radially, during its life time, that is defined by parallel loss,

$$\tau_{\text{life}} \sim \frac{L_{\parallel}}{c_s}. \quad (5.12)$$

The radial distance that the blob can propagate during its life time is given by

the radial velocity and the life time as

$$\Delta_x = V_x \tau_{\text{life}} = \frac{L_{\parallel}}{c_s} \left(\frac{\rho_s}{R} \right)^{2/5} \left(\frac{L_{\parallel}}{R} \right)^{1/5} c_s \quad (5.13)$$

$$= \frac{L_{\parallel}}{a} \left(\frac{\rho_s}{R} \right)^{2/5} \left(\frac{L_{\parallel}}{R} \right)^{1/5} a, \quad (5.14)$$

where a indicates the plasma minor radius and the other parameters have same meanings as introduced in Chapter 3. The ratio of the radial distance to the minor radius, below which the 2D blob model is valid, is given by the following equation:

$$\frac{\Delta_x}{a} = \left(\frac{R}{a} q^{6/5} \right) \left(\frac{\rho_s}{R} \right)^{2/5}, \quad (5.15)$$

where the estimation $L_{\parallel} \sim qR$ is used. It should be noted that this limit of Δ/a is proportional to $(\rho_s/R)^{2/5}$. Although ρ_s/R is a smallness parameter, the fractional power $(\rho_s/R)^{2/5}$ remains finite. Using the typical ITER parameters, $\Delta_x/a \sim 0.25$ and $\Delta_x \sim 50\text{cm}$ are evaluated. One sees that in the radial transport in the SoL, blobs can play an important role. This result indicates that the 2D model can be applied to present-day experiments as well as to fusion experiments.

The blob velocity, calculated by this formula (5.10) is compared with the ballistic propagation velocity of clumps, V_{eff} , which can be evaluated from the slope of the autocorrelation function in fig. 5.15. The value estimated from eq. (5.10) is compared with V_{eff} for the reference, low- α , and high- β cases in fig. 5.16. It was found that V_x and V_{eff} are in good agreement for all cases.

5.7 Scaling Relation of Turbulence-driven Flux

We next study the dependence of effective diffusivity D_{eff} . We performed parameter survey for several cases for different values of α and β . (See fig. 5.17, which shows α -dependence while β is fixed, and 5.18 that illustrates

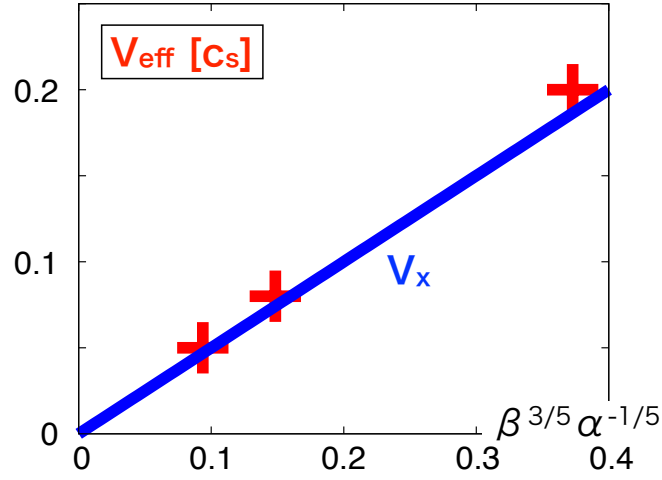


Fig. 5.16 Comparison between V_x and V_{eff} estimated from nonlinear simulation. Unit of velocity c_s indicates ion sound speed.

β -dependence with fixed α .) Dependencies are studied for the region

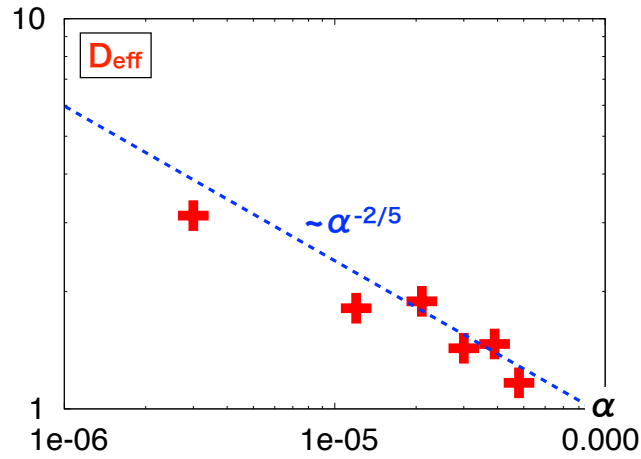


Fig. 5.17 α -dependence of effective diffusion coefficient, D_{eff} . It has a tendency proportional to $\alpha^{-2/5}$. $\beta = 6 \times 10^{-4}$ and $D = \mu = 2 \times 10^{-2}$ are fixed.

of small value of D_{eff} ($D_{\text{eff}} < 2$), while the D_{eff} is much larger than the quasi-linear estimate γ_L/k_y^2 , as was explained in the Section 5.2 of linear stability, the parameter α reduces fluctuations and transport. By increasing

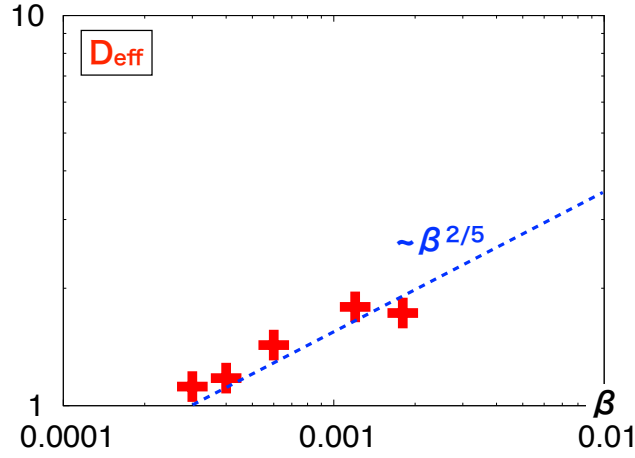


Fig. 5.18 β -dependence of effective diffusion coefficient, D_{eff} . It has a tendency proportional to $\beta^{2/5}$. $\alpha = 3 \times 10^{-5}$ and $D = \mu = 2 \times 10^{-2}$ are fixed.

α , the turbulence-driven flux, or D_{eff} , decreases as is shown in fig. 5.17. In contrast, the curvature (represented by the parameter β) is the driving source of turbulence. D_{eff} is an increasing function of β . We have the result

$$D \sim \left(\frac{\beta}{\alpha} \right)^{2/5}. \quad (5.16)$$

That is, in a dimensional form, we have the following Bohm-like dependence:

$$D_{\text{eff}} \sim \left(\frac{L_{\parallel}}{R} \right)^{2/5} \rho_s c_s \sim \left(\frac{L_{\parallel}}{R} \right)^{2/5} \frac{T}{eB}. \quad (5.17)$$

We have stressed in Section 5.4 that D_{eff} is obtained by averaging in the region $L_x/4 < x < 3L_x/4$, and that the intermittent pulses of turbulent transport have long radial correlation (fig. 5.15). Thus D_{eff} is not a local diffusion coefficient, but is a measure of non-local transport quantity. As has been pointed out by many articles, the long range fluctuation can induce Bohm-like transport. This simulation demonstrates that the selfconsistently evolved blobs in flux-driven nonlinear simulation reproduce the Bohm-like transport. This transport is

different from local diffusive transport, so that the profile deviates from simple exponential decay as is illustrated in fig. 5.11.

The estimation equation (5.16) is given by the fitting of data of nonlinear simulations as an inductive approach. Reviewing the eqs. (5.10) and (5.16), we get the following relation:

$$D_{\text{eff}} \propto (\alpha\beta)^{-1/5} V_x. \quad (5.18)$$

By assuming that the effective transport in the SoL is dominated by the blob transport the similarity argument shows that D_{eff} has a dimension of the step size and velocity of blobs. The R.H.S. of eq. (5.18) includes following factors: blob velocity, correlation length in the radial direction, the blobs amplitude, and the passing frequency of blobs. The velocity is already separated in the eq. (5.18). Remaining factors may determine the dependence in R.H.S. of eq. (5.18), $(\alpha\beta)^{-1/5}$. We investigated space-time correlation of turbulence-driven flux for low- α ($\alpha = 3.0 \times 10^{-6}, \beta = 6.0 \times 10^{-4}$), high- α ($\alpha = 4.8 \times 10^{-5}, \beta = 6.0 \times 10^{-4}$), low- β ($\alpha = 3.0 \times 10^{-5}, \beta = 3.0 \times 10^{-4}$), and high- β ($\alpha = 3.0 \times 10^{-5}, \beta = 6.0 \times 10^{-3}$) cases. The dependence of radial correlation length on α or β is not distinctive. Thus, α and β dependencies on the D_{eff} through the radial correlation length might be weak. Next, we plot the time evolution of the poloidally averaged turbulence-driven flux at $x = L_x/4$. The α -dependence is shown in fig. 5.19, and the β -dependence is illustrated in fig. 5.20. It can be observed that the amplitude of the turbulence-driven flux decreases as the value of α increases, thus large value of α makes the D_{eff} small. This is consistent with eq. 5.18. Although the origin of β -dependence of the coefficient $(\alpha\beta)^{-1/5}$ in eq. 5.18 has not been completely identified, one observes that frequently appearing blobs enhance the time averaged value of D_{eff} in higher- β case (fig. 5.20). It is suggested that the origin of the factor $(\alpha\beta)^{-1/5}$ might be attributed to the amplitude of the flux and the frequency qualitatively.

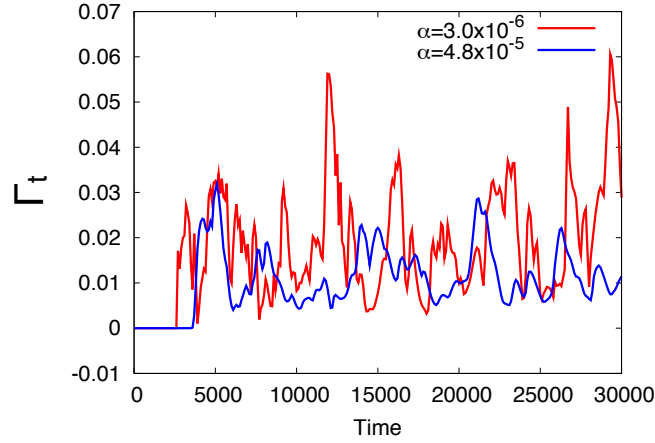


Fig. 5.19 Time evolution of poloidally averaged turbulence-driven flux at $x = L_x/4$. Low- α case and high- α are plotted.

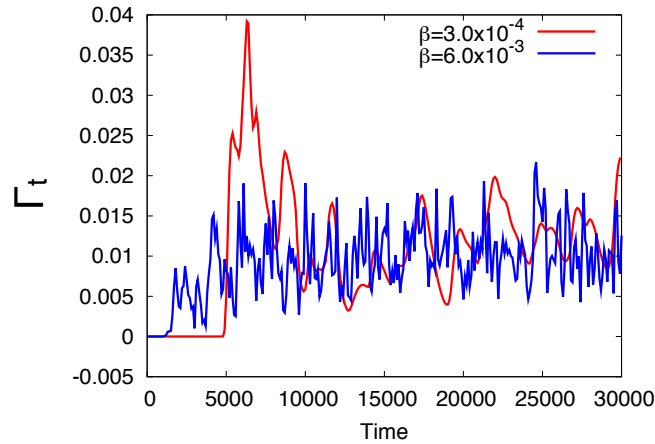


Fig. 5.20 Time evolution of poloidally averaged turbulence-driven flux at $x = L_x/4$. Low- β and high- β case are plotted.

A few comment is made on the formula (5.18) here. The coefficient $(L_{\parallel}/R)^{2/5}$, reflects impacts of geometrical parameters. First, the dependence of $R^{-2/5}$ comes from the fact that the bad curvature of the main magnetic field is the origin of driving force. The dependence $L_{\parallel}^{2/5}$, coming from $\alpha^{-2/5}$ dependence in collisionless form, indicates the fact that the conductivity is finite at

the sheath. When geometrical configuration is common, the factor $(L_{\parallel}/R)^{2/5}$ does *not* depend on the system size. The factor $(L_{\parallel}/R)^{2/5}$ is influenced by geometrical factor. For instance, it becomes smaller if the safety factor q is smaller. Equation (5.18) implies the dependence like $D_{\text{eff}} \propto q^{2/5} T/eB$.

5.8 Summary of Chapter 5

The turbulent transport in SoL is investigated using the flux-driven nonlinear simulation that employs the 2D Hasegawa-Wakatani-type resistive interchange model. Linear stability is analyzed, and influences of coefficients α , β , D , and μ are studied. Spontaneous generation of blobs is demonstrated. During the turbulence evolution, radially-elongated mesoscale convective cells are formed autonomously and they dominate the SoL radial transport. Scaling law of velocity is obtained as $V_x \propto (\rho_s/R)^{2/5} (L_{\parallel}/R)^{1/5} c_s$. The effective transport coefficient is found to obey the Bohm-like dependence, $D_{\text{eff}} \sim (L_{\parallel}/R)^{2/5} T/eB$. The origin of the parameter dependence of the effective transport coefficient is investigated. The transport velocity, amplitude of the turbulence, and the frequency of the turbulence are controlling factors.

Chapter 6

Extension to Edge Turbulence: on Stiffness of Transport

6.1 Introduction

In the preceding chapters, we have discussed the transport in scrape-off-layers, where blob-like structures play a dominant role. It is shown that the perturbations, which have long correlation length in the radial direction, enhance the transport so that the Bohm-like effective transport coefficient is deduced. This result stimulates the study of the edge plasma, where the long-range fluctuations can have influences on the transport. We here report an extension of the research to the study of edge plasmas. (edge plasma: plasma within the last closed magnetic surface and close to it) The evolution of the plasma, which is specified by the source flux at the boundary, is studied by the flux-driven nonlinear simulations in Chapter 5. As an extension of the studies in Chapter 5, we performed flux-driven nonlinear simulation of turbulence in the edge. Of course, the studies of the edge transport have its own history and a large number of literatures. This chapter illustrates an ongoing effort to combine our view of global transport in SoL with the global transport problem in the edge [120].

6.2 Illustration of the Problem

Recent experiments at JET have shown that driving a large-scale rotation of the plasma improves confinement. In particular, the so-called stiffness level of the ion temperature profile is found to decrease with plasma rotation [121]. (The ‘stiffness’, χ_s , is introduced phenomenologically by the relation $Q \propto \chi_s \nabla p (\nabla p - \nabla p_{\text{crit}})$, where Q indicates the energy flux, ∇p the pressure gradient, and ∇p_{crit} the threshold of the pressure gradient.) The level of stiffness characterizes the fact that above a critical value of the ion temperature inverse gradient length ($R/L_{T_i} = R|\nabla T_i|/T_i$, where R indicates the tokamak major radius) the ion heat flux increases strongly with R/L_{T_i} . This behavior is attributed to turbulent transport. The observations have been made in the core [121], but the ingredients to investigate this problem are included in any turbulence and rotation simulation. As one step in understanding the problem of ‘stiffness’ in transport relation, a fluid edge turbulence code is applied here.

In this work we report on 3D turbulent simulations for the edge plasmas using a flux-driven-configuration (i.e., with selfconsistent profile evolution under the prescribed heat source at an edge), revealing the behavior of the mean pressure gradient as a function of the (imposed) total energy flux. Different scenarios with respect to plasma rotation (spontaneous rotation, artificially suppressed rotation or controlled rotation) are studied [120].

