

Study on Plasma Transport Including Impurity in Tokamak Improved Confinement

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Study on Plasma Transport Including Impurity
in Tokamak Improved Confinement

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Chapter 1

Introduction

1.1 Nuclear fusion device

Various methods such as thermal power, nuclear power, hydroelectric power, and new energy generation are used to produce energy for human consumption. However, stability, depletion of energy resources, and global warming caused by increased energy consumption have become major concerns. While nuclear power generation is used in many countries due to its stable energy production, it has its own problems, such as environmental pollution and limited resources. Fusion energy, on the other hand, is expected to be a sustainable energy source as it can produce more energy with less environmental impact than nuclear power and has almost inexhaustible fuel resources. The experiment of the JT-60SA device is beginning in Japan, and the experiment of the International Thermonuclear Experimental Reactor (ITER) is approaching internationally, which shows a step towards the realization of a fusion reactor.

1.1.1 Nuclear fusion reaction

Fission reaction, as shown in Fig. 1.1(a), is a process in which a neutron collides with a fissile nuclide, causing it to split into two nuclei and releasing energy. Uranium fission reactions typically release about 2.4 neutrons. These neutrons then react with more uranium nuclei, causing the number of reactions to increase exponentially. For this reason, nuclear reactors must have a control mechanism of the reactions. In contrast, as shown in Fig. 1.1(b), a fusion reaction occurs when two light nuclei combine to form a heavier one, releasing energy in the process. This energy is then converted into electrical energy, which is called fusion energy. Unlike fission reactions, which are caused by neutral neutrons, fusion reactions occur when positively charged nuclei collide with each other against Coulomb forces. While one way to cause fusion reactions is to collide high-energy charged particle beams, this method requires a lot of input power. Therefore, thermonuclear fusion reactors are expected to maintain the reaction in high-temperature plasma.

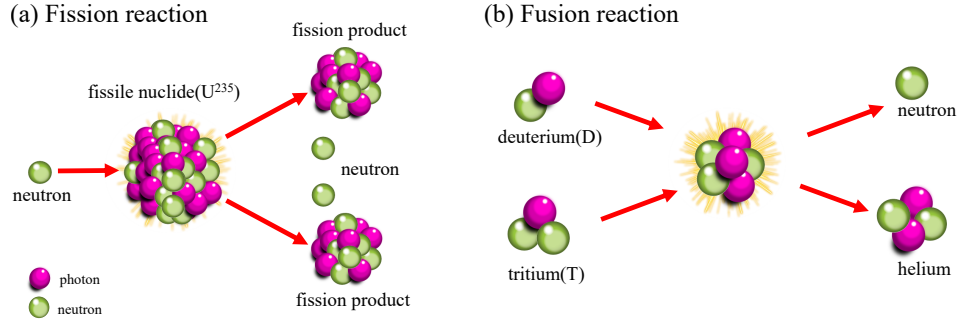
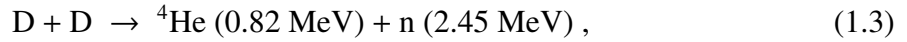
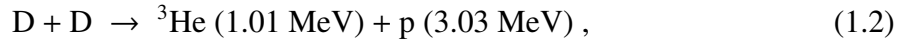
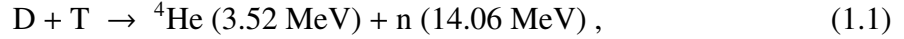


Figure 1.1: Schematic diagram of (a) fission reaction and (b) fusion reaction.

Hydrogen, helium, and lithium are low mass number atoms, which have low bond energy and thus can become stable through fusion reactions. Deuterium (D) and tritium (T) are hydrogen isotopes that can undergo fusion reactions, which can generate power and are currently under study. The following reaction formulas are being considered for fusion reactions in thermonuclear reactors.



where 'n' is the neutron and 'p' is the proton. The ${}^4\text{He}$ produced by the fusion reaction is also known as an α -particle. As shown in Fig. 1.2, the DT reaction occurs more easily at lower temperatures than the DD reaction, and the sustaining conditions of the reaction are more accessible to achieve. Consequently, the DT fusion is now actively being studied. In terms of fuel resources, tritium does not exist in nature and is produced from lithium (Li) in fusion reactors. The tritium loaded in the initial operation phase is considered to be produced by using nuclear reactors [1]. Deuterium is easily obtained from seawater, making it an almost unlimited energy resource. The hydrogen isotopes are used as fuel in fusion power generation, which means that there is no release of carbon dioxide or generation of high-level radioactive waste.

1.1.2 Nuclear fusion device system

Fusion reactors need to have systems to sustain the reaction, extract the energy, and convert it to thermal energy to generate electricity. A schematic diagram of a fusion reactor system is shown in Fig. 1.3. The 'blanket', which is installed around the plasma, receives the neutrons and converts kinetic energy into thermal energy, which is then converted into electrical energy by the turbine through an intermediate heat exchanger. Lots of tritium is required for fusion power systems, estimated at 55.8 kg per year per 1000 MW of fusion power generation [3]. The blanket in the reactor also has the function of producing tritium fuel during operation. This is achieved by loading the lithium in the blanket and using the following nuclear reactions to produce tritium.



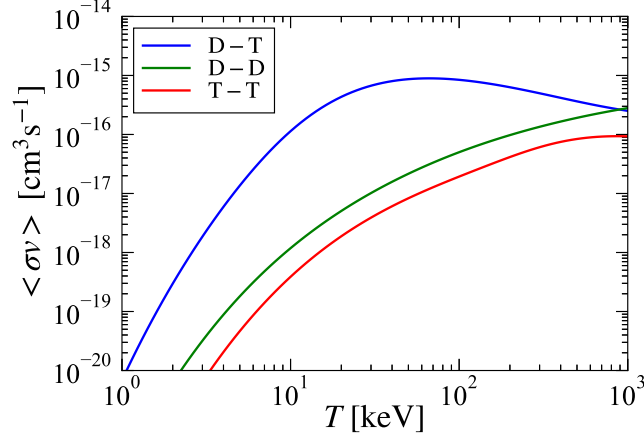


Figure 1.2: Dependence of the fusion reaction rate on the plasma temperature. The fusion reaction rates are calculated using the fitting equation in Ref. 2. Here, σ is the cross-section of the fusion reactions and v is the velocity of the bulk plasma, where the velocity $v = \sqrt{\frac{2T}{m}}$ is assumed the Maxwell distribution function.

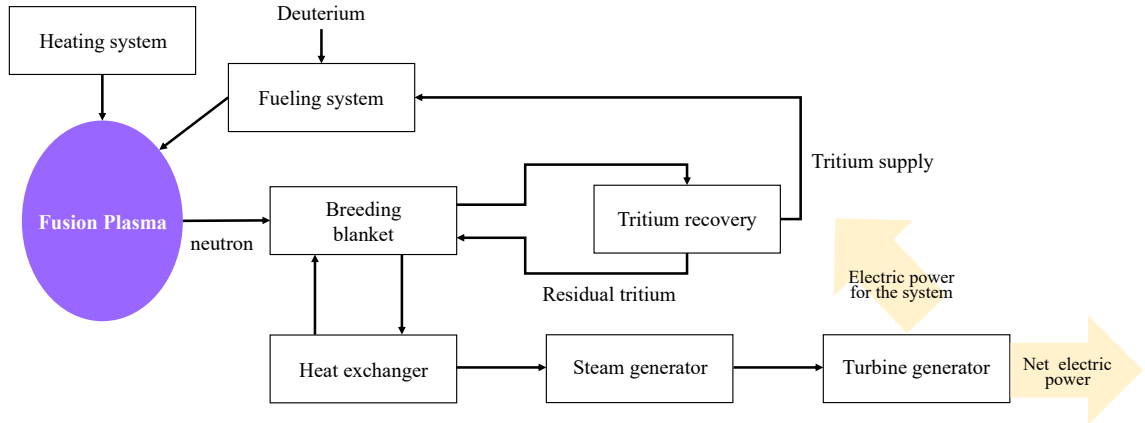


Figure 1.3: Schematic diagram of a fusion power generation system.

The neutrons that react with lithium are those produced by fusion reactions, the reaction of ^7Li , and those produced in neutron multiplication materials such as beryllium (Be) and lead (Pb). The tritium, a newly produced and recovered fuel, is supplied to the core plasma. To test the blanket, a test module called TBM will be installed in a part of the vacuum vessel in the ITER device. The TBM is being developed in collaboration with participating countries in the ITER project. The tritium burnup experiment will prepare a maximum of 4 kg of tritium, and the tritium plant will be prepared for circulating and processing the unburned fuel [4]. Based on the results of the ITER experiments, the tritium breeding system will be verified to achieve a tritium breeding ratio (TBR) of over 1.05 [5], which is necessary for the entire blanket system in steady-state operation in the demonstration power station (DEMO).

In order to sustain high-temperature and high-density fusion plasmas, it is necessary to control actuators for heating, particle supply, and magnetic field generation. The two main heating methods are neutral beam injection (NBI), which involves injecting a neutral beam

of several tens of MW into the plasma, and radio frequency (RF) heating, which is based on the resonant interaction between waves and particles.

The NBI method is a technique used for heating, current drive, and particle supply in plasma. It involves injecting fast-neutral particles into the plasma, which then collide with ions and electrons, resulting in ionization through charge exchange reactions. Positive-ion-based NBI (P-NBI) is a commonly used method in many devices, with beam energy $E_b \sim 100$ keV. However, for $E_b > 100$ keV, the neutralization efficiency significantly decreases. Since the ITER system uses NBI with $E_b = 1$ MeV, a negative-ion-based NBI (N-NBI) is required for high neutralization efficiency. N-NBI systems have been used in JT-60U tokamak and other tokamaks [6]. The NBI technique is useful in feeding particles to the plasma center region, while gas puffing and pellet injection methods are used to feed particles to the plasma periphery.

There are different types of RF heating used in fusion plasma, classified by their frequency. Electron cyclotron resonance heating (ECH) is driven at several tens to several hundreds of GHz, while ion cyclotron resonance heating (ICRH) is driven at several tens to several hundreds of MHz. Lower hybrid resonance heating (LHH) is driven at several hundreds of MHz to several GHz. The heating mechanism in plasma is based on the resonant acceleration of charged particles rotating around magnetic field lines by electric fields at a frequency that approximately matches the rotational frequency of the charged particles. This results in increased collisions with slow electrons. The current drive mechanism is based on the acceleration of electrons in the direction of the RF wave injection, when the RF wave is injected in the opposite direction of the plasma current.

To effectively control plasma, measurement equipment is necessary. Various plasma diagnostic techniques can be used, including methods that utilize properties of electromagnetic, electromagnetic waves, particle, or plasma-neutral particle interactions. These diagnostic methods include Thomson scattering, microwave diagnostics, spectroscopy, interferometry, coil-based diagnostics, neutron cameras, high-speed cameras, and many others. Over time, plasma control researches have been performed using the diagnostic tools to improve plasma performance, including plasma density, temperature, and energy confinement time.

1.1.3 Requirements for fusion reactors

To achieve a commercial fusion reactor, the output energy must exceed the input energy. The following equation gives the fusion power in case of the DT reaction:

$$P_{nf} = n_D n_T \langle \sigma v \rangle E_{nf} , \quad (1.6)$$

where n_D , n_T are the densities of D and T, $E_{nf} = 17.58$ MeV is the fusion energy, and $\langle \sigma v \rangle$ depends on temperature as shown in Fig. 1.2. The power from fission neutrons is $14.06/17.58 = 4/5$ of the total fusion power, while the power from the α -particles is $3.52/17.58 = 1/5$ of the total fusion power. This means that 20% of the total power is used to heat the plasma itself. The energy multiplication factor is $Q = P_{nf}/P_{IN}$, where P_{IN} is the heating power. The condition where $Q = 1$ is called the break-even condition, where the fusion power equals the external heating power. The condition where $Q = \infty$ is called the self-ignition condition, where no external heating is required and a high temperature can be maintained only by the α -heating. The ITER device aims to maintain $Q \geq 10$ for over 400 seconds. To increase the Q value, the fusion triple product $nT\tau_E \geq 10^{21} \text{ m}^{-3}\text{keVs}$ is required

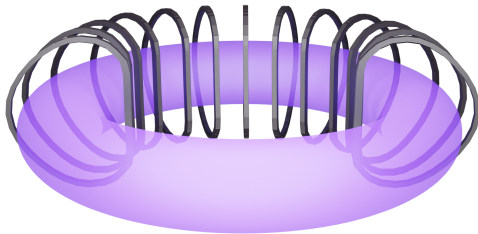
for the core plasma. Good confinement states with small energy losses and particle losses are essential to achieve this condition. Various devices and external control methods have been proposed to maintain high plasma performance for a long duration.

1.2 Tokamak device

1.2.1 Tokamak configuration

Charged particles have the ability to move freely along the direction of magnetic field lines while rotating around them due to the electromagnetic force. This property is utilized to confine high-temperature plasmas to cause nuclear fusion. There are different types of devices that have been developed for plasma confinement, such as mirror-type, tokamak-type, and helical-type. In a mirror-type device, the plasma can be confined in the direction perpendicular to the magnetic field line. However, the plasma tends to escape from the edges in the parallel direction due to the limited length of the linear device, which makes it difficult to maintain high-temperature plasma. On the other hand, as shown in Fig. 1.4, tokamak and helical type devices confine the plasma in a donut shape. The spatial inhomogeneity of the magnetic fields in these devices causes charge separation of particles due to drift motion, which induces an electric field in the perpendicular direction to the torus midplane. The induced electric fields cause charged particles to $E \times B$ drift outside the plasma. In the torus device, the plasma is stationary confined by creating a magnetic field cage by twisting the magnetic field lines. Torus devices are classified into the tokamak and helical types according to how twisted magnetic field lines are created. The helical type uses twisted external coils to create the magnetic field cage, which has excellent plasma stability and is superior for long-time operation. However, since the plasma itself becomes a spiral, there are issues in engineering and particle confinement performance. In the tokamak, the plasma current and coil current create twisting magnetic field lines. Tokamaks have higher plasma confinement performance, making them the most feasible type for a fusion reactor. Some of the typical tokamak devices used for experimental research include JT-60U [7], JET [8], ASDEX-UPGRADE [9], and DIII-D [10]. Further experiments are began in JT-60SA [11] and ITER [12]. DEMO designs are also being developed in Japan [13], EU [14], China [15], and other countries. Representative helical experimental devices are LHD [16] and Wendelstein 7-X [17].

(a) Tokamak



(b) Helical

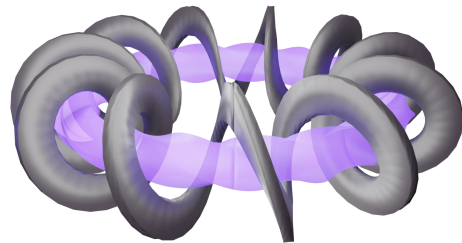


Figure 1.4: Schematic diagram of plasma in (a) tokamak device and (b) helical device.

The Fig. 1.5 shows the cross-section of a tokamak device. The toroidal direction is the circumferential direction of the torus, the poloidal cross-section is the cross-section of the torus and the poloidal direction is the circumferential direction of the poloidal cross-section. The major radial direction is the radial direction from the central axis of the torus to the outside (R) and the minor radial direction is the direction from the center of the poloidal cross-section to the outside of the vacuum vessel (r). A toroidal coil circling the poloidal direction generates a magnetic field in the toroidal direction. By changing the current of the center solenoid coil with time, the magnetic flux through the center of the vacuum vessel changes rapidly, creating an induced electric field in the toroidal direction. This induced electric field drives the plasma current in the toroidal direction and creates a magnetic field in the poloidal direction. The combined poloidal and toroidal magnetic fields form the spiral magnetic field lines that confine the plasma. However, Coulomb collisions between charged particles cause a finite plasma resistivity in tokamak plasmas. The flow of the plasma current generates Joule heat, so the center solenoid coil is also called an ohmic coil. But this plasma current drive method has a limitation in steady-state operation. Therefore, a non-inductive current drive is essential to achieve steady-state fusion plasma. The bootstrap current induced by the plasma itself, NBI or RF wave are the candidates for current drive. Hence, external heating systems are required for current drive, making it almost impossible to achieve $Q = \infty$. To realize an economical fusion reactor, efficient plasma heating is required to increase the bootstrap current.

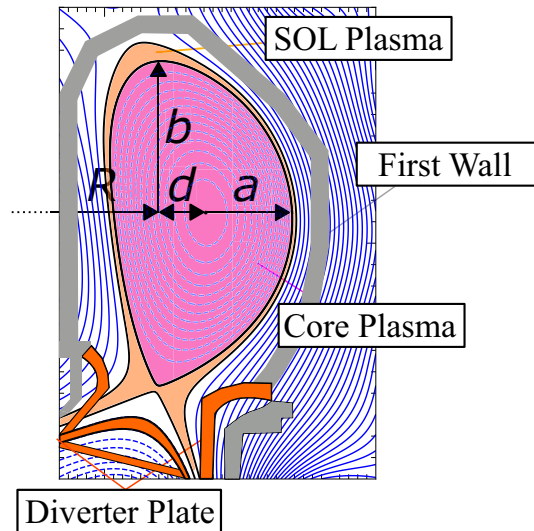


Figure 1.5: Poloidal cross section in tokamak device.

1.2.2 Plasma transport in tokamak core

Plasma transport refers to the transfer of particles, energy, and momentum. Controlling and understanding the physical mechanisms of this transport is crucial for sustaining high-temperature and high-density plasmas. Plasma transport phenomena are classified into classical, neoclassical, and anomalous transport.