6.3 Model

6.3.1 Model equation

The two-field reduced-MHD equations are solved to reproduce the plasma edge turbulence in the tokamak geometry [122]. The following equations are em-

ployed for the normalized electrostatic potential ϕ and pressure p .

$$\frac{\partial}{\partial t} \nabla_{\perp}^2 \phi + [\phi, \nabla_{\perp}^2 \phi] = -\nabla_{\parallel} \phi - \mathbf{G}p + \mu \nabla_{\perp}^4 \phi, \quad (6.1)$$

$$\begin{aligned} \frac{\partial}{\partial t} \nabla_{\perp}^2 \phi_{00} + [\phi, \nabla_{\perp}^2 \phi_{00}] &= -\nabla_{\parallel} \phi_{00} - \mathbf{G}p + \mu \nabla_{\perp}^4 \phi_{00} \\ &+ \hat{\mu}_{\text{control}} \nabla_{\perp}^2 (\phi_{\text{imp}} - \phi_{00}), \end{aligned} \quad (6.2)$$

$$\frac{\partial p}{\partial t} + [\phi, p] = \epsilon_c \mathbf{G}\phi + \chi_{\parallel} \nabla_{\parallel}^2 p + \chi_{\perp} \nabla_{\perp}^2 p + S(r), \quad (6.3)$$

The notations and symbols are explained as follows. ϕ_{00} indicates the $(m, n) = (0, 0)$ Fourier component of the electrostatic potential ϕ , where m is the poloidal and n is the toroidal mode numbers of Fourier decomposition respectively. ϕ_{imp} indicates the artificially imposed $\mathbf{E} \times \mathbf{B}$ velocity shear to ϕ_{00} component, which is discussed later. In this MHD model, the diamagnetic velocity is neglected in comparison with the $\mathbf{E} \times \mathbf{B}$ velocity, and the parallel current is evaluated by using a simplified electrostatic Ohm's law, $\eta_{\parallel 0} j_{\parallel} = -\nabla_{\parallel} \phi$, where $\eta_{\parallel 0}$ is a reference value of the parallel resistivity, $j_{\parallel}$ is the parallel current, and $\nabla_{\parallel}$ is the gradient parallel to the field line. Magnetic flux surfaces are modeled in toroidal coordinates (r, θ, ζ) that correspond to the tokamak minor radius, poloidal angle, and toroidal angle, respectively. The system is normalized using the interchange time τ_{int} , ballooning length ξ_{bal} , and magnetic shear length L_s :

$$\begin{aligned} \frac{r - r_0}{\xi_{\text{bal}}} &\rightarrow x, & \frac{r_0 \theta}{\xi_{\text{bal}}} &\rightarrow y, & \frac{R_0 \zeta}{L_s} &\rightarrow z, \\ \frac{t}{\tau_{\text{int}}} &\rightarrow t, & \frac{\tau_{\text{int}} \phi}{B_0 \xi_{\text{bal}}^2} &\rightarrow \phi, & \frac{L_p p}{\xi_{\text{bal}} p_0} &\rightarrow p, \end{aligned}$$

where

$$\begin{aligned} \tau_{\text{int}}^2 &= R_0 L_p / (2c_{s0}^2) && : \text{resistive interchange time,} \\ \xi_{\text{bal}}^2 &= m_i n_0 \eta_{\parallel 0} L_s^2 / (\tau_{\text{int}} B_0^2) && : \text{resistive ballooning length.} \end{aligned}$$

Here, all coefficients in the normalizations are evaluated with reference values in the vicinity of a flux surface r_0 . The domain of simulation is explained

later. In eqs. (6.1)-(6.3) $[f, g] = (\partial f / \partial x)(\partial g / \partial y) - (\partial f / \partial y)(\partial g / \partial x)$ is the Poisson bracket. $\nabla_{\perp}^2 = \partial^2 / \partial x^2 + \partial^2 / \partial y^2$ and $\nabla_{\parallel} = \partial / \partial z - (k_z / q)(\partial / \partial y)$ are the gradients perpendicular and parallel to the field line. The curvature operator $\mathbf{G} = \sin(k_y y)(\partial / \partial x) + \cos(k_y y)(\partial / \partial y)$ is introduced. μ indicates the viscosity, $\chi_{\parallel}$ and $\chi_{\perp}$ are parallel and perpendicular heat diffusivities. ϵ_c is the parameter that is related to the ratio between the pressure gradient length and the major radius R_0 . Time is normalized by the resistive interchange time $\tau_{\text{int}} = (R_0 L_p / 2)^{1/2} / c_s$, where c_s is the sound speed, and L_p is the pressure gradient length. Perpendicular length is normalized by the resistive ballooning length $\xi_{\text{bal}} = (\rho \eta_{\parallel} / \tau_{\text{int}})^{1/2} L_s / B_0$, where ρ is the mass density, $\eta_{\parallel}$ is the parallel resistivity, and parallel length is normalized to the magnetic shear length L_s . The parameter $\hat{\mu}_{\text{control}}$ (the last term in eq. (6.2)) is introduced in order to control the mean velocity. The explanation is given later.

6.3.2 Geometry and simulation conditions

The model equation is solved numerically by the finite difference method in the radial direction and by a Fourier expansion in the poloidal and toroidal directions. In the (m, n) space, all modes are simulated except for the $(n = 0, m > 0)$ modes to keep the neoclassical terms consistent (especially for $(1, 0)$ mode) [123].

By assuming that the safety factor $q(r)$ is monotonically increasing, the simulation domain covers a region between $q = 2.5$ and $q = 3.5$. The regions $q < 2.5$ and $q > 3.5$ are used for a dumping buffer. In the dumping buffer region, the dissipation rate is set to be extraordinarily large. Gaussian type pressure source is introduced in the inner buffer region, i.e., $q \lesssim 2.5$. We vary the amplitude of the source such that the integral, i.e. the total energy flux varies from 1 to 32 in normalized unit (see fig. 6.1). The simulations are performed for the following three scenarios concerning the poloidal rotation:

spontaneous rotation case, artificially suppressed case, and controlled rotation case. In the case of spontaneous rotation, the coefficient $\hat{\mu}_{\text{control}}$ in eq. (6.2) is set to be zero. In the suppressed case and in the artificially imposed case, we take the limit of $\hat{\mu}_{\text{control}} \rightarrow \infty$, i.e., the condition $\phi_{00} = \phi_{\text{imp}}$ is employed, ϕ_{imp} indicates the imposed mean potential profile.

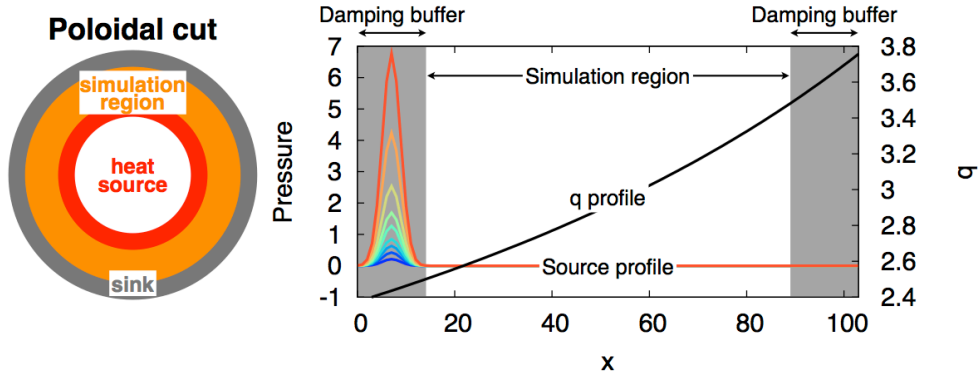


Fig. 6.1 Radial profiles of pressure source and safety factor q . Region of $2.5 \leq q \leq 3.5$ is main simulation region and outside of this region is used as a damping buffer. Pressure source is given in this damping buffer.

6.4 Simulation Results

6.4.1 Evolutions of pressure gradient and flow (case of spontaneous rotation)

Nonlinear simulations are performed for the case of spontaneous rotation. Evolutions are observed until a quasi-steady state is reached, and the temporal and spatial (poloidal and toroidal) averages are taken in the quasi-steady state. We use the word ‘quasi-steady’, because this state is a dynamical steady state in which intermittent enhancement of flux is included. After this averages, we use the word ‘stationary’. The pressure profile in a stationary state is illustrated in fig. 6.2. Associated with the global pressure gradient, the mean poloidal flow

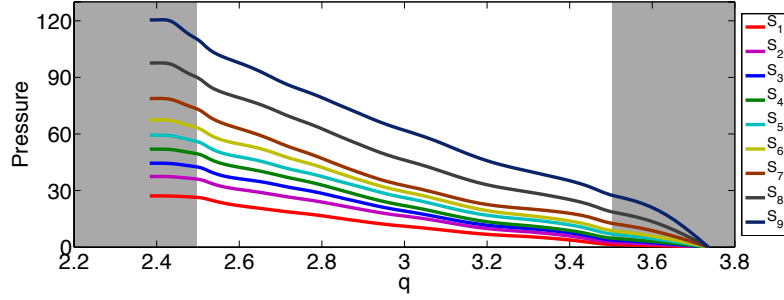


Fig. 6.2 Time averaged profiles of pressure on q in spontaneous rotation case. Averaged temporal length is $\tau = 1000$ in steady state. S_n indicate pressure source variation. From S_1 to S_9 , source integrals are 1, 2, 3, 4, 6, 8, 12, 20, 32. Main simulation region is $2.5 \leq q \leq 3.5$

is generated. Figure 6.3 shows the radial profile of the $\mathbf{E} \times \mathbf{B}$ velocity shear in the selfconsistent case with spontaneous rotation. It is evident that the shear amplitude is increasing with the total energy flux.

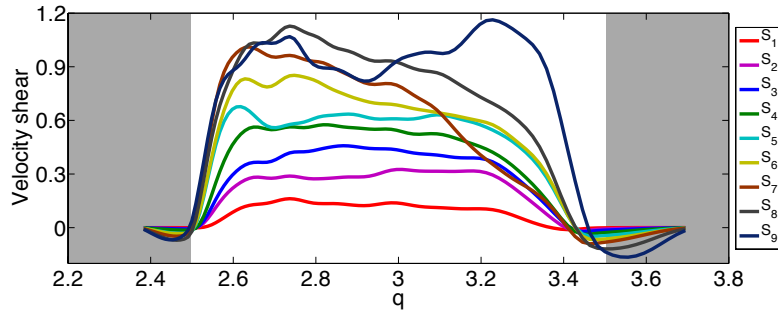


Fig. 6.3 Time averaged profiles of poloidal velocity shear on q in spontaneous rotation case.

Temporal evolution of global parameters is shown in fig. 6.4 in order to illustrate that the plasma reaches the quasi-steady state. The mean pressure and flux (which are averaged over magnetic surface at $x = 52$, i.e., the center of the simulation domain) are drawn. The simulation is flux-driven and the pressure source is localized in the region $x < 14$. The flux at the center ($x = 52$) shows intermittent oscillations. In figs. 6.2 and 6.3, time averages of pressure profile and the poloidal velocity shear ($4000 \leq t \leq 5000$) are shown.

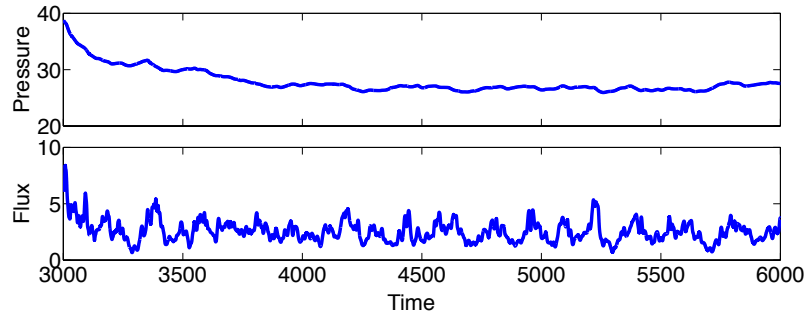


Fig. 6.4 Temporal evolution of pressure and electrostatic flux. Sampling radial position is the center. Spontaneous rotation case. Pressure source integral is 4.

A snap shot of two-dimensional structure of fluctuations are shown in fig. 6.5. Note that long-range stripes are observed in the contour of the perturbed pressure and in that of flux.

6.4.2 Fitting to gradient-flux relation

In figure 6.6 the mean pressure gradient is plotted versus the total imposed energy flux. The spontaneous case is compared to the case where the rotation is artificially suppressed (i.e., $\phi_{00} = 0$). It is clearly shown that the rotation

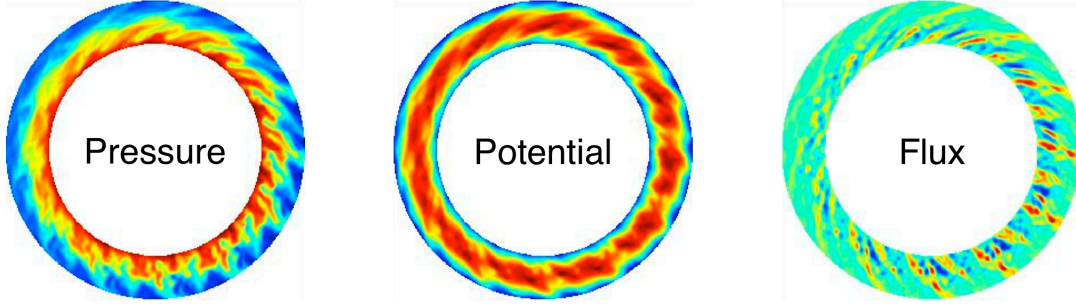


Fig. 6.5 Poloidal cuts of pressure, electrostatic potential, and electrostatic flux. Snap shot is taken at $t = 5000$ in spontaneous rotation case. Pressure source integral is 4.

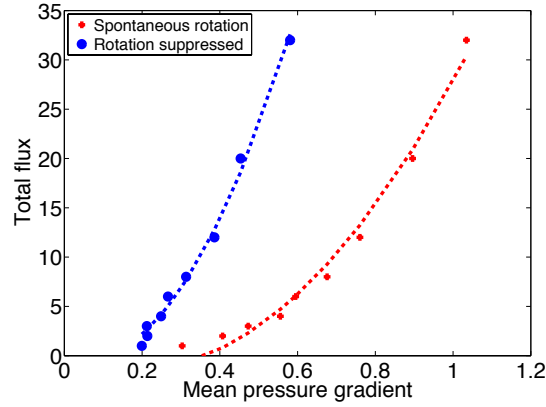


Fig. 6.6 Mean pressure gradient versus total imposed flux with fitting curves for no-rotation case and spontaneous rotation case.

keeps the pressure gradient steep for constant heat flux. For the fitting curve, we use the gyro-Bohm type relation between the total energy flux and the pressure gradient, parameterized by the stiffness level and the critical gradient (eq. (6.4)) [124]. That is,

$$Q_{\text{tot}} = Q^{\text{res}} + \chi_{\text{g-B}} q^{1.5} \frac{\tau_{\text{int}}}{\xi_{\text{bal}}^2} \chi_s \frac{dp}{dr} \left(\frac{dp}{dr} - \frac{dp}{dr} \Big|_{\text{crit}} \right), \quad (6.4)$$

Table. 6.1 Values of stiffness χ_s and the critical gradient $dp/dr|_{\text{crit}}$ for the cases where the rotation is spontaneous and where the rotation is artificially suppressed.

	χ_s	$dp/dr _{\text{crit}}$
spontaneous rotation	1.3614	0.38446
suppressed rotation	3.6942	0.11537

where Q_{tot} is the total imposed source flux, Q^{res} the residual flux (including the neoclassical flux), $\chi_{\text{g-B}} = T_e \rho_i / (eBR)$ the gyro-Bohm coefficient, q the safety factor, χ_s the stiffness level, and $dp/dr|_{\text{crit}}$ the threshold of the pressure gradient. The fitting gives χ_s and $dp/dr|_{\text{crit}}$ for each curve. The fitting is done in the statistically stationary state (not the relaxation oscillation or simple diffusion). Solid lines in fig. 6.6 are drawn according to eq. (6.4) with the fitted constants χ_s and $dp/dr|_{\text{crit}}$. If one does the fitting on fig. 6.6, one obtains values of stiffness χ_s and the critical gradient $dp/dr|_{\text{crit}}$ for the cases where the rotation is spontaneous and where the rotation is artificially suppressed. The values are shown in table 6.1.