1.2.2.1 Classical transport

The movement of plasma fluid across a magnetic field includes diamagnetic flow, $E \times B$ flow, and diffusion, except for the microscopic collective motion of the plasma. The classical transport refers to the transport of charged particles by two-body collisions. The classical diffusion coefficient can be expressed as $D_c = \rho_e^2 \nu_{ei}$, where the Larmor radius $\rho_e = \frac{\sqrt{m_e T_e}}{eB}$, and ν_{ei} is the electron-ion collision frequency. In a uniform magnetic field, charged particles can freely move by winding around the magnetic field lines while Larmor motion. Coulomb collisions cause the position of the guiding center of the Larmor motion to change. These collisions continue towards lower-density regions, leading to the movement of particles across the magnetic fields. The Larmor radius of ions is larger than that of electrons due to their higher mass, and the hydrogen ion diffusion coefficient is approximately 40 times larger. The classical diffusion coefficient is inversely proportional to B^2 , which means that it can be reduced by increasing the magnetic fields.

1.2.2.2 Neoclassical transport

In non-uniform magnetic fields, particles that are in spiral motion around magnetic field lines undergo drift motion. Tokamak plasmas have inhomogeneous magnetic fields, where the magnetic field is stronger in the inner region and weaker in the outer region. As plasma temperature increases, particle collisions decrease, and the particles become more susceptible to magnetic field structures. Variations in the magnetic field B cause a ∇B drift, called a toroidal drift, and the particle orbits in the poloidal surface are shifted from the magnetic flux surface. With high particle velocities in the magnetic field line direction, particles move in toroidal and poloidal directions as untrapped particles. On the other hand, particles with low velocities in the magnetic field direction are trapped in the weak magnetic field region outside the tokamak by the drift motion and become banana orbits. The trapped particle orbits differ from those of the untrapped particles, which are localized in the poloidal direction while rotating slowly in the torus direction. The shift from the orbital plane is about $\Delta \sim \rho q$ for the untrapped particles and $\Delta \sim \rho q / \sqrt{\epsilon}$ for the trapped particles to the minor radial direction, where $\rho = r/R$ is the normalized minor radius, $q = \frac{rB_t}{RB_p}$ is the safety factor, B_t is the toroidal magnetic field, B_p is the poloidal magnetic field, $\epsilon = r/R$ is the aspect ratio. The shift is particularly large for the trapped particles because they are affected by the toroidal drift over a long time. The transport driven by Coulomb two-body collisions, which includes the effects of significant orbital shifts due to the drift and the increase in the collision step size due to trapped particles, is called neoclassical diffusion. Figure 1.6 shows the dependence of the neoclassical diffusion coefficient on the collision frequency in the axisymmetric torus plasma [18]. Depending on the collision frequency, the diffusion coefficient is divided into three regimes: banana regimes, plateau regime, and Pfirsch-Schlüter (PS) regime. In the case of high plasma density and temperature, there are few collisions ($\nu_{ei} < \nu_b = \epsilon^{-\frac{3}{2}} \frac{v_{th,e}}{qR}$), so the particles can move with the banana orbits, and the diffusion coefficient is $D_{BP} \simeq \epsilon^{-\frac{3}{2}} \rho_e^2 \nu_{ei}$. The plateau regime is relatively narrow, the diffusion coefficient is independent of the collision frequency, and the banana and PS regimes are smoothly connected [19]. In the PS regime, where the collisional frequency is sufficiently large ($\nu_{ei} > \frac{v_{th,e}}{qR}$), the electrons cannot move within a banana orbit, and the diffusion coefficient is $D_{PS} \simeq q^2 \rho_e^2 \nu_{ei}$, including the torus effect. Thus, the neoclassical diffusion coefficient is larger than the classical one. Drift motions perpendicular to the

magnetic fields include ∇B drift, $E \times B$ drift, and curvature drift, but they are small compared to the parallel particle velocities. The ∇B drift causes charge separation, and the electrons move in a direction that will cancel it so that the diamagnetic drift is driven in the opposite direction in electrons and ions. The $E \times B$ drift is driven by electric fields E perpendicular to magnetic fields B . The electrons and ions drift inwardly at the inside of the torus and outwardly at the outside with the same magnitude. The curvature drift (helical drift) is caused by the centrifugal force of charged particles due to the curvature of the magnetic fields. The direction is opposite for electrons and ions, and the drift velocity is larger for ions than for electrons. Figure 1.7 shows the collision frequencies of bulk ions for the typical parameters in tokamak devices. In medium-large devices, the bulk ion is in the banana regime for all radial regions except the plasma center region. This means that most deuterium particles are trapped and are not much involved in neoclassical transport. For the parameters of the small device, collisions are in the plateau regime for the broad radial region because the main ion species is hydrogen, and the ion temperature is low. Several models have been developed to calculate neoclassical transport, including NEO [20] and NCLASS [21]. NEO is the first-principle simulation code for solving the full drift equation of motion. NCLASS uses the Hirschman-Sigmar collision operator based on pitch angle scattering, which enables faster calculations, although the accuracy is degraded.

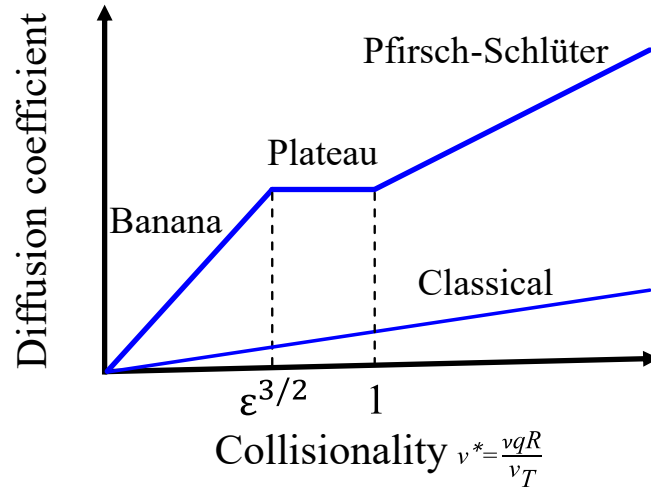


Figure 1.6: Schematic diagram of dependence of the neoclassical diffusion coefficient on collisionality in tokamaks.

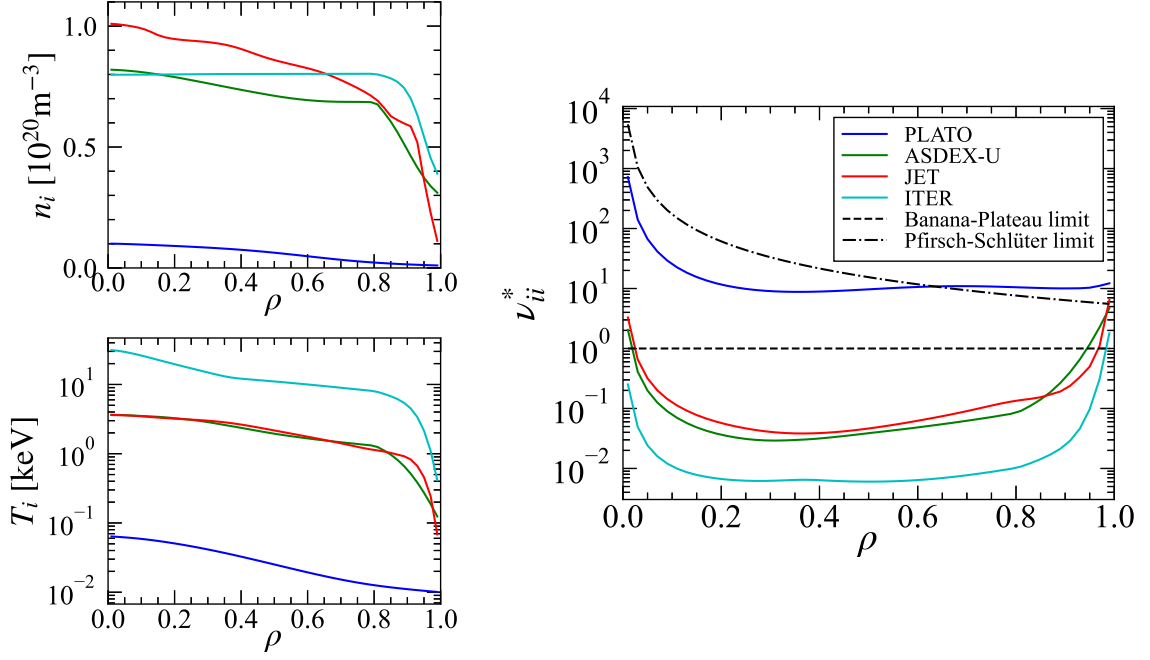


Figure 1.7: Comparison of the radial profiles of ion density n_i , ion temperature T_i , and ion collisionality ν_{ii}^* . Results are shown for the small tokamak PLATO ($n_i = n_H$), the medium tokamak ASDEX-UPGRADE ($n_i = n_D$), and JET ($n_i = n_D$), and the large tokamak ITER ($n_i = n_D + n_T$, $m_i = (m_D + m_T)/2$). The main ion collisionality is defined as $\nu_{ii}^* = \frac{qR\nu_{ii}}{v_{th,i}\epsilon^{\frac{3}{2}}} = \frac{qR\nu_{ii}}{v_{th,i}\epsilon^{\frac{3}{2}}} \frac{1}{2^{\frac{3}{2}}3^{\frac{3}{2}}\sqrt{\pi}} \frac{Z_i^4 e^4 n_i \ln \Lambda}{\epsilon_0^2 \sqrt{m_i T_i^{\frac{3}{2}}}}$, where $\nu_{ii}^* = 1.0$ denotes the banana-plateau limit, and $\nu_{ii}^* = \epsilon^{-\frac{3}{2}}$ denotes the Pfirsch-Schlüter limit.

When there is a gradient in plasma density, electrons move in a banana-shaped orbit and collide with other particles. This collision produces excess momentum in the direction parallel to the magnetic field. The trapped electrons then collide with untrapped electrons, transferring this excess momentum through a viscous force. The momentum of the trapped electrons is lost due to friction with ions, but this is balanced with the excess momentum (viscous force = friction force). This creates a plasma current parallel to the magnetic fields, called the bootstrap current ($j_{BS} \simeq -\epsilon^{\frac{1}{2}} \frac{1}{B_p} \frac{dp}{dr}$). The shape of the pressure gradient affects the bootstrap current, and it is not easily generated at the plasma center. However, it is self-driven in the plasma. Increasing the bootstrap current fraction ($f_{BS} = \frac{I_{BS}}{I_p}$) to the total plasma current (I_p) improves the reactor's economy because the external coils' inductive current can be reduced. This f_{BS} is proportional to the poloidal beta value of the plasma, which requires maintaining a high-temperature and high-density plasma.

1.2.2.3 Turbulent transport

Experimental observations have shown that the transport of plasma is not solely determined by the Coulomb two-body collision process. Anomalous transport, also known as turbulent transport, is a dominant phenomenon in a plasma in a turbulent state [22]. Electromagnetic fluctuations caused by various instabilities in the plasma, such as the drift wave instability induced by temperature and density gradients, and magnetohydrodynamics

(MHD) instability, create microscopic fluctuations that grow and saturate with nonlinear processes. This coupling of velocity fluctuations with density and temperature fluctuations generates radial particle and heat fluxes, leading to anomalous plasma transport. The effect of turbulent transport is much more significant than neoclassical transport, resulting in lower plasma confinement performance than expected. Various theoretical models, such as ion temperature gradient (ITG) modes [23], electron temperature gradient (ETG) modes [24], and trapped electron modes (TEM) [25], have been proposed, but they have yet to explain all experimental results. Depending on the plasma conditions, these modes have different scales and can exist individually or simultaneously in different regions. Fluid instabilities, such as ballooning modes driven by pressure gradients [26] and neoclassical tearing modes [27] driven by self-driven currents, also contribute to anomalous transport. Typical simulation codes for predicting turbulent transport are GKW [28] and GKV [29], based on gyrokinetic theory, and Weiland [30], GLF23 [31], TGLF [32], and QuaLiKiz [33], based on simplified gyrokinetic theory.

1.2.3 Plasma confinement mode

Around 1980, there was a problem: the confinement of plasma was degraded, and high temperatures could not be obtained due to the turbulence generated by extremely unstable plasma in the case of strong plasma heating. This resulted in a plasma with low density and temperature, known as low confinement mode (L-mode). In 1982, the high confinement mode (H-mode), in which the confinement time is improved by a factor of two or more, was observed in the German experimental device ASDEX tokamak [34]. H-mode improved the confinement time by a factor of two or more. This mode suppresses turbulence near the plasma boundary, creating an Edge Transport Barrier (ETB) that prevents particles and energy from leaving the plasma. This increases the density and temperature of the plasma. This breakthrough made it practical to achieve the core plasma conditions required for fusion reactors. The transition from L-mode to H-mode is called the L-H transition. Plasma poloidal rotation or rotation velocity shear due to the radial electric field contribute to the confinement improvement during this transition. H-mode operation is planned as the standard state of the plasma in the ITER device. Based on the results of various experimental devices, it is known that a lower divertor is required to achieve H-mode, and this has been taken into account in the ITER design.

The region that forms the ETB is called the pedestal region. The height and width of the pedestal pressure determine the performance of the entire core plasma. Standard turbulent transport levels are maintained between the top of the pedestal and the magnetic axis. The core plasma temperature tends to be proportional to the pedestal height because the temperature profile is limited by the turbulence driven by temperature gradients such as ITG/TEM (profile stiffness). However, in the pedestal region, an Edge Localized Mode (ELM) causes a relaxation oscillation due to MHD instabilities localized at the plasma boundary, releasing the stored particles and energy. This leads to a decrease in pedestal height, as well as a decrease in density and temperature of the core plasma region, which ultimately leads to a degradation of the plasma confinement performance. Additionally, the considerable energy that is exhausted intermittently can cause damage to the divertor plate.

Databases in many experimental devices has established a scaling law of proportionality for confinement time. The databases are useful for engineering design of fusion devices. The ITER89-P confinement scaling [35] for L-mode tokamak plasmas and the ITER98(y,2)

ELMy H-mode confinement scaling [36, 37] for H-mode tokamak plasmas are as follows:

$$\tau_E^{ITER89-P} = 0.0480 I_p^{0.85} B_t^{0.20} n^{-0.10} P_{IN}^{-0.50} R^{1.50} \kappa^{0.50} \epsilon^{0.30} m_i^{0.50}, \quad (1.7)$$

$$\tau_E^{IPB98y2} = 0.0562 I_p^{0.93} B_t^{0.15} n^{-0.41} P_{IN}^{-0.69} R^{1.97} \kappa^{0.78} \epsilon^{0.58} m_i^{0.19}. \quad (1.8)$$

Here, I_p [MA] is the plasma current, B_t [T] is the toroidal magnetic field, n [10^{20}m^{-3}] is the line averaged plasma density, P_{IN} [MW] is the absorbed power, R [m] is the major radius, κ is the elongation, $\epsilon = r/R$ is the inverse aspect ratio, m_i is the hydrogen isotope mass. The energy confinement time is a measure of the heat transport level of the core plasma. The L-mode thermal diffusion coefficient χ , represented by the ITER89-P scaling, follows the Bohm-type ($\chi_B \propto \frac{T}{eB_t}$). On the other hand, the H-mode thermal diffusion coefficient represented by the ITER98(y,2) scaling is gyro-Bohm-type ($\chi_{GB} \propto \frac{T}{eB_t} \frac{\rho_i}{L_p}$), where ρ_i is the ion gyro radius and L_p is the normalized pressure gradient length. The confinement improvement factor $H_{H98(y,2)} = \frac{\tau_E}{\tau_E^{IPB98y2}}$ is the ratio of the actual confinement time τ_E to the confinement time calculated from the scaling law. This factor is used to indicate how much the confinement is improved and $H_{H98(y,2)} = 1$ is expected in the standard operation of ITER.

In the 1990s, confinement performance was significantly improved in the JT-60 tokamak due to a significant reduction in plasma turbulent transport. In this experiments, it was observed for the first time that the formation of an internal transport barrier (ITB), which is a steep temperature gradient in the confinement region [38]. Plasmas with ITBs were seen in discharges called high- β mode [38–40] and negative shear mode [41–43], and studies have been conducted to determine the mechanisms that determine the position, width, and strength of ITB [44–47]. Several physical mechanisms have been considered for the ITB formation, such as the velocity shear of poloidal rotation due to the radial electric fields, the reversed magnetic shear due to the current drive, the Shafranov-shift due to the pressure gradient, and the toroidal rotation velocity shear. Techniques for ITB formation have been confirmed using intense external heating and rapid current ramp-up [48], as well as pellet injection [49, 50]. The ITB formation mechanism has the property that it is further enhanced due to the increase in pressure gradient, self-driven current, and toroidal rotation velocity, which contribute to the enhancement and maintenance of the ITB as a positive feedback loop. However, the plasma can become unstable if the confinement becomes excessively enhanced and the pressure gradient becomes excessively steep [51–53]. In addition, the plasma can collapse due to the accumulation of impurities inside the internal transport barrier. Therefore, optimization of the ITB plasma is necessary for future advanced operation scenarios.

The plasma current density in L-mode and H-mode plasmas usually peaks at the magnetic axis. However, in plasmas with ITBs, the current density profile is either flat or has a peak at the off-magnetic axis. Bootstrap currents are driven by pressure gradients, so improved confinement mode plasmas with ITBs and ETBs have a higher bootstrap current fraction. The safety factor profile indicates the peaking of the plasma current profile. The confinement modes are characterized by the safety factor profile, as shown in Fig. 1.8. In L-mode and H-mode plasmas, the safety factor increases from the plasma center to the plasma periphery, and the profile of magnetic shear $s = \frac{r}{q} \frac{dq}{dr}$, which represent the rate of change in the radial direction of the safety factor q , have a positive value throughout the radial direction. On the other hand, in high- β mode and negative shear mode, the safety factor profiles have a minimum value except at the magnetic axis, and a region where $s < 0$ is formed. ITBs are easily formed in the $s \leq 0$ region, which is reversed from the monotonic q profile, so

the negative shear mode is also known as the reversed magnetic shear mode. The plasma density and temperature inside the ITB can significantly increase compared to the L-mode and H-mode plasma. The NBI or RF heating with plasma current profile control is an effective method for controlling the ITB formation position.