6.4.3 Role of rotation in gradient-flux relation

Figure 6.6 clearly shows that the spontaneous rotation reduces the transport which is expressed by the slope of curves. The phenomenological fitting (tbl. 6.1) may suggest that χ_s is reduced by spontaneous rotation. The mechanism on this reduction is studied by artificially controlling the velocity shear profile. We impose different rotation velocities with broad velocity shears, and simulations are performed for variety of amplitudes (fig. 6.7). In each case of specified flow (e.g., the case of V_5 in fig. 6.7), the magnitude of source is varied from S_1 to S_9 . As a whole, including the case of zero velocity, $11 \times 9 = 99$ sim-

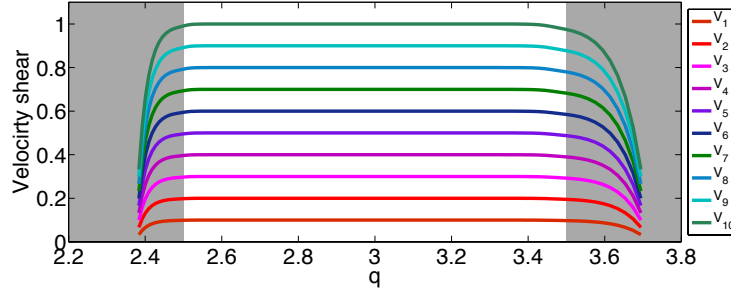


Fig. 6.7 Imposed poloidal velocity shear profile on q in imposed rotation case.

ulations are performed as controlled cases. In each simulation, mean pressure gradient $\overline{\nabla p}$ is obtained. The gradient-flux relation $(\overline{\nabla p}, S)$ is plotted in fig. 6.8. In fig. 6.8a, the result of the controlled simulations is summarized. It should

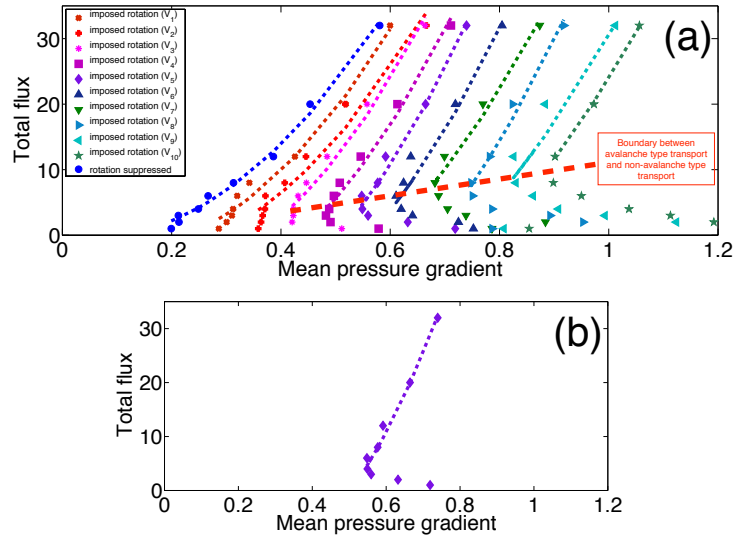


Fig. 6.8 Mean pressure gradient versus total imposed flux with fitting curves

be noticed that the gradient-flux relation is not monotonic for fixed imposed velocity shear. For instance, we illustrate one line (case of V_5) in fig. 6.8b. In the region where the given flux is very small, the transport is strongly intermittent in the time evolution. We here focus on the case of strong heat source, where flux is an increasing function of gradient. In this region of interest, the data $(\overline{\nabla p}, S)$ are highlighted by dotted lines in fig. 6.8a.

The fitting parameters χ_s and $dp/dr|_{\text{crit}}$ are also deduced for the results of controlled cases. The fitting is shown by dotted lines in fig. 6.8a. Now we see how the imposed flow influences on χ_s and $dp/dr|_{\text{crit}}$. Figure 6.9 shows the stiffness level χ_s versus the velocity shear amplitude and fig. 6.10 shows the critical gradient $dp/dr|_{\text{crit}}$ versus the velocity shear amplitude, both in the controlled case. We find that the critical gradient is almost linearly

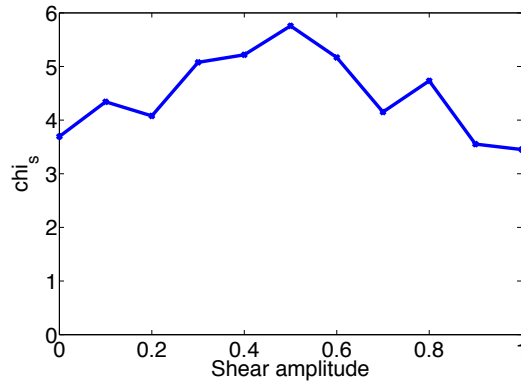


Fig. 6.9 Stiffness level χ_s obtained from fitting of data in fig. 6.8 with eq. (6.4) versus imposed shear amplitude.

proportional to the strength of the velocity shear, but the influence on the stiffness level is weak.

The relevance of this controlled simulation in understanding the selfconsistent cases is confirmed for comparing figs. 6.6 and 6.8a. Data of spontaneous

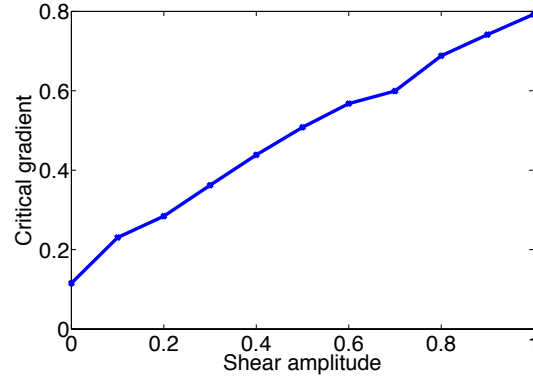


Fig. 6.10 Critical gradient of the pressure $dp/dr|_{\text{crit}}$ obtained from fig. 6.8 versus imposed shear amplitude.

simulations are superimposed on result of controlled simulations in fig. 6.11. The spontaneous cases follow the results of controlled cases by increasing im-

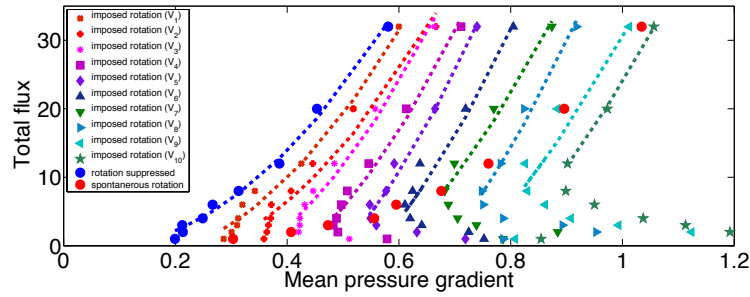


Fig. 6.11 The case of spontaneous rotation (fig. 6.6) is overlaid on the case of controlled rotation (fig. 6.8)

posed velocity as the flux S increased.

6.5 Summary of Chapter 6

We here studied the global simulation of edge plasma by flux-driven nonlinear simulations. The gradient-flux relation is fitted to a model with parameters of stiffness and critical gradient, eq. (6.4). The slope of the grad- p -v.s.-flux-curve is clearly lower in the presence of spontaneous rotation, compared to the case where rotation is artificially suppressed. The origins of this difference are studied by imposing specified velocity shear. The slopes of Q_{tot} v.s. $\overline{\nabla p}$ are similar for cases with different amplitudes of controlled rotation. That is, the stiffness does not change significantly, while the critical gradient changes. Therefore, the change of the slope observed in the case of the spontaneous rotation has to be attributed to the fact that the velocity shear amplitude itself is increasing with the pressure source. The increase of the velocity shear is considered to cause the change of the critical gradient in the case of the spontaneous rotation, while it does not cause the change of the stiffness. A tentative conclusion is given that the origin of the effect on the gradient-flux curve given by the spontaneous rotation is found to be an increment of critical gradient.

Chapter 7

Summary

In this thesis we first made an overview of the transport problems in the Scrape-off Layer region of tokamak plasmas. It is found that, in SoL plasmas, the transport across the magnetic field lines is often dominated by the radial motion of filamentary structures which are elongated along the field lines but are localized across the magnetic field. These filaments contain excess particles and heat compared to background plasma, and are thus referred to as plasma blobs.

The dynamics of blobs and the associated radial transport are the subject of this thesis. Particularly, the blob velocity and the evolution of size during radial fast propagation are important issues. Characteristics of the transport are also crucial. For instance, the velocity of blob, transport coefficient, radial profile associated with the transport, and so on must be clarified.

We here summarize the findings in this thesis as follows.

In Chapter 3, the blob model and the model equations are discussed. The model equation is characterized by dimensionless parameters α and β , $\alpha = \rho_s/L_{\parallel}$ and $\beta = \rho_s/R$, where α represents the impact of finite conductivity at the sheath and β denotes the bad magnetic curvature. In addition, particle diffusivity and viscosity by collisional process, D and μ , are introduced. The model system is controlled by these four parameters. We choose the ‘Conserved

model', and the comparison with others (Linearized and Simplified models) is explained.

In Chapter 4, we investigated the motion of a single blob by solving eqs. (3.22) and (3.23), i.e., the Hasegawa-Wakatani-type model equation for blobs. It was found that the density nonlinearity in the vorticity equation has a dominant effect on the blob propagation and that the nonlinearity in the continuity equation causes the asymmetry in the poloidal direction. The dependence of the propagation velocity on the blob size was investigated. When the initial blob size δ_b is large, the propagation velocity scales as δ_b^{-2} . In contrast, when the initial blob size is small, the blob is broken into the pieces due to the Kelvin-Helmholtz instability, and the size of blob approaches to $\delta_{\text{crit}} = \rho_s \left(L_{\parallel}^2 / 2\rho_s R \right)^{1/5}$.

In Chapter 5, the turbulent transport in SoL is investigated using the flux-driven nonlinear simulation that employs the 2D Hasegawa-Wakatani-type resistive interchange model. Linear stability is analyzed, and influences of coefficients α , β , D and μ are studied. Spontaneous generation of blobs is demonstrated. During the turbulence evolution, radially-elongated meso-scale convective cells are formed autonomously and they dominate the radial transport in SoL region. Scaling law of velocity is obtained as $V_x \sim (\rho_s/R)^{2/5} c_s$. The effective transport coefficient is found to obey the Bohm-like dependence $D_{\text{eff}} \propto (L_{\parallel}/R)^{2/5} T/eB$. The origin of the parameter dependence of the effective transport coefficient is investigated. The velocity, amplitude, and the frequency of appearance of blobs are controlling parameters of D_{eff} . The nonlinear simulation clarifies the 'non local' feature of the transport in SoL region.

The formula of blob dynamics and associated transport are clarified in this thesis. The radial motion of blob-like structures in the SoL of magnetically confined plasmas causes large cross-field transport of particles and heat, and strongly influences the properties of the SoL and divertor operation. In particular, the blobs cause a broadening of the radial plasma profile, induce en-

hancement of plasma-wall interaction, and thereby increase erosion of (and generation of impurities at) the main chamber walls. The transport mediated by such blob structures is likely related to a number of outstanding problems for the steady operation of a fusion experiment, including recycling at main chamber, material migration, plasma rotation and the density disruption limit.

The results in Chapters 3-5 have clarified the role of long-range perturbations and ‘non local’ transport. These findings stimulate the study of the edge plasma, where the long-range fluctuations can have influences on the transport. For the extension of the studies in Chapter 5, we performed a nonlinear simulation of turbulence in the edge by flux-driven model in Chapter 6. We here studied the global transport simulation of edge plasma by flux-driven nonlinear simulations. The gradient-flux relation is fitted to the model with parameters of stiffness and critical gradient, eq. (6.4). The slope of the grad- p -v.s.-flux-curve is clearly lower in the presence of turbulence-generated rotation, compared to the case where rotation is artificially suppressed. The origins of this difference are studied by imposing specified velocity. The slopes of Q_{tot} v.s. $\overline{\nabla p}$ are similar for cases with different amplitudes of controlled rotation. That is, the stiffness does not change significantly, while the critical gradient changes. The change in the slope observed in the case of spontaneous rotation has to be attributed to the fact that the velocity shear itself is increasing with the pressure source. A tentative conclusion is given that the origin of the effect on the gradient-flux curve given by the spontaneous rotation is found to be an increment of critical gradient.

Future Work

The model derived in Chapter 3 and used in Chapters 4 and 5 is an approximate one in 2D modelling. In this model, the physical quantities except the current are assumed to be uniform in the parallel direction. This model is valid to

analyze the condition that the radial transport of blob is much larger than the parallel loss. We analyzed this 2D model and obtained the propagation characteristics of an isolated blob, the transport velocity dominated by blobs in the SoL plasma, characteristics of the effective blobby transport coefficient, etc.

In the next step of the study, the extension to 3D model is necessary, to include parallel dynamics. In the 3D simulation, a couple of 3D effects on the blob behavior are expected: 1) During the long time evolution of blobs, the density of blobs reduces because of the parallel loss. 2) the connection length in the SoL, $L_{\parallel}$, changes with the change of radial coordinate. This has an effect on the parameter $\alpha = 2\rho_s/L_{\parallel}$. 3) the flux tube along the magnetic field line may be deformed with the time evolution and this affects the curvature parameter β .

Some divertor physics are expected to affect the parameter of the conductivity at divertor plates α ; For instance, the neutral effect, particle vacuum, etc., can influence α . In addition, we assumed the velocity distribution of the electron flux at the sheath at the divertor plate to be Maxwellian to derive the α parameter. This simplification might be corrected as well.

Finally, the theoretical results obtained in this thesis must be tested by experiments. An initial comparison has been made on the blob velocity by, e.g., the TORPEX Group [118]. An initial report was promising, so that detailed comparison is looked for. The future comparison includes the correlation length of flux, and so on.

These are left for future work.

Appendix A

Dependence on Model and Numerical Scheme

A.1 Partially Linearized Model

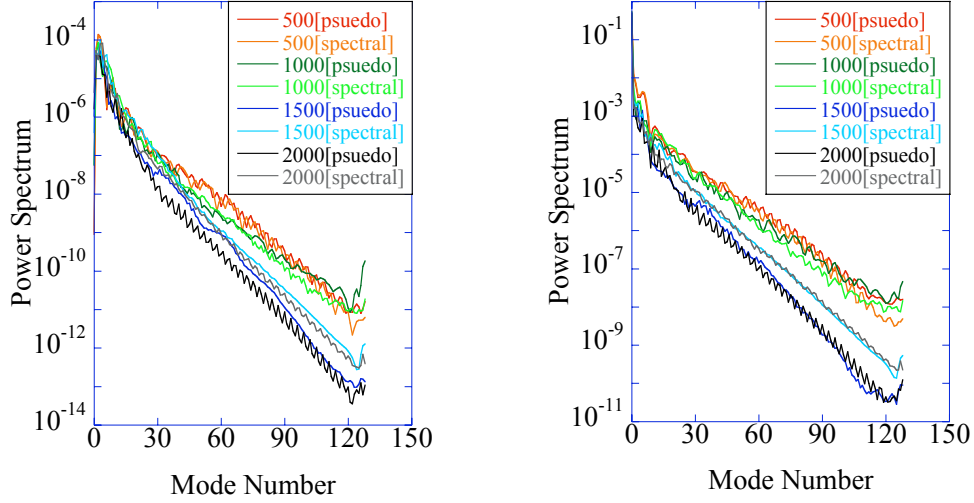
In this section, we show the simulation results solving the equations (3.30) and (3.31), in which the nonlinear terms of density are linearized: L-model. The partially linearized model equations can be solved by both spectral and pseudo-spectral methods. The pseudo-spectral code is compared with the spectral code for code validation, then we use the pseudo-spectral code for simulations demonstrated in the subsequent sections.

A.1.1 Comparison between spectral and pseudo-spectral codes

In this subsection, we compare the simulation results of the power spectra when spectral and pseudo-spectral codes are employed.

Figures A.1 shows the power spectra of electrostatic potential and density in the small blob simulations ($\delta_b = 5$). In the early phase of the nonlinear evolution, the power spectra coincide with each other. As time goes by, the difference between the spectral method and the pseudo-spectral method is observable. However, the difference is negligible in low mode region, so that we

may consider that both codes work correctly.



(a) Power spectra of the electrostatic potential in the small blob propagation simulations solved by the spectral and the pseudo-spectral codes.

(b) Power spectra of the plasma density in the small blob propagation simulations solved by the spectral and the pseudo-spectral codes.

Fig. A.1 Line colors refer to the time step of the simulation results by the spectral and the pseudo-spectral codes. Vertical axis and horizontal axis indicate the energy of the power spectra and the Fourier mode number in the poloidal direction, respectively.