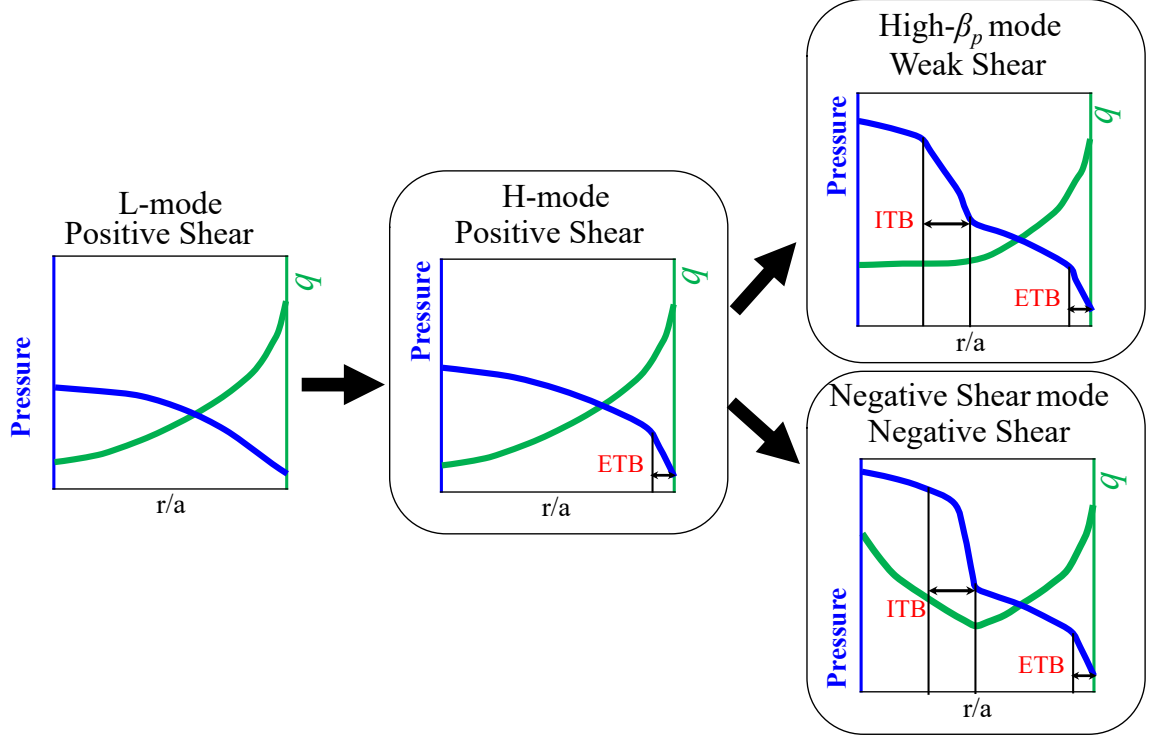


Figure 1.8: Schematic diagram of plasma confinement modes.

Different operation scenarios have been proposed for ITER, including the baseline (inductive) scenario, the hybrid scenario, and the advanced steady-state (fully non-inductive) scenario [54]. The baseline scenario aims to maintain the ELMy H-mode plasma, with a goal of $Q = 10$ for 400 seconds, using the inductive method to start up the plasma current $I_p = 15$ MA. The hybrid scenario aims to maintain $Q = 5 - 10$ for 1000 seconds, characterized by improved confinement compared to the H-mode plasma operation scenario. The plasma is ramped up without causing a sawtooth crash ($q > 1$) at the beginning of the flat-top phase, but the plasma density at the flat-top phase is decreased compared to the baseline scenario. The engineering capabilities, such as the long pulse operation at high neutron fluence, will be demonstrated for future reactor operations. For that purpose, the plasma current decreases compared to the baseline scenario, with $I_p = 12.5$ MA. The plasma current is driven by a combination of induced and non-induced currents to control the current profile. In the steady-state scenario, $Q > 5$ is maintained in steady state, and R&D is carried out to establish a concept of an economically superior demonstration plant for power generation. The plasma current is decreased to $I_p = 9$ MA, which requires using the NBI and RF, and the bootstrap current needs to be increased. The formation of ITBs is essential to obtaining high f_{BS} and high fusion power. Increasing the radiation fraction in the plasma periphery to reduce the equipment material heat load is also a requirement in the steady-state operation scenario. Thus, the steady-state operation is an advanced operation scenario that achieves the plasma conditions required for a plant to demonstrate power generation.

1.3 Impurities in tokamak core plasmas

1.3.1 Impact of impurities

Radiation loss caused by impurities produced by the vessel wall material is one of the factors that reduces the performance of the core plasma. In the tokamak magnetic field configuration, external coils generate magnetic field lines that circle the core to avoid material damage from physical contact between the vacuum vessel and the plasma. These magnetic field lines create a divertor region, which allows ions leaving the core plasma to be transported along the magnetic field lines to the divertor region. This concentration of hot particles in one region results in less damage to device materials and less impurity contamination. However, the divertor plates are exposed to very high heat loads and particle fluxes, which creates a tradeoff between the damage to the wall and the damage to the divertor plate. When the energy of the particles hitting the material exceeds the threshold energy to move the surface atoms, then the surface atoms are sputtered away, causing physical sputtering. As a result, as shown in Fig. 1.9, damaged material mixes with the core plasma as impurity ions. Impurities ejected from the plasma collide with the walls, and some of them sputter again or cause reflection or implantation. Additionally, the divertor plates need cooling, which is achieved by injecting rare gases such as Ne, Ar, and Xe into the divertor region. However, the rare gas particles may mix into the core plasma as impurity ions. Although fuel dilution by impurity is another problem, radiation loss is a serious issue because it significantly degrades plasma performance [55–57]. Therefore, research is being carried out to develop materials with higher strength and methods to eject impurities from the SOL region to the divertor region.

In early tokamak studies, the main impurity species were oxygen, which is absorbed on the walls, and carbon, which is used as the device material. Oxygen can be reduced by cleaning the vacuum vessel wall, and carbon has low radiation loss power. However, carbon materials have disadvantages such as high depletion due to plasma, unfavorable cooling performance, and high thermal conductivity degradation due to neutron irradiation. On the other hand, tungsten (W), which is planned to be used as a wall material for ITER and other experimental devices, has high heat load resistance, low erosion yield, and low deuterium retention [58]. Tungsten, however, causes significant radiation losses and can significantly affect plasma performance, as observed in various experimental devices [59–61]. Even a small amount of tungsten contamination can be severe, as shown in Figure 1.10, which illustrates the dependence of confinement performance on impurity concentration. The acceptable concentration of carbon is $10^{-2} - 10^{-1}$, whereas the acceptable concentration of tungsten is $10^{-4} - 10^{-3}$. Establishing a method to control the impurity accumulation in the core plasma is one of the most important topics. Plasma heating using RF has been experimentally confirmed to be effective in controlling impurities in the core plasma [62–65].

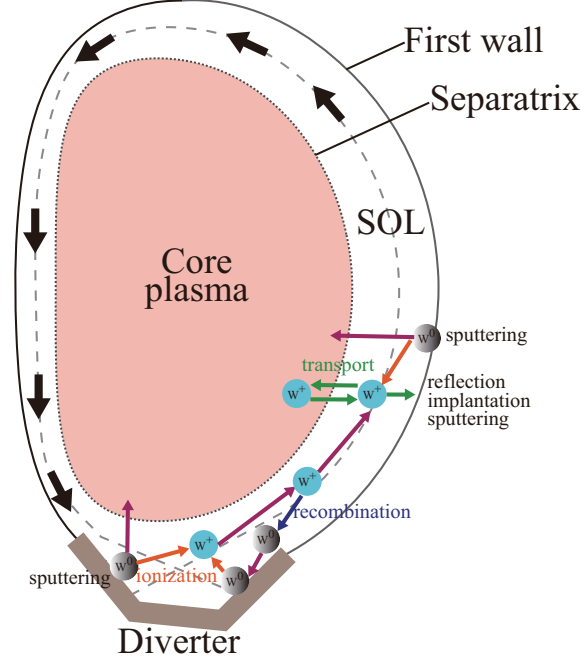


Figure 1.9: Schematic diagram of the impurity contamination processes.