Figure A.2 shows the time evolution of the density energy for the largest 10 Fourier modes. Although the similar behavior is observed for the spectral and the pseudo-spectral code for $t < 1000$, the dissipation in the results of pseudo-spectral code are somewhat stronger than that in the spectral code. For $t > 1000$, the complicated behavior of high modes is observed in the pseudo-spectral code, which is different from that in the spectral code. This behavior might be related with the chaotic nature of the system. Since the dealiasing is not applied in this version of the pseudo-spectral code, there exists noise in the high k_y region, which might lead to different behavior of high k_y modes due to

the sensitivity of a initial value.

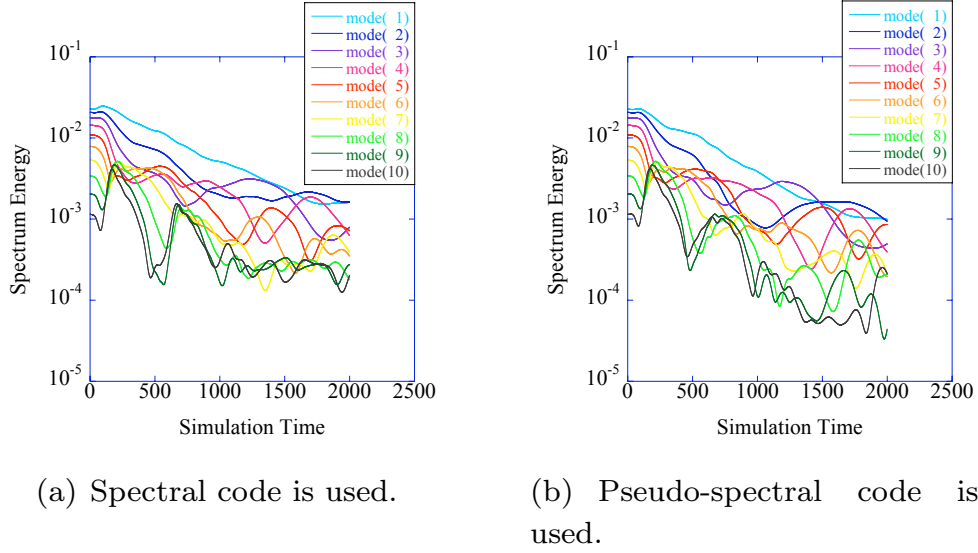
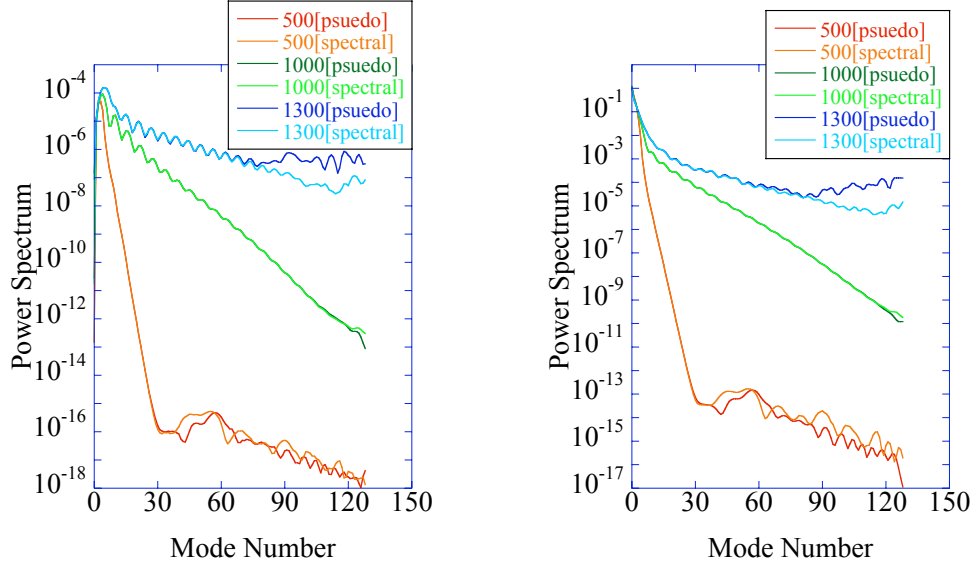


Fig. A.2 Time evolution of the spectrum energies of the plasma density in the small blob ($\delta_b = 5$) propagation simulations. Line colors refer to the spectral and the pseudo-spectral mode numbers.

Figure A.3 shows the power spectra in the large blob simulations ($\delta_b = 60$). In this case, the agreement between two codes is fairly well compared to the case with the small blob. However, we can see the accumulation of energies in the high-mode region as time goes by, in the pseudo-spectral case. Some care should be taken into account to get rid of the aliasing error in the pseudo-spectral code for long time simulations.



(a) Power spectra of the electrostatic potential in the large blob propagation simulations solved by the spectral and the pseudo-spectral codes.

(b) Power spectra of the plasma density in the large blob propagation simulations solved by the spectral and the pseudo-spectral codes.

Fig. A.3 Line colors refer to the time step of the simulation results by the spectral and the pseudo-spectral codes. Vertical axis and horizontal axis have the same meanings as in fig. A.1.

A.1.2 Evolution of small blobs

In the previous section, we validate the spectral and the pseudo-spectral codes. In this subsection, we focus on our simulation results of the blob propagation in the SoL using the pseudo-spectral code. The blob instabilities that occur during propagation in the SoL and are affected by the blob size are investigated by Aydemir [67]. The characteristic blob size δ_{crit} is introduced, which gives the categorization of the type of blob instabilities

$$\delta_{\text{crit}} = \rho_s \left(\frac{L^2}{2\rho_s R} \right)^{1/5}, \quad (\text{A.1})$$

where R is the major radius, $L = qR$ is the connection length of the magnetic field line in the SoL.

Using large tokamak parameters: $c_s = 6 \times 10^4 \text{m/sec}$, $\Omega_i = 10^8 \text{sec}^{-1}$, $L = 40 \text{m}$, and $R = 2 \text{m}$ (similar to the parameters used by Yu and Krasheninnikov [61]), we get $\alpha = 3 \times 10^{-5}$, and $\beta = 6 \times 10^{-4}$. These gives the dimensionless length scale $\delta_{\text{crit}} = 15$ (in units of $\rho_s = 6 \times 10^{-4} \text{m}$). We simulate the process of nonlinear evolution in blob propagation with respect to δ_{crit} .

First, we show the result of small blob propagation in fig. A.4. Here, we take the blob size as $\delta_b = 5$ ($\delta_b/\delta_{\text{crit}} < 1$). The rapid convection occurs and the velocity shear is generated, which gives rise to the formation of Kelvin-Helmholtz vortices and vortex shedding. As a result, small blobs do not survive as coherent structures.

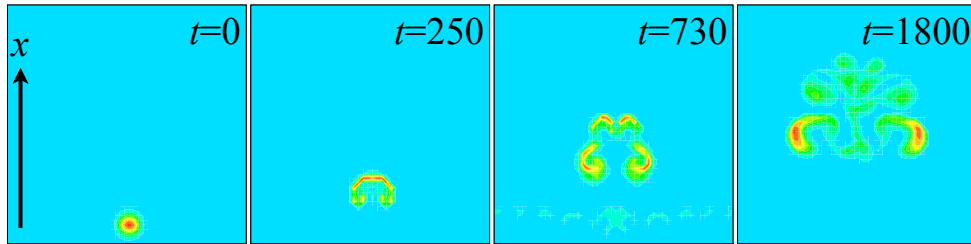


Fig. A.4 Nonlinear evolution of the small blob ($\delta_b=5$) with the partially linearized model. In this and following figures, the bottom boundary represents the separatrix (or the last closed flux surface). The horizontal and vertical coordinates are the poloidal and radial directions, respectively. The red and yellow colors correspond to high density, blue and green to low. The field size of this result is 128 in x and is 128 in y . The lengths are in units of $\rho_s = 6 \times 10^{-4} \text{m}$. Time is normalized with $\Omega_i^{-1} = 1 \times 10^{-8} \text{sec}$.

Figure A.5 shows the radial density profile of $m = 0$ mode in the early stage of nonlinear evolution. In this stage, the blob remains as a coherent structure. The initial Gaussian profile is advected radially as a pulse with a front and tail. The peak of the Gaussian profile convects rapidly due to the formation of

large electrostatic potential which gives rise to the large $\mathbf{E} \times \mathbf{B}$ drift velocity. Accompanied with its effect, the shape of the density profile gets gradually steepening. The radial profile of the propagating blob resembles “shock-wave” solutions of the Burgers’ equation, as mentioned in earlier works [67], [59].

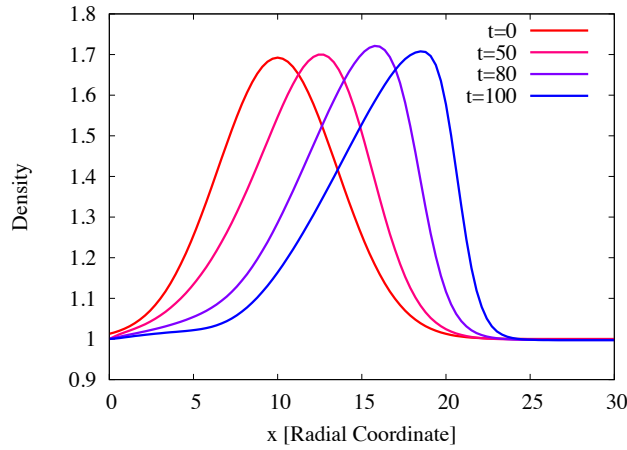


Fig. A.5 Radial density profile of $m = 0$ mode in the small blob ($\delta_b = 5$) propagation using the partially linearized model equations. In this stage of evolution, the blob form remains the coherent structure.

A.1.3 Evolution of large blobs

Large blobs are not affected by Kelvin-Helmholtz instabilities since their convective velocities are small. Figure A.6 shows the simulation results, with $\delta_b = 60$ ($\delta_b/\delta_{\text{crit}} \gg 1$). It is seen that “Rayleigh-Taylor fingers” develop. In this case, interchange instabilities play a dominant role in the large blob propagation. We also find that rapid dissipation occurs.

During the evolution of large blobs, their initial structure survives longer than that of small blobs. The peak value decreases according to the deformation of blob due to interchange instability. The steepening effect is observed here, too.

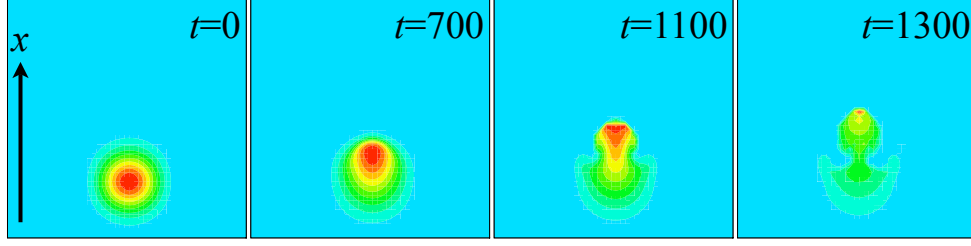


Fig. A.6 Nonlinear evolution of the large blob of size $\delta_b = 60$ with the partially linearized model. The other parameters are same as in fig. A.4. The field size is 512 in each direction.

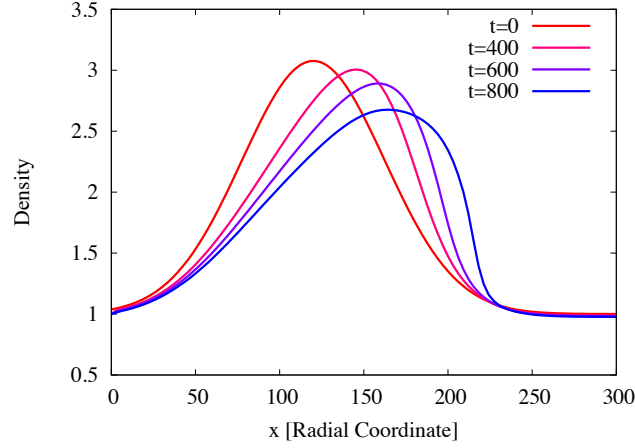


Fig. A.7 Radial density profile of $m = 0$ mode in the large blob ($\delta_b = 60$) propagation using the partially linearized model equations. In this stage of evolution, the blob form remains the coherent structure.

A.2 Full Nonlinear Model

In this section, we investigate the effect of density nonlinearity in eqs. (3.22) and (3.23): C-model. Such nonlinearity can be treated by the pseudo-spectral method. We compare the results with those in the previous section and clarify the effect of density nonlinearity. Simulation results are discussed in detail in the following subsections.

A.2.1 Evolution of small blobs

The Kelvin-Helmholtz vortices and the instabilities are observed in fig. A.8. The time scale of propagation is quite slow in the case of density nonlinearity. For example, the blob structure at $t = 600$ in fig. A.8 indicates the case which is similar to the one at $t = 200$ in fig. A.4. Comparing figs. A.8 and A.4, one finds that the density nonlinearity has a tendency to slow down the evolution of Kelvin-Helmholtz instability. The vortex breaks up into pieces at $t = 1100$. In fig. A.4, vortices broken up are dragged by the front head of vortex and form a complicated convective structure. On the other hand, broken vortices are left behind (and then dissipated) in fig. A.8.

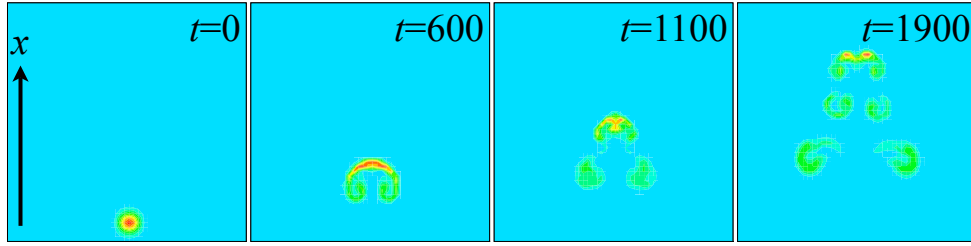


Fig. A.8 Nonlinear evolution of small blob ($\delta_b = 5$, $\delta_b/\delta_{\text{crit}} < 1$) with full nonlinear model equations. Other parameters are the same as in fig. A.4. The system size is 128 in each direction.

Figure A.9 shows the radial density profile of $m = 0$ mode in this simulation. This figure indicates the effect of density nonlinearity that makes the blob propagation velocity slow. For example, the x coordinate of the peak in the radial density profile at $t = 80$ in fig. A.9 is clearly smaller than that in fig. A.5. We can see that the steepening effect is weaker in this full nonlinear case. For example, the heads and tails of the blob radial profile at $t = 100$ in fig. A.5 and at $t = 160$ in fig. A.9 bear a strong resemblance each other, however the peak of the profile in fig. A.9 clearly shows decay.

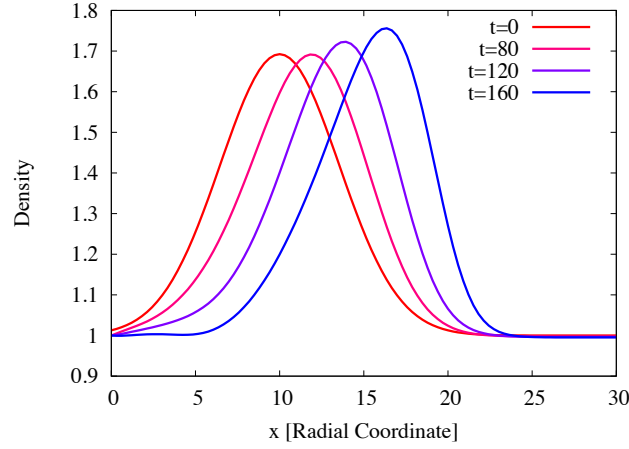


Fig. A.9 Radial density profile of $m = 0$ mode in the small blob ($\delta_b = 5$) propagation using the full nonlinear model in the early stage of nonlinear evolution. At this stage, the blob remains a coherent structure.

We focus on the blob propagation velocities here. In fig. A.10, the location of the peak of density shows (over the range $t = 0$ -100 in the partially linearized model and $t = 0$ -160 in the full nonlinear model) of small blob propagation. In these ranges, the blob keeps the coherent structure. (Because the peak position is indicated by the grid coordinate in x direction, the point sets are not smooth.) We introduce a quadratic function to fit these data sets with the least-square method. The equations and parameters used are follows:

$$\begin{aligned}
 f(t) = & (6.43168 \times 10^{-4} \pm 1.019 \times 10^{-5}) t^2 \\
 & + (2.33722 \times 10^{-2} \pm 1.053 \times 10^{-3}) t \\
 & + (9.80925 \pm 2.278 \times 10^{-2}) \quad (\text{A.2})
 \end{aligned}$$

$$\begin{aligned}
 g(t) = & (1.98346 \times 10^{-4} \pm 3.149 \times 10^{-6}) t^2 \\
 & + (9.32051 \times 10^{-3} \pm 5.208 \times 10^{-4}) t \\
 & + (9.88098 \pm 1.803 \times 10^{-2}) \quad (\text{A.3})
 \end{aligned}$$

where $\pm$ indicates the error, $f(t)$ fits the partially linearized model data set,

and $g(t)$ fits the full nonlinear model data set. The velocity is given by time derivative of $f(t)$ and $g(t)$. Equations (A.2) and (A.3) are given by the initial behavior. We observe the propagation velocities of small blobs which relax to a constant value after enough time passed. For the model without density nonlinearity, we have $\dot{f}(80) \sim 1.26 \times 10^{-1}$ and for the model with density nonlinearity, we have $\dot{g}(140) \sim 6.49 \times 10^{-2}$, where ‘ $\cdot$ ’ means time derivative. This result indicates that the propagation velocity with density nonlinearity is about 3 times slower than that without the density nonlinearity in the case of the small blobs. These estimates are used in next subsection to compare the large and small blobs propagation velocities based on the equation (2.18).

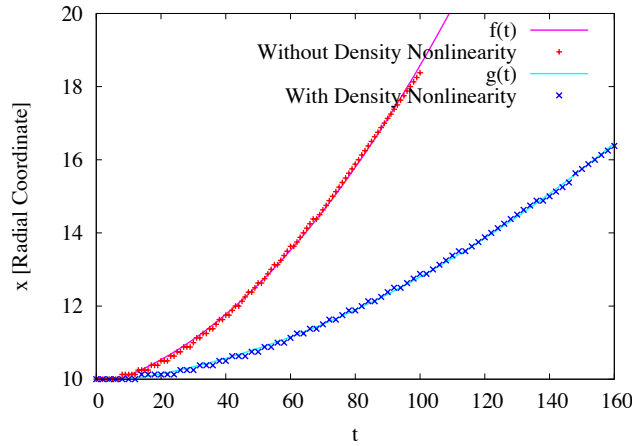


Fig. A.10 The peak locations of $m = 0$ mode for the density versus time, fitted with the quadratic functions using the least-square method.

A.2.2 Evolution of large blobs

Figure A.11 shows the nonlinear evolution of large blob in the case with the density nonlinearity. The propagation velocity is slow enough to develop the Rayleigh-Taylor fingers in this case.

The steepening effect weakly appears in this case as well, compared to the results in fig. A.7. The blob keeps the coherent structure of itself in the time

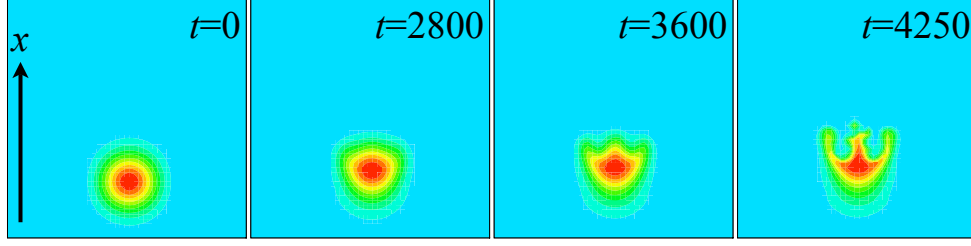


Fig. A.11 Nonlinear evolution of the large blob ($\delta_b = 60$, $\delta_b/\delta_{\text{crit}} \gg 1$) with the full nonlinear model. Other parameters are same as in fig. A.4. The system size is 512 in each direction.

stage shown in fig. A.12. After the last time step, they lose their coherent structure because of interchange instability.

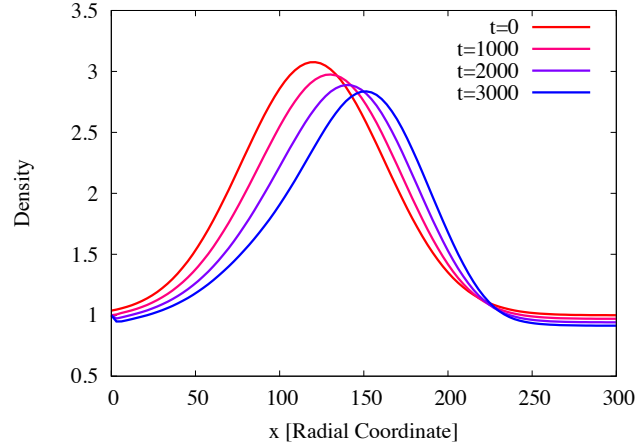


Fig. A.12 Radial density profile of $m = 0$ mode in the large blob ($\delta_b = 60$) propagation using the full nonlinear model in the early stage of nonlinear evolution. At this stage, the blob remains a coherent structure.

Here, we have the propagation velocity of large blobs which is estimated as in previous subsection. In fig. A.13, the point set shows the time series of the peak of density (over the range $t = 0-600$ in the partially linearized model and $t = 0-3000$ in the full nonlinear model) of large blob propagation. In these ranges, the blob keeps the coherent structure. The quadratic fitting to the data

in the least-square method are as follows:

$$\begin{aligned}
 h(t) = & (7.50336 \times 10^{-6} \pm 2.048 \times 10^{-6}) t^2 \\
 & + (6.37899 \times 10^{-2} \pm 1.27 \times 10^{-3}) t \\
 & + (118.365 \pm 0.1648) \quad (\text{A.4})
 \end{aligned}$$

$$\begin{aligned}
 k(t) = & (2.28985 \times 10^{-7} \pm 2.041 \times 10^{-8}) t^2 \\
 & + (9.65625 \times 10^{-3} \pm 6.329 \times 10^{-5}) t \\
 & + (120.077 \pm 4.107 \times 10^{-2}) \quad (\text{A.5})
 \end{aligned}$$

where $\pm$ indicates the error, $h(t)$ fit the partially linearized model data set, and $k(t)$ fit the full nonlinear model data set. The velocity is given by time derivative of $h(t)$ and $k(t)$.

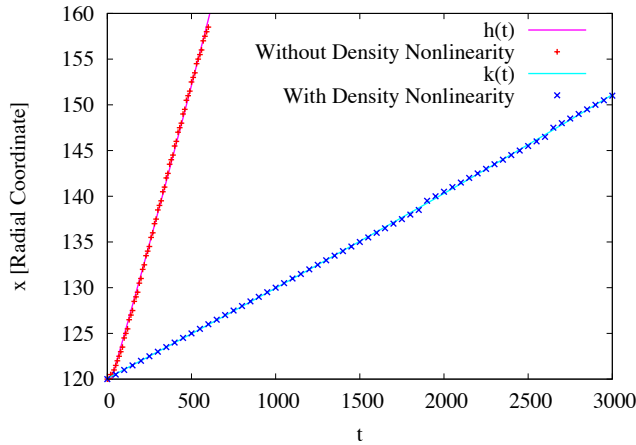


Fig. A.13 The peak location of $m = 0$ mode for the density versus time, fitted with the quadratic functions using the least-square method.

As is in previous subsection, we also observe that large blobs reach constant velocities. For the model without density nonlinearity, we have $\dot{h}(80) \sim 6.5 \times 10^{-2}$ and, for the model with density nonlinearity, we have $\dot{k}(140) \sim 9.72 \times 10^{-3}$. These result indicate that the propagation velocity with the density

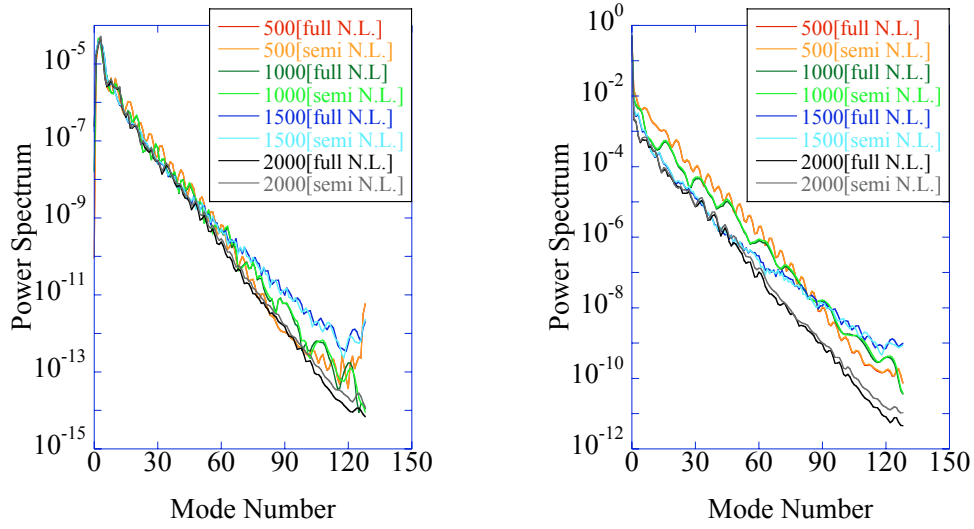
nonlinearity is about 6.5 times slower than that without the density nonlinearity in the case of large blobs.

Now we consider the blob propagation velocity equation (2.18), which indicates that the propagation velocity is inversely proportional to square of blob size. In this work, we use $\delta_{\text{b small}} = 5$ and $\delta_{\text{b large}} = 60$, thus the velocity ratio must be $V_{x \text{ large}}/V_{x \text{ small}} \sim 12^2$. We compare this simple estimate with simulation results. In the later phase of evolution, the velocity ratio between small and large blobs (models without density nonlinearity) is $\dot{h}(80)/\dot{f}(80) \sim 1.9$ and the ratio is $\dot{k}(140)/\dot{g}(140) \sim 13$ (for the model with density nonlinearity). Both of them are far from linear estimate of eq. (2.18). This is probably because the derivation of eq. (2.18) is based on the assumption that the shape of the blob is given and fixed. In the real simulation, the shape of blob is strongly deformed. Taking this effect into consideration, comparison of the observed velocity with the formula is given in the main text.

A.3 Detailed Discussion on Effects of Density Nonlinearity

In the previous section, we discussed on the effects of the density nonlinearity. However it is not revealed that which density nonlinearity is effective, $n^{-1} \partial n / \partial y$ in eq. (3.22) or $n \partial \varphi / \partial y$ in eq. (3.23). We show the comparison of the power spectra with the full nonlinear model ((3.22) and (3.23)) and with the semi-nonlinear model, in which $n \partial \varphi / \partial y$ term is replaced by $\partial \varphi / \partial y$ in the continuity equation. Figures A.14 and A.15 show the power spectra obtained by the small and the large blob propagation simulations, respectively, with full nonlinear model and semi-nonlinear model.

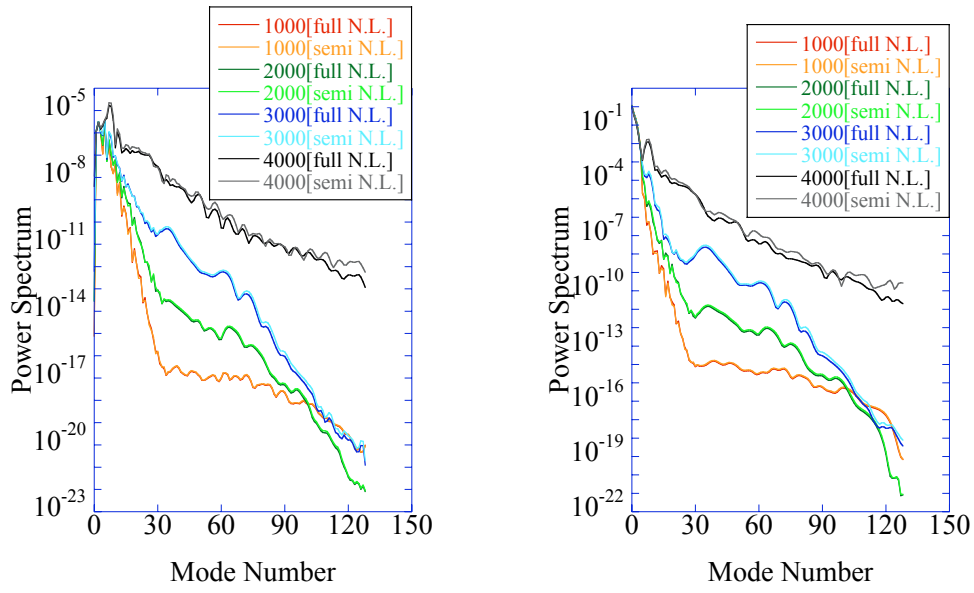
The good agreement of the power spectra in each time step is observed, which indicates the dominant contribution of the density nonlinearity is $n^{-1} \partial n / \partial y$ in the vorticity equation.



(a) Power spectra of the electrostatic potential by the small blob propagation simulations with the full nonlinear model and the semi-nonlinear model.

(b) Power spectra of the density by the small blob propagation simulations with the full nonlinear model and the semi-nonlinear model.

Fig. A.14 Line colors refer to the time step of the full nonlinear model and the semi-nonlinear model. Vertical and horizontal axes indicates the strength of the power spectrum and the Fourier mode number in the poloidal direction, respectively.



(a) Power spectra of the electrostatic potential by the large blob propagation simulations with full nonlinear model and semi-nonlinear model.

(b) Power spectra of the plasma density by the large blob propagation simulations with full nonlinear model and semi-nonlinear model.

Fig. A.15 Line colors refer to the time step of the full nonlinear model and the semi-nonlinear model. Vertical and horizontal axes have the same meanings as in fig. A.14.

Appendix B

Spectral and Pseudo-Spectral Method

In this thesis, we mainly use the pseudo-spectral code to solve model equations. In Appendix A, we have compared the spectral and the pseudo-spectral codes observing power spectra. In this appendix, we briefly explain how to solve a nonlinear equation using spectral and pseudo-spectral methods.

To explain both methods, we employ the 1-dimensional nonlinear advection equation, as an example,

$$\frac{\partial u}{\partial t} + u \frac{\partial u}{\partial x} = 0, \quad (\text{B.1})$$

where the following conditions are assumed:

$$\text{Initial Condition:} \quad u(x, 0) = f(x), \quad (\text{B.2})$$

$$\text{Boundary Condition:} \quad u(x, t) = u(x + 2\pi, t). \quad (\text{B.3})$$

Expanding $u(x, t)$ in eq. (B.1) with Fourier series,

$$u(x, t) = \sum_{k=-N}^N \hat{u}_k(t) e^{ikx}, \quad (\text{B.4})$$

where N is the upper bound of the retained components, which determines the numerical resolution. Substituting eq. (B.4) into the L.H.S. of eq. (B.1), the

residual R is given by

$$R(x, t) = \sum_{k=-N}^N \frac{d\hat{u}_k}{dt} e^{ikx} + \sum_{k=-N}^N \sum_{l=-N}^N il\hat{u}_l\hat{u}_k \delta e^{ikx} e^{ilx}. \quad (\text{B.5})$$

Multiplying the weighting function e^{-imx} to both sides, integrating 0 to 2π with x , and assuming the integral of the residual is 0, then we obtain the ordinary differential equation,

$$\frac{d\hat{u}_m}{dt} + \sum_{k=-N}^N \sum_{l=-N}^N il\hat{u}_l\hat{u}_k \delta_{k+l,m} = 0. \quad (\text{B.6})$$

By taking the Kronecker delta into account eq. (B.6) is expressed as

$$\frac{d\hat{u}_m}{dt} + \sum_{l=\max(-N, -N+m)}^{\min(N, N+m)} il\hat{u}_{m-l}\hat{u}_l = 0 \quad (m = -N, \dots, N), \quad (\text{B.7})$$

where the upper and lower limits of the summation on l are derived by taking account of the fact that the range of k and l varies as $-N \leq k, l \leq N$ and by using the relation $\delta_{k+l,m} = 1$ when $k = m - l$.