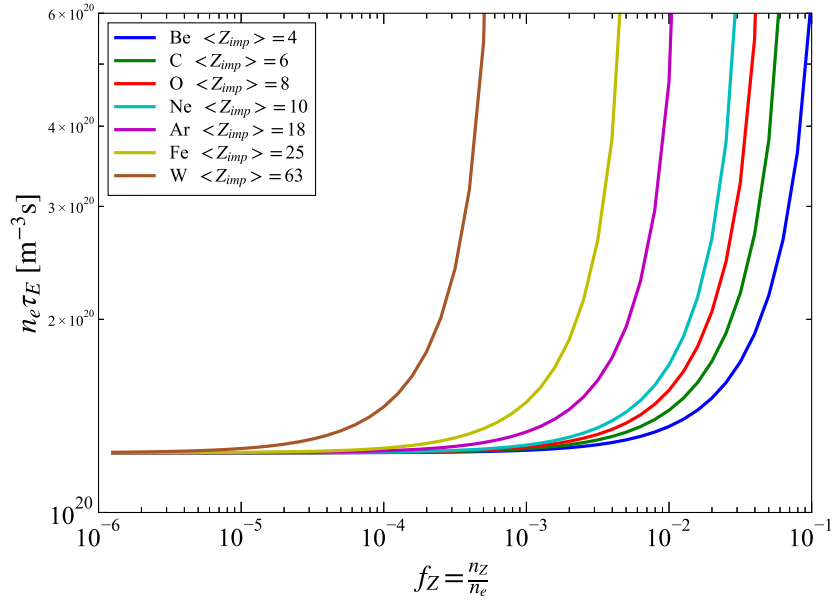


Figure 1.10: 0D analysis of the dependence of $n\tau_E$ on impurity concentration $f_Z = n_Z/n_e$.

The formula of $n_e\tau_E = \frac{\frac{3}{2}T\left(\frac{2}{f_i} + f_Z - Z_Z^{eff} f_Z - \frac{f_{He}}{f_i}\right)}{\frac{AE_{nf}(\sigma v)f_i}{4} - (L_i + f_{He}L_{He} + f_Z L_Z)}$ is used in Ref. 66, and the plasma parameters of $Q = 17.5$, $T_e = T_i = T_Z = 15$ keV, $f_{He} = 0.07$ are used JA-DEMO from Ref. 67. Here, the Q -value for the beam is used $P_{BF}/P_B = 1.0$ ($A \sim 0.27$). The cooling factor L_Z and effective charge number of impurity Z_Z^{eff} are evaluated with Ref. 68. The figure shows the minimum value of $n_e\tau_E$ required for ignition. In the case of tungsten, a high $n\tau_E$ must be achieved for $f_Z > 10^{-4}$ to achieve $Q = 17.5$ and $T = 15$ keV.

1.3.2 Impurity transport in tokamak core plasma

Impurity transport in plasma is governed by both collisional and turbulent transport mechanisms, which are influenced by the charge (Z) and mass (M) of the impurity ion as well as the main plasma conditions. Collisional transport is particularly important as it can cause central peaking in the density profile of high- Z impurities. Impurity collisional transport can be explained by neoclassical transport theory as well as in bulk plasmas. As shown in Fig. 1.11, for the ion temperature $T_i > 1$ keV and the bulk ion density $n_i > 10^{19} \text{ m}^{-3}$ with medium tokamak plasma parameters, the collision of W is in the plateau regime, unlike other particle species. High- Z impurities such as W tend to be in the plateau regime, although W collisions can be in the Pfirsch-Schlüter regime due to the decreased temperature near the plasma boundary. The neoclassical transport dominates in the plasma center where the contribution of turbulent transport is not significant. The flow driven by neoclassical transport is the key mechanism for the accumulation of impurities in the plasma center. The $E \times B$ drift and the diamagnetic drift of bulk ions usually cancel out because they have similar magnitudes. However, since the diamagnetic drift velocity is inversely proportional to the charge number, the poloidal flow velocities of the bulk and impurity ions are different, resulting in a frictional force in the poloidal direction. This causes the flow to be driven towards the plasma center where the bulk ion density is increased, while the thermal force drives the flow towards the outside of the plasma where the temperature is decreased. The balance between the flow due to the density gradient of the main ions and the flow due to the temperature gradient of main ions determines the net radial transport. Heavy impurities such as W, which are high- M and high- Z species, are affected by strong inertial force and electrostatic force, resulting in a distribution with a strong poloidal asymmetry. Through collisions with the toroidally rotating bulk plasma, the impurities also rotate in the toroidal direction. It is experimentally and theoretically known that the impurity localizes at the low field side (LFS) due to enhanced neoclassical transport by the centrifugal effect (CF) [69–73]. In the presence of particles with temperature anisotropy caused by the ICRH, a poloidal electric field localized at the high field side (HFS) can be formed, decreasing impurity accumulation [74].

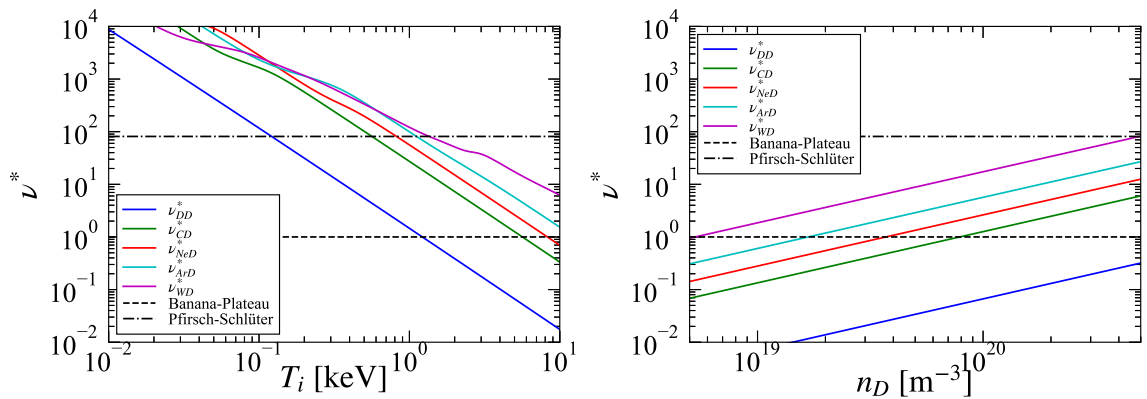


Figure 1.11: Dependence of collisionality between main ion (D) and various particle species on (a) ion temperature and (b) ion density. Here, the plasma parameters are $q = 1$, $R = 3.0$ m, and $r = 0.16$ m. In figure (a), the constant bulk ion density $n_D = 10^{20} \text{ m}^{-3}$ is used, and in figure (b), the constant ion temperature $T_i = T_Z = 5$ keV is used. The banana-plateau limit and the Pfirsch-Schlüter limit are same as Fig. 1.7.

On the other hand, the turbulent transport of the impurity plays an important role near the plasma boundary, especially for low- Z impurity species. In fact, experiments have confirmed that turbulent transport exceeds the prediction of neoclassical transport [75–77]. The main drivers of turbulent transport are the ITG mode and TEM instabilities. At the plasma periphery, turbulent diffusion is expected to be significantly larger than neoclassical diffusion. Fluid theory can describe turbulent convective transport in the impurity trace limit. It is divided into three types of thermo-diffusion pinch, parallel compression pinch, and curvature pinch [78]. The directions of each impurity turbulent pinch have been reviewed according to the dominant instability [79]. The thermo-diffusion pinch has an outward direction driven by the ITG mode and an inward direction driven by the TEM, and its magnitude is proportional to $1/Z$. The parallel compression pinch has an inward direction driven by the ITG mode and an outward direction driven by the TEM, and its magnitude is proportional to Z/A . The curvature pinch has the same direction in both ITG mode and TEM, and its magnitude is independent of Z . It is important to note that the magnetic field gradient is proportional to the magnetic shear. Note that the magnetic field gradient is proportional to the magnetic shear, so the curvature pinch has an inward direction with the positive magnetic shear and an outward direction with the negative shear. Several experiments and theories have investigated the detailed turbulent transport and its dependence on Z [78, 80, 81]. Researchers have also analyzed the impact of toroidal rotation on impurity turbulent transport. It has been reported that rotational shear reduces the peaking of low- Z impurities [82]. On the other hand, impurity transport is also sensitive to MHD events, such as sawtooth oscillation [83–85] and saturated (1,1) mode [74], which can decrease the impurity accumulation.

1.4 Integrated transport modeling

The quantitative prediction of plasma behavior is crucial for the successful development of fusion reactors. To achieve a comprehensive understanding of the physical phenomena that occur in these reactors, it is important to consider various factors, including both physical and engineering problems. Simulation codes that integrate various physical and engineering models provide an analytical tool for experiments and help in understanding plasma phenomena and predicting their behavior. Many integrated codes such as TASK [86], TOPICS [87], TOTAL [88], JINTRAC [89], ASTRA [90], CRONOS [91], and TRANSP [92] are being developed worldwide. A benchmark of integrated transport simulation codes has been performed to confirm their reliability [93]. To facilitate the development of entire integrated codes, it is necessary to create a common framework that defines the data exchange interface between modules. This framework enables mutual use of modules and comparative verification among multiple models and experiments. The ITER organization has developed a data interface called IMAS (Integrated Modelling and Analysis Suite) [94], which is internationally recognized. In Japan, BPSD (Burning Plasma Simulation Data) [95] is being developed based on the BPSI (Burning Plasma Simulation Initiative) [96], and work is underway to link it with IMAS.