For the spectral method, the nonlinear term in eq. (B.7) is evaluated in the Fourier transform and $O(N^2)$ operations are needed for the summation. However, the use of pseudo-spectral transform methods enables the evaluation in $O(N \log_2 N)$ operations. This technique was developed independently by Orszag [125, 126] and Eliassen, Machenhauer and Rasmussen [127]. The scheme is illustrated as:

1. Discrete inverse Fourier transforms of $\hat{u}_k$ and $ik\hat{u}_k$

$$u_j = \sum_{k=-N}^N \hat{u}_k e^{ikx_j} \quad (j = 0, \dots, J-1), \quad (\text{B.8})$$

$$\left(\frac{\partial u}{\partial x} \right)_j = \sum_{k=-N}^N ik\hat{u}_k e^{ikx_j} \quad (j = 0, \dots, J-1) \quad (\text{B.9})$$

2. estimation of the nonlinear term in the real space

$$F_j = \left(u \frac{\partial u}{\partial x} \right)_j = u_j \left(\frac{\partial u}{\partial x} \right)_j \quad (j = 0, \dots, J-1) \quad (\text{B.10})$$

3. then discrete Fourier transform of F_j

$$\hat{F}_k = \frac{1}{J} \sum_{j=0}^{J-1} F_j e^{-ikx_j} \quad (k = -N, \dots, N) \quad (\text{B.11})$$

Appendix C

Linear Stability Analysis of Interchange Mode in SoL

In this Appendix, we discuss the linear stability analysis of the interchange mode in SoL plasma. Equation (5.1) in Chapter 5 shows the dispersion relation, which is derived from model equations (3.22) and (3.23). Here, we reproduce the eq. (5.1) as follows:

$$k_{\perp}^2 \left(\omega + i\hat{\mu} + \frac{i\alpha}{k_{\perp}^2} \right) \left(\omega + i\hat{D} + i\alpha - \beta k_y \right) + k_y^2 (\beta\Lambda - \beta^2) = 0, \quad (\text{C.1})$$

where $\hat{\mu} = \mu k_{\perp}^2$, $\hat{D} = Dk_{\perp}^2$, and $\Lambda \equiv 1/\lambda$.

We first illustrate the interchange instability, which is driven by the β term. Taking a limit of $\alpha \rightarrow 0$, then the equation becomes

$$\omega^2 + \left(+i\hat{\mu} + i\hat{D} - \beta k_y \right) \omega - \hat{\mu} \left(\hat{D} + i\beta k_y \right) + \frac{k_y^2}{k_{\perp}^2} (\beta\Lambda - \beta^2) = 0. \quad (\text{C.2})$$

Solving eq. (C.2) for ω we have

$$\omega = \frac{-i \left(\hat{\mu} + \hat{D} \right) + \beta k_y}{2} + \frac{i}{2} \sqrt{\frac{4k_y^2}{k_{\perp}^2} (\beta\Lambda - \beta^2) + \left(-\hat{\mu} + \hat{D} + i\beta k_y \right)^2}. \quad (\text{C.3})$$

The βk_y term indicates the finite real frequency, which is caused by the ∇B drift. Usually, the relation $\Lambda \gg \beta$ (i.e., the density gradient is larger than

$\nabla B/B$) can be assumed. In a region of small k_y , we put $\mu = D$ for the simplicity and have

$$\omega = \frac{\beta k_y}{2} + i \left(\frac{k_y}{k_\perp} \sqrt{\beta \Lambda} - D k_\perp^2 \right). \quad (\text{C.4})$$

The linear growth rate of the interchange instability is given by

$$\gamma = \text{Im}(\omega) \sim \frac{k_y}{k_\perp} \sqrt{\beta \Lambda} - D k_\perp^2. \quad (\text{C.5})$$

The interchange instability is linearly unstable for $k_\perp < k_{\perp c}$, where

$$k_{\perp c}^2 = \frac{1}{D} \frac{k_y}{k_\perp} \sqrt{\beta \Lambda}. \quad (\text{C.6})$$

We here note that if one of D or μ vanishes, $k_{\perp c}$ is given by the different expression.

Next, we consider the finite α effect. The eigenfrequency is obtained by putting

$$\begin{cases} \mu k_\perp^2 \rightarrow \mu k_\perp^2 + \alpha/k_\perp^2 \\ D k_\perp^2 \rightarrow D k_\perp^2 + \alpha \end{cases} \quad (\text{C.7})$$

in the previous formula (C.3). For the transparency of argument, we put $\mu = D$ and $k_\perp \sim k_y$. In this limiting case, we have

$$\begin{aligned} \omega = & \frac{-i(2Dk_\perp^2 + \alpha + \alpha/k_\perp^2) + \beta k_y}{2} \\ & + \frac{i}{2} \sqrt{\frac{4k_y^2}{k_\perp^2} (\beta \Lambda - \beta^2) + \left(\alpha \left(1 - \frac{1}{k_\perp^2} \right) + i\beta k_y \right)^2}. \end{aligned} \quad (\text{C.8})$$

Away from the stability boundary, the first term is dominant in the $\sqrt{\quad}$ term for small α . One has the linear growth rate of the interchange instability:

$$\gamma = \text{Im}(\omega) \sim \frac{k_y}{k_\perp} \sqrt{\beta \Lambda} - \left(D k_\perp^2 + \frac{\alpha}{2} + \frac{\alpha}{2k_\perp^2} \right). \quad (\text{C.9})$$

The stabilizing term becomes minimum at

$$k_{\perp} = \left(\frac{\alpha}{2D} \right)^{1/4}. \quad (\text{C.10})$$

The wave number for the most unstable mode eq. (C.10) is weakly dependent on β . The maximum growth rate is given from eq. (C.9) as

$$\gamma_{\max} \simeq \sqrt{\beta\Lambda} - \sqrt{2\alpha D} - \frac{\alpha}{2} + \dots. \quad (\text{C.11})$$

The following relation is given for the stability boundary:

$$\beta\Lambda = \alpha D, \quad (\text{C.12})$$

for the case of the small- α limit. Equation (C.9) shows that D and α stabilize the linear instability of the interchange mode in SoL. This argument agrees with the indication that is given by figs. 5.2 and 5.3 in Chapter 5. Equation (C.10) gives the wave number for the most unstable mode. This wave number is proportional to α and is inversely proportional to D . If α increases, the wave number increases. If D increases, the wave number decreases. This argument also agrees with the indication of figs. 5.2 and 5.3.

More general argument is made for the stability boundary. If one neglects the weak stabilization by the finite real frequency βk_y (the last term in the R.H.S. of eq. (C.8)), then the marginal stability condition $\gamma = 0$ in eq. (C.8) is written as

$$\frac{4k_y^2}{k_{\perp}^2} \beta\Lambda + \left(1 - \frac{1}{k_{\perp}^2} \right)^2 \alpha^2 = \left(2Dk_{\perp}^2 + \alpha \left(1 + \frac{1}{k_{\perp}^2} \right) \right)^2, \quad (\text{C.13})$$

i.e.,

$$\frac{k_y^2}{k_{\perp}^2} \beta\Lambda = D^2 k_{\perp}^4 + D (1 + k_{\perp}^2) \alpha + \frac{\alpha^2}{k_{\perp}^2}. \quad (\text{C.14})$$

The critical condition is deduced from eq. (C.14). Setting $z \equiv \alpha/D$, we have

$$\frac{k_y^2}{k_{\perp}^2} \frac{\beta\Lambda}{D^2} = k_{\perp}^4 + (1 + k_{\perp}^2) z + \frac{z^2}{k_{\perp}^2}. \quad (\text{C.15})$$

If $z \ll 1$, the R.H.S. of eq. (C.15) takes the minimum at

$$k_{\perp}^2 = \frac{z^{2/3}}{2^{1/3}}, \quad (\text{C.16})$$

which gives

$$[\text{the R.H.S. of eq. (C.15)}] = \frac{1}{2^{2/3}} z^{4/3} + z + \frac{1}{2^{1/3}} z^{5/3} + z^{4/3}, \quad (\text{C.17})$$

$$\approx z. \quad (\text{C.18})$$

If $z \gg 1$, the R.H.S. of eq. (C.15) takes the minimum at

$$k_{\perp}^2 = \sqrt{z} \quad (\text{C.19})$$

which gives

$$[\text{the R.H.S. of eq. (C.15)}] \approx 2z^{3/2}. \quad (\text{C.20})$$

Thus, we have a more complete expression for the stability boundary as

$$\frac{\beta\Lambda}{D^2} = \begin{cases} \alpha/D & \text{if } \alpha/D \ll 1 \\ 2(\alpha/D)^{3/2} & \text{if } \alpha/D \gg 1. \end{cases} \quad (\text{C.21a})$$

$$(\text{C.21b})$$

The stability boundary for the small α/D limit, eq. (C.21a), agrees with eq. (C.12). The stability boundary is illustrated on the (α, D) plane, with fixed β and λ values in fig. C.1.

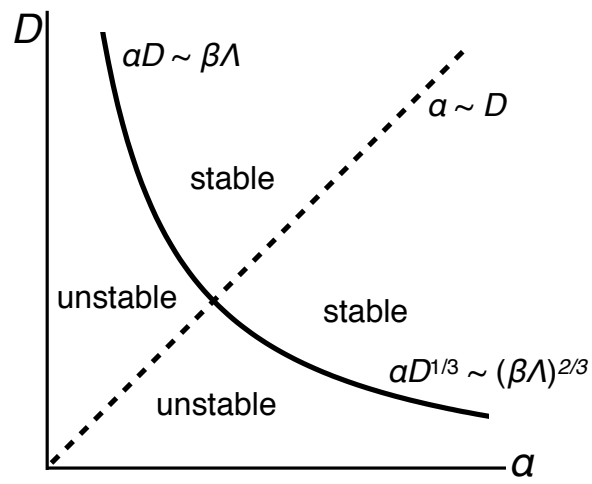


Fig. C.1 Linear stability boundary on (α, D) plane with fixed β and λ .

Acknowledgements

Firstly, I would like to express my sincere gratitude to my supervisor, Professor Sanae-I. Itoh, for providing me this precious study opportunity as a Ph.D student of Kyushu University in her laboratory. She has given me continuous encouragement and thought-provoking advice on plasma physics. I would like to deeply appreciate Professor Masatoshi Yagi for a warm and heartfelt guidance during the course of my research. He guided me to the numerical simulation of fusion, supervised me with my Master thesis and has continued to be the excellent adviser during my Ph.D course.

During my Doctor course, I had a chance to study at Provence University in France one year through great supports from Professor Sanae-I. Itoh and Professor Sadruddin Benkadda. I appreciate them to give me the opportunity. I am deeply grateful to Professor Sadruddin Benkadda for his constructive comments, warm encouragement, and continuous help of my life in France. I am indebted to Professor Peter Beyer for his elaborated guidance and invaluable discussion with my study in France.

I owe a very important debt to Professor Kimitaka Itoh for his encouragement and highly instructive guidance. I would like to express my appreciation to Professor Hideki Zushi and Professor Keisuke Matsuoka for their insightful comments and useful discussions.

I would like to thank Professor Akihide Fujisawa and Associate Professor Shigeru Inagaki for their illuminating comments and discussions. I am grateful

to Dr. Makoto Sasaki and Dr. Stella Oldenb rger for their friendly support. I am also grateful to Dr. Guillaume Fuhr. He provided me with the numerical codes for use in this thesis.

Special thanks are to Dr. Xavier Garbet, Dr. Tomonori Takizuka, Dr. Masafumi Azumi, Dr. Shinji Tokuda, Professor Patrick Diamond, Professor Wendell Horton, Dr. Gloria Falchetto, Dr. Philippe Ghendrih, Dr. Paolla Mantica, Associate Professor Olivier Agullo, and Dr. Patrick Tamain for fruitful discussions and encouragements.

I am grateful to Miss Kei Asano, Miss Fumi Yonaga, Mrs. Yoshiko Hieida, and Mrs. Hiroko Sugitani, secretaries of Itoh's laboratory, for their help of miscellaneous administrative works of my campus life. I am also grateful to Dr. Shimpei Futatani and Dr. Magali Muraglia for their kind support in my stay in France.

This work is supported by the Grant-in-Aid for Specially-Promoted Research of Ministry of Education, Culture, Sports, Science and Technology of Japan (16002005), Grant-in-Aid for Scientific Research (B) of Japan Society of Promotion of Science (19360459), Grant-in-Aid for Scientific Research (B) of Japan Society of Promotion of Science (19360415), Grant-in-Aid for Scientific Research (S) of Japan Society of Promotion of Science (21224014), and Grant-in-Aid for Japan Society of Promotion of Science Fellows (22-4534).

My stay in France was financially supported by the Collage Doctoral Franco-Japonais, Interdisciplinary Graduate School of Engineering Sciences, Kyushu University, NIFS/NINS under the project of Formation of International Network, and LIA 336 CNRS/France-Japan Magnetic Fusion Laboratory.

Bibliography

- [1] S.-I. Itoh, S. Inagaki, M. Shindo, and M. Yagi. 2ND ITER INTERNATIONAL SUMMER SCHOOL: In conjunction with the 47th Summer School of JSPF for Young Plasma Scientists: Confinement. In *American Institute of Physics Conference Series*, Vol. 1095, 2009.
- [2] J. Wesson and D.J. Campbell. *Tokamaks*. Oxford University Press, USA, 2004.
- [3] S.I. Krasheninnikov. On scrape off layer plasma transport. *Physics Letters A*, Vol. 283, No. 5-6, pp. 368–370, May 2001.
- [4] S. Sugita, M. Yagi, and S.-I. Itoh. Nonlinear Simulation of Plasma Blob Formation and Radial Propagation in the Scrape-off Layer. *Reports of Research Institute for Applied Mechanics, Kyushu University*, Vol. 136, pp. 1–18, 2009.
- [5] A. Hasegawa and M. Wakatani. Plasma edge turbulence. *Physical Review Letters*, Vol. 50, No. 9, pp. 682–686, 1983.
- [6] S. Sugita, M. Yagi, S.-I. Itoh, and K. Itoh. Propagation Velocity Analysis of a Single Blob in the SOL. *Plasma and Fusion Research*, Vol. 3, p. 40, 2008.
- [7] S. Sugita, M. Yagi, S.-I. Itoh, and K. Itoh. Nonlinear Effect on the Plasma Blob Propagation in the Scrape-off Layer. *National Committee for Theoretical and Applied Mechanics, Science Council of Japan*, Vol. 57, pp. 207–215, 2009.