1.5 Purpose of this study

The aim of this study is to clarify the impurity accumulation mechanism in improved confinement mode plasmas and identify effective physical mechanisms to control the impurity accumulation. Although improved confinement mode is required to maintain a steady-state burning plasma, the impurity accumulation can occur in the core plasma. To control the accumulation of tungsten in the core plasma, it is essential to understand the mechanism that causes the accumulation. In high-temperature plasmas, the transport process and power balance of multiple ion species need to be considered. The transport process includes collisional and turbulent transport processes. The power balance includes the interaction between the plasma main components and impurities with different charge numbers due to radiation losses. In these processes, to predict the behavior of the plasma, including impurities, the profile formation mechanism must be investigated in a self-consistent manner. Although first-principle models are effective for quantitative analyses to understand the transport mechanisms, they are not realistic for describing the entire plasma discharge due to their huge computational costs. Integrated modeling simulations are useful to investigate the evolution of the core plasma and are also used for experimental analysis and predictions for developing devices. Reliable simulation codes are needed for the development of fusion reactors to predict plasma phenomena, including impurities. Establishing a physical model with fast and high accuracy calculations is an important issue. Currently, a plasma transport simulation scheme including impurity transport is being developed in the integrated code TASK. Many analyses on plasma transport have been performed using previous TASK codes. In this study, the impurity transport analysis code is incorporated into the TASK code, and the simulation scheme is validated.

On the other hand, impurity transport in plasma is influenced by both collisional (neoclassical) and turbulent transport, and it depends on the plasma state, the charge numbers (Z), and the mass. In improved confinement plasmas, where turbulent transport is suppressed, neoclassical transport plays a key role in describing the behavior of high- Z impurities. The effect of large poloidal asymmetry, due to toroidal rotation, is known to be important in impurity accumulation.

The purpose of this study is to develop an integrated simulation code that can describe the entire plasma discharge, including the spatial direction. A neoclassical transport model considering the effect of toroidal rotation is introduced to enable fast calculations. By introducing poloidal asymmetry due to the centrifugal force of impurity ions, as a geometric coefficient, it is demonstrated that the experimental tungsten distribution can be reproduced. This study also investigates the formation of impurity distribution in improved confinement plasmas with an internal transport barrier, to clarify the impurity accumulation mechanism. By simulating both main plasma and impurity transport simultaneously, their interaction and variations of impurity transport in response to plasma evolution are clarified. Finally, an actuator to suppress impurity accumulation is identified and an optimal improved confinement plasma configuration is presented.

1.6 Summary of the chapter

Fusion energy is expected to be a next-generation energy source because of its advantages, such as easier availability of fuel resources, high safety, and less environmental impact. However, there are some unsolved issues that need to be addressed to maintain continuous

plasma burning and make nuclear fusion reactors a reality. One of these issues is turbulent transport, which makes plasma confinement worse. Improved confinement mode operations, which suppress turbulent transport, are effective in achieving high plasma performance in tokamaks. The improved confinement is crucial for the steady-state operation of fusion reactors. Additionally, another one of the critical issues is the impurity accumulation in the core plasma emerged from the vessel wall materials when the high-temperature plasma interacts with them. In the current fusion device design, tungsten is planned to be used as the wall material, but it can cause strong radiation losses in the core regions and degrade plasma performance. One of the improved confinement mode plasmas forms, a characteristic profile called internal transport barrier, which increases the density and temperature of the plasma core but can also possibly lead to increased impurity accumulation. Therefore, keeping the tungsten contamination as low as possible and developing methods to control impurity density are necessary. To control tungsten accumulation, it is essential to understand the mechanism that leads to it. Since the behavior of impurities in the core region is determined by transport processes that are strongly dependent on the background plasma, an approach using an integrated code that deals with multiple physical processes is effective. This study aims to perform plasma transport simulations, including impurities, in the improved confinement mode plasma using the integrated code TASK and clarify the impurity accumulation mechanism and its reduction. The structure of this document is as follows.

In Chapter 2, the integrated code TASK, and the module structure consisting of the MHD equilibrium analysis modules, the plasma transport analysis module, and the impurity transport analysis module are described. Here, the impurity neoclassical transport models including the toroidal rotation effect are presented.

In Chapter 3, the transport simulation scheme is applied to the L-mode plasma in a laboratory tokamak. Here, the plasma performance and impurity effects are predicted for the experimental tokamak device PLATO.

In Chapter 4, the transport simulation scheme including the impurity is applied to the improved confinement mode, the high- β_p mode and the negative magnetic shear mode. The formation mechanism of the internal transport barriers in the plasma profile is clarified. We also investigate the mechanism of the formation of the tungsten profile as a result of the formation of the internal transport barrier.

In Chapter 5, it is shown that the tungsten profile obtained by simulation reproduces the experimental results using the impurity transport model including the effect of toroidal rotation. The effect of toroidal rotation on impurity transport and its mechanism in improved confinement mode plasmas are clarified. Then, the transport mechanism that reduces impurity accumulation in improved confinement mode plasmas is investigated.

In Chapter 6, the results of this study are summarized.

Chapter 2

Simulation model

2.1 TASK code

Integrated code TASK (Transport Analyzing System for tokamaK) is a modular code with various functions such as transport analysis, equilibrium analysis, wave heating analysis, etc., as shown in Fig. 2.1. Although each module can operate by itself, it is possible to realize self-consistent simulations by exchanging data between the modules through the data interface BPSD. The TASK code is open access and source code is available from a git server.

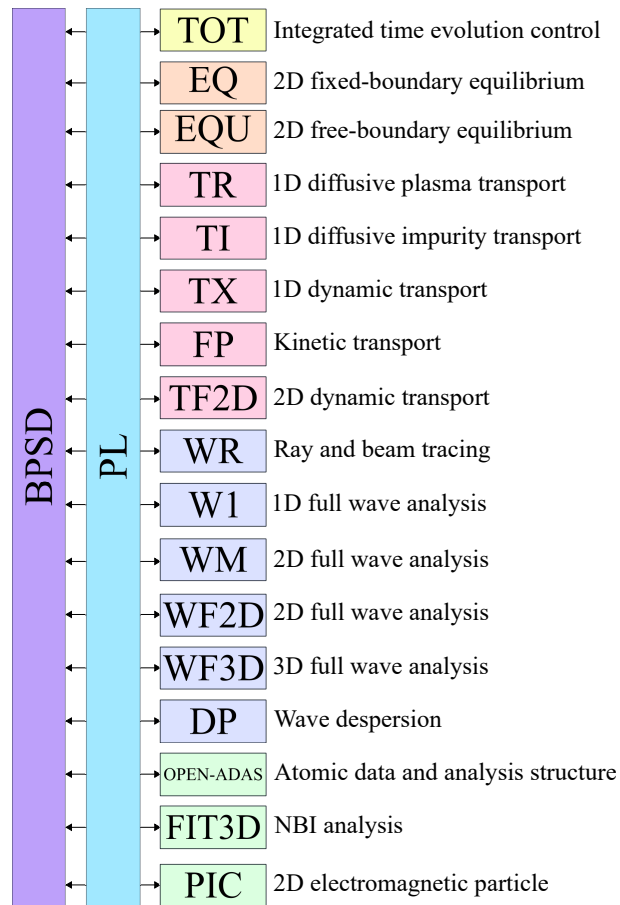


Figure 2.1: Modular structure of the TASK code.

2.2 Transport simulation scheme of core plasma including impurity

A 1.5D plasma transport simulation scheme including impurity transport is being developed in the TASK code. The scheme is composed of the modules shown in Fig. 2.2. The TASK/TR code is used for plasma main component transport, the TASK/EQ code and TASK/EQU code are used for two-dimensional MHD equilibrium, the TASK/TI code is used for impurity transport, and the TASK/WR code is used for EC ray tracing. The OPEN-ADAS database is used to evaluate the atomic processes of impurities. The TASK/TOT code controls these modules and manages their time evolution.

From the TASK/TR code to the TASK/TI code, the radial profiles of the plasma density (n_e, n_i), the plasma temperature (T_e, T_i), the fast ion density (n_{fast}), and the fast ion energy (T_{fast}) are transferred, and conversely, from the TASK/TI code to the TASK/TR code, the radial profiles of the impurity density (n_Z), the effective charge number (Z_Z^{eff}), and the radiation loss power (P_{rad}) are transferred. At present, only impurity particle transport is calculated in the TASK/TI code, and the electron density is determined as $n_e = n_i + Z_Z^{eff} n_Z$ to satisfy the quasi-neutrality condition in the TASK/TR code. Quantities regarding the plasma shape, the metric, and the magnetic-surface averages are obtained from the equilibrium calculations and are used in the transport calculations. On the other hand, in the equilibrium calculation, the plasma pressure profile obtained from the transport calculation is used as an initial condition, and the poloidal current profile is derived from the safety factor q . In the TASK/WR code, the radial absorption profile of ECH is calculated using the equilibrium magnetic field configuration data and the plasma profile data. ECH data is used in the plasma transport code. Using the OPEN-ADAS, the ionization, the recombination, and the radiation loss coefficients of impurities are evaluated and used in the TASK/TI code.