- [8] S. Sugita, M. Yagi, S.-I. Itoh, and K. Itoh. Interchange Turbulence and Radial Transport in Tokamak Scrape-off Layer Dominated by Meso Scale Structure. *Journal of the Physical Society of Japan*, Vol. 79, No. 4, p. 4502, 2010.
- [9] S. Sugita, M. Yagi, S.-I. Itoh, and K. Itoh. Study of Radial Particle Transport Accompanied with Plasma Blob and Self-organized Meso-scale Structure in Tokamak Scrape-off Layer. *Preprints of 23rd International Atomic Energy Agency Fusion Energy Conference*, 2010.
- [10] M. Shimada, DJ Campbell, V. Mukhovatov, M. Fujiwara, N. Kirneva, K. Lackner, M. Nagami, VD Pustovitov, N. Uckan, J. Wesley, et al. Overview and summary. *Nuclear Fusion*, Vol. 47, p. S1, 2007.
- [11] A.V. Nedospasov. Edge turbulence in tokamaks. *Journal of Nuclear Materials*, Vol. 196, pp. 90–100, 1992.
- [12] N. Ohno. Plasma Blob Transport in Scrape-Off Layer. *J. Plasma Fusion Res.*, Vol. 82, No. 4, pp. 205–214, 2006.
- [13] M Yagi, S.-I. Itoh, K. Itoh, and Diamond P.H. Disparate Scale Nonlinear Interactions in Edge Turbulence. *Contributions to Plasma Physics*, Vol. 48, No. 1-3, pp. 13–22, 2008.
- [14] O.E. Garcia. Blob Transport in the Plasma Edge: a Review. *Plasma and Fusion Research*, Vol. 4, No. 0, pp. 19–19, 2009.
- [15] N. Asakura. Recent Progress of SOL and Divertor Plasma Studies in Tokamaks. *Journal of plasma and fusion research*, Vol. 80, No. 3, pp. 190–200, 20040325.
- [16] N. Asakura and N. Ohno. Recent Progress of ELM Plasma Transport Study in SOL and Divertor. *Journal of plasma and fusion research*, Vol. 82, No. 9, pp. 582–589, 20060925.
- [17] S.J. Zweben, J.A. Boedo, O. Grulke, C. Hidalgo, B. LaBombard, R.J. Maqueda, P. Scarin, and J.L. Terry. Edge turbulence measurements in

-
- toroidal fusion devices. *Plasma Physics and Controlled Fusion*, Vol. 49, p. S1, 2007.
- [18] B.D. Scott. Tokamak edge turbulence: background theory and computation. *Plasma Physics and Controlled Fusion*, Vol. 49, p. S25, 2007.
- [19] S.I. Krasheninnikov, D.A. D'Ippolito, and J.R. Myra. Recent theoretical progress in understanding coherent structures in edge and SOL turbulence. *Journal of Plasma Physics*, Vol. 74, No. 05, pp. 679–717, 2008.
- [20] V. Naulin. Turbulent transport and the plasma edge. *Journal of Nuclear Materials*, Vol. 363, pp. 24–31, 2007.
- [21] M. Endler. Turbulent SOL transport in stellarators and tokamaks. *Journal of Nuclear Materials*, Vol. 266, pp. 84–90, 1999.
- [22] P.C. Stangeby. *The Plasma Boundary of Magnetic Fusion Devices (Series in Plasma Physics)*. Bristol: Institute of Physics, 2000.
- [23] M. Kikuchi, S. Nishio, G. Kurita, K. Tsuzuki, M. Bakhtiari, H. Kawashima, H. Takenaga, Y. Kusama, and K. Tobita. Blanket–plasma interaction in tokamaks Implication from JT-60U, JFT-2M and reactor studies. *Fusion Engineering and Design*, Vol. 81, No. 8-14F, pp. 1589–1598, 2006.
- [24] M.V. Umansky, S.I. Krasheninnikov, B. LaBombard, and J.L. Terry. Comments on particle and energy balance in the edge plasma of Alcator C-Mod. *Physics of Plasmas*, Vol. 5, No. 9, pp. 3373–3376, 1998.
- [25] A. Kirk, H.R. Wilson, R. Akers, N.J. Conway, G.F. Counsell, S.C. Cowley, J. Dowling, B. Dudson, A. Field, F. Lott, et al. Structure of ELMs in MAST and the implications for energy deposition. *Plasma Physics and Controlled Fusion*, Vol. 47, p. 315, 2005.
- [26] R.J. Maqueda, G.A. Wurden, S. Zweben, L. Roquemore, H. Kugel, D. Johnson, S. Kaye, S. Sabbagh, and R. Maingi. Edge turbulence measurements in NSTX by gas puff imaging. Vol. 72, pp. 931–934. AIP, 2001.

- [27] O. Grulke, J.L. Terry, B. LaBombard, and S.J. Zweben. Radially propagating fluctuation structures in the scrape-off layer of Alcator C-Mod. *Physics of Plasmas*, Vol. 13, No. 1, p. 012306, 2006.
- [28] J.A. Boedo, D.L. Rudakov, R.A. Moyer, G.R. McKee, R.J. Colchin, M.J. Schaffer, P.G. Stangeby, W.P. West, S.L. Allen, T.E. Evans, R.J. Fonck, E.M. Hollmann, S.I. Krasheninnikov, A.W. Leonard, W. Nevins, M.A. Mahdavi, G.D. Porter, G.R. Tynan, D.G. Whyte, and X. Xu. Transport by intermittency in the boundary of the DIII-D tokamak. *Physics of Plasmas*, Vol. 10, No. 5, pp. 1670–1677, 2003.
- [29] H. Zushi, N. Nishino, K. Hanada, H. Honma, H.Q. Liu, Y. Higashizono, M. Sakamoto, Y. Nakashima, M. Ishiguro, T. Ryoukai, S. Tashima, K. Nakamura, H. Idei, M. Hasegawa, A. Fujisawa, O. Mitarai, A. Fukuyama, Takeiri, Y. Takase, T. Maekawa, Y. Kishimoto, K. Toi, M. Kikuchi, and the QUEST group. Study of Edge Turbulence from the Open to Closed Magnetic Field Configuration during the Current Ramp-up Phase in QUEST. *Preprints of IAEA-FEC 2010*, pp. EXS/P2–22, 2010.
- [30] H. Tanaka, N. Ohno, N. Asakura, Y. Tsuji, H. Kawashima, S. Takamura, Y. Uesugi, et al. Statistical analysis of fluctuation characteristics at high- and low-field sides in L-mode SOL plasmas of JT-60U. *Nuclear Fusion*, Vol. 49, p. 065017, 2009.
- [31] N. Asakura, N. Hosogane, K. Itami, A. Sakasai, S. Sakurai, K. Shimizu, M. Shimada, H. Kubo, S. Higashijima, H. Takenaga, et al. Role of divertor geometry on detachment and core plasma performance in JT60U. *Journal of Nuclear Materials*, Vol. 266, pp. 182–188, 1999.
- [32] J.L. Terry, R. Maqueda, C.S. Pitcher, S.J. Zweben, B. LaBombard, E.S. Marmor, A.Y. Pigarov, and G. Wurden. Visible imaging of turbulence in the SOL of the Alcator C-Mod tokamak. *Journal of Nuclear Materials*,

Vol. 290, pp. 757–762, 2001.

- [33] J.A. Boedo, D. Rudakov, R. Moyer, S.I. Krasheninnikov, D. Whyte, G. McKee, G.R. Tynan, M. Schaffer, P. Stangeby, P. West, S. Allen, T. Evans, R. Fonck, E. Hollmann, A. Leonard, A. Mahdavi, G. Porter, M. Tillack, and G. Antar. Transport by intermittent convection in the boundary of the DIII-D tokamak. *Physics of Plasmas*, Vol. 8, No. 11, pp. 4826–4833, 2001.
- [34] N. Asakura, M. Takechi, N. Oyama, and T. Nakano. Parallel and radial transport of ELM plasma in the SOL and divertor of JT-60U. *Journal of Nuclear Materials*, Vol. 337, pp. 712–716, 2005.
- [35] J.M. Dewhurst, B. Hnat, N. Ohno, R.O. Dendy, S. Masuzaki, T. Morisaki, and A. Komori. Statistical properties of edge plasma turbulence in the Large Helical Device. *Plasma Physics and Controlled Fusion*, Vol. 50, p. 095013, 2008.
- [36] G.S. Xu, V. Naulin, W. Fundamenski, C. Hidalgo, J.A. Alonso, C. Silva, B. Gonçalves, A.H. Nielsen, J.J. Rasmussen, S.I. Krasheninnikov, et al. Blob/hole formation and zonal-flow generation in the edge plasma of the JET tokamak. *Nuclear Fusion*, Vol. 49, p. 092002, 2009.
- [37] Y. Xu, R.R. Weynants, M.V. Schoor, M. Vergote, S. Jachmich, M.W. Jakubowski, M. Mitri, O. Schmitz, B. Unterberg, P. Beyer, et al. Reduction of the turbulent blob transport in the scrape-off layer by a resonant magnetic perturbation in TEXTOR. *Nuclear Fusion*, Vol. 49, p. 035005, 2009.
- [38] R.J. Maqueda, D.P. Stotler, et al. Intermittent divertor filaments in the National Spherical Torus Experiment and their relation to midplane blobs. *Nuclear Fusion*, Vol. 50, p. 075002, 2010.
- [39] I. Cziegler, J.L. Terry, J.W. Hughes, and B. LaBombard. Experimental studies of edge turbulence and confinement in Alcator C-Mod. *Physics*

- of Plasmas*, Vol. 17, p. 056120, 2010.
- [40] J. Cheng, L.W. Yan, W.Y. Hong, K.J. Zhao, T. Lan, J. Qian, A.D. Liu, H.L. Zhao, Y. Liu, Q.W. Yang, et al. Statistical characterization of blob turbulence across the separatrix in HL-2A tokamak. *Plasma Physics and Controlled Fusion*, Vol. 52, p. 055003, 2010.
- [41] O.E. Garcia, J. Horacek, R.A. Pitts, A.H. Nielsen, W. Fundamenski, J.P. Graves, V. Naulin, and J.J. Rasmussen. Interchange turbulence in the TCV scrape-off layer. *Plasma Physics and Controlled Fusion*, Vol. 48, p. L1, 2006.
- [42] B. Labit, A. Diallo, A. Fasoli, I. Furno, D. Iraj, S.H. Muller, G. Plyushchev, M. Podesta, F.M. Poli, P. Ricci, et al. Statistical properties of electrostatic turbulence in toroidal magnetized plasmas. *Plasma Physics and Controlled Fusion*, Vol. 49, p. B281, 2007.
- [43] S.K. Saha and S. Chowdhury. Intermittency in the heat and particle transports in the SINP tokamak scrape-off layer. *Physics of Plasmas*, Vol. 15, p. 012305, 2008.
- [44] N. Katz, J. Egedal, W. Fox, A. Le, and M. Porkolab. Experiments on the propagation of plasma filaments. *Physical review letters*, Vol. 101, No. 1, p. 15003, 2008.
- [45] M. Podesta, A. Fasoli, B. Labit, I. Furno, P. Ricci, FM Poli, A. Diallo, SH Muller, and C. Theiler. Cross-field transport by instabilities and blobs in a magnetized toroidal plasma. *Physical review letters*, Vol. 101, No. 4, p. 45001, 2008.
- [46] C. Theiler, A. Diallo, A. Fasoli, I. Furno, B. Labit, M. Podesta, F.M. Poli, and P. Ricci. The role of the density gradient on intermittent cross-field transport events in a simple magnetized toroidal plasma. *Physics of Plasmas*, Vol. 15, p. 042303, 2008.
- [47] D.C. Pace, M. Shi, J.E. Maggs, G.J. Morales, and T.A. Carter. Exponen-

-
- tial Frequency Spectrum in Magnetized Plasmas. *Physical review letters*, Vol. 101, No. 8, p. 85001, 2008.
- [48] I. Furno, B. Labit, M. Podesta, A. Fasoli, S.H. Muller, F.M. Poli, P. Ricci, C. Theiler, S. Brunner, A. Diallo, and J. Graves. Experimental Observation of the Blob-Generation Mechanism from Interchange Waves in a Plasma. *Physical Review Letters*, Vol. 100, No. 5, p. 055004, 2008.
- [49] I. Furno, B. Labit, A. Fasoli, F.M. Poli, P. Ricci, C. Theiler, S. Brunner, A. Diallo, J.P. Graves, M. Podesta, and S.H. Muller. Mechanism for blob generation in the TORPEX toroidal plasma. *Physics of Plasmas*, Vol. 15, No. 5, p. 055903, 2008.
- [50] C. Theiler, I. Furno, P. Ricci, A. Fasoli, B. Labit, S.H. Muller, and G. Plyushchev. Cross-Field Motion of Plasma Blobs in an Open Magnetic Field Line Configuration. *Physical review letters*, Vol. 103, No. 6, p. 65001, 2009.
- [51] R. Barni and C. Riccardi. Advection of long lived density blobs in the turbulent state of a simple magnetized torus plasma. *Plasma Physics and Controlled Fusion*, Vol. 51, p. 085010, 2009.
- [52] S.H. Muller, C. Theiler, A. Fasoli, I. Furno, B. Labit, G.R. Tynan, M. Xu, Z. Yan, and J.H. Yu. Studies of blob formation, propagation and transport mechanisms in basic experimental plasmas (TORPEX and CSDX). *Plasma Physics and Controlled Fusion*, Vol. 51, p. 055020, 2009.
- [53] T. Happel, F. Greiner, N. Mahdizadeh, B. Nold, M. Ramisch, and U. Stroth. Generation of Intermittent Turbulent Events at the Transition from Closed to Open Field Lines in a Toroidal Plasma. *Physical review letters*, Vol. 102, No. 25, p. 255001, 2009.
- [54] P. Ricci, C. Theiler, A. Fasoli, I. Furno, B. Labit, S.H. Muller, M. Podesta, and F.M. Poli. Langmuir probe-based observables for plasma-turbulence code validation and application to the TORPEX basic plasma physics

- experiment. *Physics of Plasmas*, Vol. 16, p. 055703, 2009.
- [55] H. Tanaka, N. Ohno, Y. Tsuji, and S. Kajita. 2D Statistical Analysis of Non-Diffusive Transport under Attached and Detached Plasma Conditions of the Linear Divertor Simulator. *Contributions to Plasma Physics*, Vol. 50, No. 3-5, pp. 256–266, 2010.
- [56] N. Sasini, R. Paikaray, L. Dinda, G. Sahoo, J. Ghosh, and A.K. Sanyasi. Density and temperature measurements of pulsed plasma produced inside a curved vacuum chamber. In *Journal of Physics: Conference Series*, Vol. 208, p. 012132. IOP Publishing, 2010.
- [57] Y. Sarazin and Ph. Ghendrih. Intermittent particle transport in two-dimensional edge turbulence. *Physics of Plasmas*, Vol. 5, No. 12, pp. 4214–4228, 1998.
- [58] D.A. D’Ippolito, J.R. Myra, and S.I. Krasheninnikov. Cross-field blob transport in tokamak scrape-off-layer plasmas. *Physics of Plasmas*, Vol. 9, No. 1, pp. 222–233, 2002.
- [59] N. Bian, S. Benkadda, J.-V. Paulsen, and O.E. Garcia. Blobs and front propagation in the scrape-off layer of magnetic confinement devices. *Physics of Plasmas*, Vol. 10, No. 3, pp. 671–676, 2003.
- [60] D.A. D’Ippolito and J.R. Myra. Blob stability and transport in the scrape-off-layer. *Physics of Plasmas*, Vol. 10, No. 10, pp. 4029–4039, 2003.
- [61] G.Q. Yu and S.I. Krasheninnikov. Dynamics of blobs in scrape-off-layer/shadow regions of tokamaks and linear devices. *Physics of Plasmas*, Vol. 10, No. 11, pp. 4413–4418, 2003.
- [62] S.I. Krasheninnikov and A.I. Smolyakov. On neutral wind and blob motion in linear devices. *Physics of Plasmas*, Vol. 10, p. 3020, 2003.
- [63] D.A. D’Ippolito, J.R. Myra, S.I. Krasheninnikov, G.Q. Yu, and A. Yu. Blob Transport in the Tokamak Scrape-off-Layer. *Contributions to Plasma Physics*, Vol. 44, No. 1-3, pp. 205–216, 2004.

-
- [64] D.A. D'Ippolito, J.R. Myra, D.A. Russell, and G.Q. Yu. Rotational stability of plasma blobs. *Physics of Plasmas*, Vol. 11, No. 10, pp. 4603–4609, 2004.
- [65] J.R. Myra, D.A. D'Ippolito, S.I. Krasheninnikov, and G.Q. Yu. Convective transport in the scrape-off-layer by nonthermalized spinning blobs. *Physics of Plasmas*, Vol. 11, No. 9, pp. 4267–4274, 2004.
- [66] J.R. Myra, D.A. D'Ippolito, D.A. Russell, S.I. Krasheninnikov, A.Y. Pigarov, G.Q. Yu, X.Q. Xu, W.M. Nevins, and M. Umansky. Plasma Convection by Blobs in the Scrape-off-Layer. *London, June 28-July*, Vol. 2, , 2004.
- [67] A.Y. Aydemir. Convective transport in the scrape-off layer of tokamaks. *Physics of Plasmas*, Vol. 12, No. 6, p. 062503, 2005.
- [68] O.E. Garcia, N.H. Bian, V. Naulin, A.H. Nielsen, and J.J. Rasmussen. Mechanism and scaling for convection of isolated structures in nonuniformly magnetized plasmas. *Physics of Plasmas*, Vol. 12, No. 9, p. 090701, 2005.
- [69] R.H. Cohen and D.D. Ryutov. Dynamics of an Isolated Blob in the Presence of the X-Point. *Contributions to Plasma Physics*, Vol. 46, No. 7-9, pp. 678–684, 2006.
- [70] D.A. D'Ippolito and J.R. Myra. Thermal transport catastrophe and the tokamak edge density limit. *Physics of Plasmas*, Vol. 13, No. 6, p. 062503, 2006.
- [71] O.E. Garcia, N.H. Bian, and W. Fundamenski. Radial interchange motions of plasma filaments. *Physics of Plasmas*, Vol. 13, No. 8, p. 082309, 2006.
- [72] J.R. Myra, D.A. Russell, and D.A. D'Ippolito. Collisionality and magnetic geometry effects on tokamak edge turbulent transport. I. A two-region model with application to blobs. *Physics of Plasmas*, Vol. 13, No. 11, p.

- 112502, 2006.
- [73] D.D. Ryutov. The dynamics of an isolated plasma filament at the edge of a toroidal device. *Physics of Plasmas*, Vol. 13, No. 12, p. 122307, 2006.
 - [74] G.Q. Yu, S.I. Krasheninnikov, and P.N. Guzdar. Two-dimensional modelling of blob dynamics in tokamak edge plasmas. *Physics of Plasmas*, Vol. 13, No. 4, p. 042508, 2006.
 - [75] S.I. Krasheninnikov and A.I. Smolyakov. Generation of mesoscale convective structures in tokamak edge plasma. *Physics of Plasmas*, Vol. 14, p. 102503, 2007.
 - [76] J.R. Myra. Current carrying blob filaments and edge-localized-mode dynamics. *Physics of Plasmas*, Vol. 14, p. 102314, 2007.
 - [77] D. Jovanovic, P.K. Shukla, and F. Pegoraro. Effects of the parallel electron dynamics and finite ion temperature on the plasma blob propagation in the scrape-off layer. *Physics of Plasmas*, Vol. 15, p. 112305, 2008.
 - [78] D.A. D'Ippolito and J.R. Myra. Effect of discrete coherent structures on plasma-wall interactions in the scrape-off-layer. *Physics of Plasmas*, Vol. 15, p. 082316, 2008.
 - [79] K. Bodi, S.I. Krasheninnikov, and A.I. Smolyakov. Blob dynamics in an inhomogeneous plasma. *Physics of Plasmas*, Vol. 15, p. 102304, 2008.
 - [80] W. Kerner, A. Nassigh, and O. Pogutse. Electromagnetic interchange instability induced by sheath resistivity in the SOL. *Plasma Physics and Controlled Fusion*, Vol. 39, p. 1101, 1997.
 - [81] A.Y. Pigarov, S.I. Krasheninnikov, T.D. Rognlien, M.J. Schaffer, and W.P. West. Tokamak edge plasma simulation including anomalous cross-field convective transport. *Physics of Plasmas*, Vol. 9, No. 4, pp. 1287–1299, 2002.
 - [82] Ph. Ghendrih, Y. Sarazin, G. Attuel, S. Benkadda, P. Beyer, G. Falchetto, C. Figarella, X. Garbet, V. Grandgirard, and M. Ottaviani. Theoreti-

-
- cal analysis of the influence of external biasing on long range turbulent transport in the scrape-off layer. *Nuclear Fusion*, Vol. 43, No. 10, pp. 1013–1022, 2003.
- [83] Y. Sarazin, Ph. Ghendrih, G. Attuel, C. Clement, X. Garbet, V. Grandgirard, M. Ottaviani, S. Benkadda, P. Beyer, N. Bian, et al. Theoretical understanding of turbulent transport in the SOL. *Journal of Nuclear Materials*, Vol. 313, pp. 796–803, 2003.
- [84] N. Bisai, A. Das, S. Deshpande, R. Jha, P. Kaw, A. Sen, and R. Singh. Simulation of plasma transport by coherent structures in scrape-off-layer tokamak plasmas. *Physics of Plasmas*, Vol. 11, No. 8, pp. 4018–4024, 2004.
- [85] O.E. Garcia, V. Naulin, A.H. Nielsen, and J.J. Rasmussen. Computations of Intermittent Transport in Scrape-Off Layer Plasmas. *Phys. Rev. Lett.*, Vol. 92, No. 16, p. 165003, Apr 2004.
- [86] D.A. Russell, D.A. D'Ippolito, J.R. Myra, W.M. Nevins, and X.Q. Xu. Blob Dynamics in 3D [small-caps BOUT] Simulations of Tokamak Edge Turbulence. *Physical Review Letters*, Vol. 93, No. 26, p. 265001, 2004.
- [87] Ph. Ghendrih, Y. Sarazin, G. Attuel, S. Benkadda, P. Beyer, G. Darmet, G. Falchetto, C. Figarella, X. Garbet, V. Grandgirard, et al. Statistical analysis of turbulent front propagation in the scrape-off-layer. *Journal of Nuclear Materials*, Vol. 337, pp. 347–351, 2005.
- [88] Smolyakov A.I., S.I. Krasheninnikov, and O.I. Tolstikhin. Electromagnetic instability of an axially inhomogeneous plasma. *Physics of Plasmas*, Vol. 12, No. 5, p. 052112, 2005.
- [89] N. Bisai, A. Das, S. Deshpande, R. Jha, P. Kaw, A. Sen, and R. Singh. Edge and scrape-off layer tokamak plasma turbulence simulation using two-field fluid model. *Physics of Plasmas*, Vol. 12, No. 7, p. 072520, 2005.

- [90] N. Bisai, A. Das, S. Deshpande, R. Jha, P. Kaw, A. Sen, and R. Singh. Formation of a density blob and its dynamics in the edge and the scrape-off layer of a tokamak plasma. *Physics of Plasmas*, Vol. 12, No. 10, p. 102515, 2005.
- [91] O.E. Garcia, V. Naulin, A.H. Nielsen, and J.J. Rasmussen. Turbulence and intermittent transport at the boundary of magnetized plasmas. *Physics of Plasmas*, Vol. 12, No. 6, p. 062309, 2005.
- [92] J.R. Myra and D.A. D'Ippolito. Edge instability regimes with applications to blob transport and the quasicohherent mode. *Physics of Plasmas*, Vol. 12, No. 9, p. 092511, 2005.
- [93] C.F. Figarella, Ph. Ghendrih, Y. Sarazin, G. Attuel, S. Benkadda, P. Beyer, G. Falchetto, E. Fleurence, X. Garbet, and V. Grandgirard. Control of long range turbulent transport with biasing in the tokamak scrape-off-layer. *Journal of Nuclear Materials*, Vol. 337, pp. 342–346, March 2005.
- [94] J.R. Myra, D.A. D'Ippolito, D.P. Stotler, S.J. Zweben, B.P. LeBlanc, J.E. Menard, R.J. Maqueda, and J. Boedo. Blob birth and transport in the tokamak edge plasma: Analysis of imaging data. *Physics of Plasmas*, Vol. 13, No. 9, p. 092509, 2006.
- [95] O.E. Garcia, V. Naulin, A.H. Nielsen, and J.J. Rasmussen. Turbulence simulations of blob formation and radial propagation in toroidally magnetized plasmas. *Physica Scripta*, Vol. 122, pp. 89–103, 2006.
- [96] V. Naulin, W. Fundamenski, A.H. Nielsen, J.J. Rasmussen, O.E. Garcia, B. Gonçalves, C. Hidalgo, M. Hron, et al. Turbulence Modelling of JET SOL Plasma. *Proceedings of 21st IAEA Fusion Energy Conference*, Vol. TH/P6-22, , 2006.
- [97] O.E. Garcia, J. Horacek, R.A. Pitts, A.H. Nielsen, W. Fundamenski, V. Naulin, and J.J. Rasmussen. Fluctuations and transport in the TCV

-
- scrape-off layer. *Nuclear Fusion*, Vol. 47, No. 7, pp. 667–676, 2007.
- [98] D. Reiser. Impact of large island perturbations on turbulent blob transport in tokamaks. *Physics of Plasmas*, Vol. 14, p. 082314, 2007.
- [99] D.A. Russell, J.R. Myra, and D.A. D’Ippolito. Collisionality and magnetic geometry effects on tokamak edge turbulent transport. II. Many-blob turbulence in the two-region model. *Physics of Plasmas*, Vol. 14, p. 102307, 2007.
- [100] X. Xu, M. Umansky, B. Dudson, and P. Snyder. Boundary Plasma Turbulence Simulations for Tokamaks. *Communications in Computational Physics*, vol. 4, no. 5, July 1, 2008, pp. 949-979, Vol. 4, No. LLNL-JRNL-404082, 2008.
- [101] S.I. Krasheninnikov and A.I. Smolyakov. Dynamics and generation mechanisms of mesoscale structures in tokamak edge plasmas. *Physics of Plasmas*, Vol. 15, No. 5, p. 055909, 2008.
- [102] J.R. Myra, D.A. Russell, and D.A. D’Ippolito. Transport of perpendicular edge momentum by drift-interchange turbulence and blobs. *Physics of Plasmas*, Vol. 15, No. 3, p. 032304, 2008.
- [103] J.R. Myra, D.A. Russell, and D.A. D’Ippolito. Transport of perpendicular edge momentum by drift-interchange turbulence and blobs. *Physics of Plasmas*, Vol. 15, p. 032304, 2008.
- [104] J. Anderson and E. Kim. Nonperturbative models of intermittency in edge turbulence. *Physics of Plasmas*, Vol. 15, p. 122303, 2008.
- [105] S.I. Krasheninnikov, A.Y. Pigarov, T.K. Soboleva, and D.L. Rudakov. Strongly intermittent edge plasma transport: Issues with modeling and interpretation of experimental data. *Physics of Plasmas*, Vol. 16, p. 014501, 2009.
- [106] K. Bodi, A.I. Smolyakov, and S.I. Krasheninnikov. On blob generation mechanisms in tokamak edge plasma. *Journal of Nuclear Materials*, Vol.

- 390, pp. 359–363, 2009.
- [107] G.S. Xu, V. Naulin, W. Fundamenski, J.J. Rasmussen, A.H. Nielsen, and B.N. Wan. Intermittent convective transport carried by propagating electromagnetic filamentary structures in nonuniformly magnetized plasma. *Physics of Plasmas*, Vol. 17, p. 022501, 2010.
- [108] S.-I. Itoh. Structure Formation in Turbulent Plasmas. *Plasma and Fusion Research*, Vol. 4, pp. 38–38, 2009.
- [109] P.H. Diamond, S.-I. Itoh, K. Itoh, and T.S. Hahm. Zonal flows in plasma—a review. *Plasma Physics and Controlled Fusion*, Vol. 47, p. R35, 2005.
- [110] P.H. Diamond, S.-I. Itoh, and K. Itoh. *Modern Plasma Physics: Volume 1, Physical Kinetics of Turbulent Plasmas*. Cambridge Univ Pr, 2010.
- [111] A. Fujisawa. A review of zonal flow experiments. *Nuclear Fusion*, Vol. 49, p. 013001, 2009.
- [112] G.R. Tynan, A. Fujisawa, and G. McKee. A review of experimental drift turbulence studies. *Plasma Physics and Controlled Fusion*, Vol. 51, p. 113001, 2009.
- [113] T. Yamada, S.-I. Itoh, T. Maruta, N. Kasuya, Y. Nagashima, S. Shinohara, K. Terasaka, M. Yagi, S. Inagaki, Y. Kawai, et al. Anatomy of plasma turbulence. *Nature Physics*, Vol. 4, No. 9, pp. 721–725, 2008.
- [114] Y. Nagashima, S.-I. Itoh, S. Shinohara, M. Fukao, A. Fujisawa, K. Terasaka, Y. Kawai, G.R. Tynan, P.H. Diamond, M. Yagi, et al. Observation of the parametric-modulational instability between the drift-wave fluctuation and azimuthally symmetric sheared radial electric field oscillation in a cylindrical laboratory plasma. *Physics of Plasmas*, Vol. 16, p. 020706, 2009.
- [115] H. Arakawa, K. Kamataki, S. Inagaki, T. Maruta, Y. Nagashima, T. Yamada, S. Shinohara, K. Terasaka, S. Sugita, M. Yagi, et al. Observations of abrupt changes in the fluctuation spectrum on LMD-U. *Plasma Physics*

-
- and Controlled Fusion*, Vol. 51, p. 085001, 2009.
- [116] H. Arakawa, S. Inagaki, Y. Nagashima, T. Yamada, K. Kamataki, T. Kobayashi, S. Sugita, M. Yagi, N. Kasuya, A. Fujisawa, et al. Bifurcation of the plasma turbulence on LMD-U. *Plasma Physics and Controlled Fusion*, Vol. 52, p. 105009, 2010.
- [117] K. Itoh. Summary of IAEA Technical Meeting on Theory of Plasma Instabilities. *Nuclear Fusion*, Vol. 50, p. 054001, 2010.
- [118] C. Theiler, I. Furno, P. Ricci, A. Fasoli, B. Labit, and TORPEX basic plasma physics group. Study of plasma filament motion and their active control. In *36th European Physical Society Conference on Plasma Physics*, 6 2009.
- [119] B.B. Kadomtsev. Plasma turbulence. 1965.
- [120] S. Sugita, P. Beyer, G. Fuhr, F. Jaulmes, S. Benkadda, X. Garbet, M. Yagi, S.-I. Itoh, and K. Itoh. 3D Turbulence Simulations of the Impact of Plasma Rotation on Stiffness Level. *Proceedings of European Physical Society 37th Conference on Plasma Physics*, 2010.
- [121] P. Mantica, D. Strintzi, T. Tala, C. Giroud, T. Johnson, H. Leggate, E. Lerche, T. Loarer, A.G. Peeters, A. Salmi, et al. Experimental Study of the Ion Critical-Gradient Length and Stiffness Level and the Impact of Rotation in the JET Tokamak. *Physical review letters*, Vol. 102, No. 17, p. 175002, 2009.
- [122] P. Beyer, S. Benkadda, G. Fuhr-Chaudier, X. Garbet, P. Ghendrih, and Y. Sarazin. Turbulence simulations of transport barrier relaxations in tokamak edge plasmas. *Plasma Physics and Controlled Fusion*, Vol. 49, p. 507, 2007.
- [123] G. Fuhr, P. Beyer, S. Benkadda, and X. Garbet. Evidence from Numerical Simulations of Transport-Barrier Relaxations in Tokamak Edge Plasmas in the Presence of Electromagnetic Fluctuations. *Physical review letters*,

- Vol. 101, No. 19, p. 195001, 2008.
- [124] X. Garbet, P. Mantica, C. Angioni, E. Asp, Y. Baranov, C. Bourdelle, R. Budny, F. Crisanti, G. Cordey, L. Garzotti, et al. Physics of transport in tokamaks. *Plasma Physics and Controlled Fusion*, Vol. 46, p. B557, 2004.
- [125] S.A. Orszag. *Physics of Fluids Supplement II*, Vol. 12, p. 250, 1969.
- [126] S.A. Orszag. *Journal of the Atmospheric Science*, Vol. 27, p. 890, 1970.
- [127] E. Eliassen, B. Machenhauer, and E. Rasmussen. *On a numerical method for integration of the hydrodynamical equations with a spectral representation of the horizontal fields*. Københavns Universitet, Institut for Teoretisk Meteorologi, 1970.

RÉSUMÉ

Il est maintenant accepté expérimentalement que les filaments de plasma alignés sur le champ magnétique, appelés “blobs”, jouent un rôle important dans le transport dans le plasma de bord. Dans cette thèse, les phénomènes fondamentaux du transport dans le plasma de bord sont étudiés en mettant l'accent sur le phénomène de filaments plasma. Dans un premier temps, les mécanismes de propagation de blobs uniques sont envisagés. Puis la génération de blobs par la turbulence de bord est étudiée, et le transport turbulent est discuté en tant que phénomène collectif. Des particularités du transport turbulent, incluant les blobs auto-organisés, sont reliées à un transport de type Bohm (c'est à dire des perturbations avec des corrélations radiales longues, et un coefficient de transport effectif qui suit la dépendance Bohm). De plus, en prolongement de ce travail, un effort initial vers une transposition du transport non-local au plasma de bord est décrite.

ABSTRACT

Recently, it has been accepted that magnetic-field-aligned plasma filaments, referred to as ‘blobs’ play important roles in the transport of Scrape-off Layer (SoL) plasmas. In this thesis, putting an emphasis on the plasma blob phenomenon, we study fundamental processes of SoL transport using numerical simulation. At first, we investigate the propagation mechanisms of single and isolated blobs. Next, we study the generation of blobs from edge turbulence, and discuss the SoL turbulent transport as a collective phenomenon. Features of turbulent transport, which includes the self-organized blobs in SoL, are identified as Bohm-like transport (i.e., the perturbation with long radial correlations and the effective transport coefficient that follows the dependence of Bohm-like transport). Additionally, as an advancement of study, we describe an initial effort to extend the view of nonlocal transport to edge plasmas.