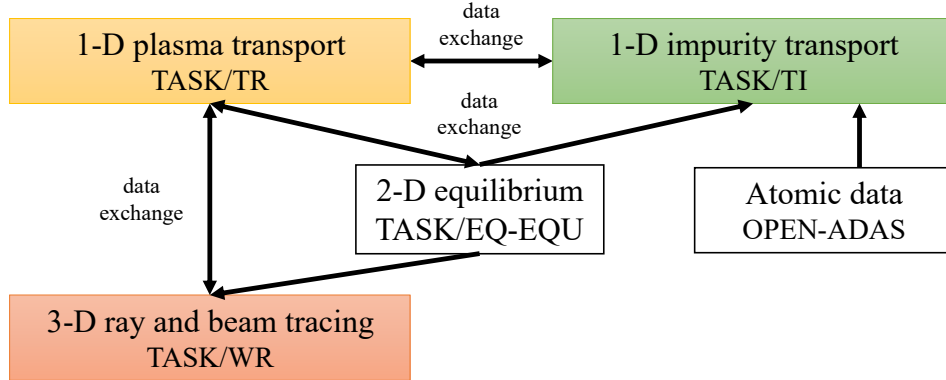


Figure 2.2: Schematic diagram of the plasma transport simulation scheme including impurity in the TASK code.

2.3 MHD equilibrium calculation

The TASK code includes EQ and EQU modules for equilibrium analysis. Both of them are codes to obtain the magnetic configuration in the plasma equilibrium state. The magnetic field configuration of the tokamak can be described by a two-dimensional plasma

cross-section. The macroscopic behavior of the electromagnetic fluid in equilibrium can be described by the following ideal MHD equations:

$$\nabla \cdot \rho_m \mathbf{v} = 0 , \quad (2.1)$$

$$\mathbf{j} \times \mathbf{B} = \nabla p , \quad (2.2)$$

$$\mathbf{E} + \mathbf{v} \times \mathbf{B} = 0 , \quad (2.3)$$

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{j} , \quad (2.4)$$

$$\nabla \cdot \mathbf{B} = 0 . \quad (2.5)$$

Equation 2.1 is the continuity equation, Eq. 2.2 is the momentum balance equation, Eq. 2.3 is the Ohm's law, and Eqs. 2.4 and 2.5 are the Maxwell's equations. Here, ρ_m is the mass density, $\mathbf{v}$ is the velocity, $\mathbf{j}$ is the current density, p is the plasma pressure, $\mathbf{B}$ is the magnetic field, $\mathbf{E}$ is the electric field, and μ_0 is the permeability. Equation 2.2 is the equation of the plasma equilibrium assuming electrical quasi-neutrality conditions. These equations of macroscopic equilibrium can be summarized into a single differential equation called the Grad-shafranov equation using the flux function. In tokamak plasma, surfaces formed by magnetic field lines are nested and are called the magnetic flux surfaces labeled by the flux function ψ . Since charged particles can move freely along the magnetic field lines, the plasma pressure on the magnetic surfaces is constant, so it is also called a constant-pressure surface. The vector perpendicular to the magnetic surface is orthogonal to the magnetic field, so $\mathbf{B} \cdot \nabla \psi = 0$. Furthermore, the plasma current is $\mathbf{j} = \nabla \times \mathbf{B}$, so $\mathbf{j} \cdot \nabla \psi = 0$. With these relations, the Grad-Shafranov equation is given as

$$\nabla \frac{1}{R^2} \nabla \psi = -4\pi^2 \mu_0 \frac{dp}{d\psi} - \frac{I}{R^2} \frac{dI}{d\psi} . \quad (2.6)$$

Here, ψ , p , I , R are the poloidal flux function, the plasma pressure, the poloidal current, and the major radius, respectively. By solving this Grad-Shafranov equation, the plasma equilibrium can be obtained. It is necessary to solve the transport equation to determine the profiles of p and I on the right-hand side with respect to ψ . When EQ or EQU are operated independently, equilibrium can be obtained by assuming an appropriate functional form. In EQU, the Grad-Shafranov equation is solved by considering the magnetic field generated by the external coil current, which can control the plasma position. In the free-boundary equilibrium calculation, ψ includes contributions both from the plasma current ψ_p and the external vertical coil current ψ_{coil}^i as $\psi = \psi_p + \sum_i \psi_{coil}^i$, where i denotes the index number of the coils. The magnetic flux surface quantities of the macroscopic density and temperature of the plasma can be determined by using the following magnetic flux surface averages:

$$\langle X \rangle = \frac{\int \frac{X}{B_p} dl_p}{\int \frac{1}{B_p} dl_p} , \quad (2.7)$$

where, X is the plasma quantities, B_p is the poloidal magnetic field and dl_p is the line element along the poloidal magnetic field.

2.4 Transport calculation

2.4.1 Transport equations

Plasma transport

The macroscopic time evolution of plasma in the TASK code can be obtained by using the TASK/TR code. The continuity equation and the energy conservation equation are derived by velocity integral of the equations of motion in the three-dimensional real space and the three-dimensional velocity space. By averaging the particle and energy equations over the magnetic surface, the time evolution of density and pressure can be described. The following particle transport equation, heat transport equation, and magnetic diffusion equation are solved for the plasma species s .

$$\frac{1}{V'} \frac{\partial}{\partial t} (n_s V') = -\frac{1}{V'} \frac{\partial}{\partial \rho} \left(V' \langle |\nabla \rho| \rangle n_s V_s - V' \langle |\nabla \rho|^2 \rangle D_s \frac{\partial n_s}{\partial \rho} \right) + S_s, \quad (2.8)$$

$$\begin{aligned} \frac{1}{V'^{\frac{5}{3}}} \frac{\partial}{\partial t} \left(\frac{3}{2} n_s T_s V'^{\frac{5}{3}} \right) = & -\frac{1}{V'} \frac{\partial}{\partial \rho} \left(V' \langle |\nabla \rho| \rangle \frac{3}{2} n_s T_s V_{E_s} - V' \langle |\nabla \rho|^2 \rangle \chi_s \frac{\partial n_s T_s}{\partial \rho} \right. \\ & \left. - \langle |\nabla \rho|^2 \rangle \left(\frac{3}{2} D_s - \chi_s \right) T_s \frac{\partial n_s}{\partial \rho} \right) + P_s, \end{aligned} \quad (2.9)$$

$$\frac{\partial}{\partial \rho} \left(\frac{\partial \psi}{\partial \rho} \right) = \frac{\partial}{\partial \rho} \left(\frac{\eta_{||}}{\mu_0} \frac{I}{V' \langle R^{-2} \rangle} \frac{\partial}{\partial \rho} \left(\frac{V'}{I} \left\langle \frac{|\nabla \rho|^2}{R^2} \right\rangle \frac{\partial \psi}{\partial \rho} \right) - \frac{\eta_{||}}{I \langle R^{-2} \rangle} \langle (J_{CD} + J_{BS}) B_\phi \rangle \right). \quad (2.10)$$

Here, n_s , T_s , S_s , P_s , D_s , V_s , χ_s , V_{E_s} , B_p , $\eta_{||}$, J_{CD} , J_{BS} , I , $\rho = r/a$, are the plasma density, the plasma temperature, the particle source, the heating source, the particle diffusion coefficient, the particle convective velocity, the thermal diffusion coefficient, the thermal convective velocity, the poloidal magnetic field, the plasma resistivity along the magnetic field lines, the induced current density, the bootstrap current density, $B_t R$ and the normalized minor radius, respectively, and $V' = dV/d\rho$ is the radial derivative of plasma volume. The symbols in $\langle \rangle$ denote a magnetic flux surface average defined in Eq. 2.7. The quasi-neutrality condition is assumed in the particle transport equation. Transport coefficients include neoclassical and turbulent contributions and are described by linear summation. J_{CD} includes the contributions of the current driven by NBI and the current driven by RF. The transport equations are solved by the implicit difference method.

Impurity transport

The difference between the transport analysis of fusion plasma main components and that of impurity ions is whether the particles are fully or partially ionized, and therefore whether the charge number of the particles is one or more than one. For impurity ions of valence Z , the steady state is called the coronal equilibrium, where the amount of increase in the ionization ($Z - 1 \rightarrow Z$) and the recombination ($Z + 1 \rightarrow Z$) is balanced by the amount of decrease in the ionization ($Z \rightarrow Z + 1$) and the recombination ($Z \rightarrow Z - 1$). The coronal equilibrium is based on the following assumptions:

- Neglecting the radiative transitions from the upper energy levels.

- Neglecting the electron impact excitation to other energy levels.
- The excitation, ionization, and recombination time scales are sufficiently shorter than the plasma time scales.
- Neglecting the effect of the absorbed photons on the spectra and atomic processes.

Since the coronal equilibrium is effective for stationary ions, the ionization state of impurity ions in the plasma results in a spatial deviation from coronal equilibrium due to the transport phenomena, as shown in Fig. 2.3. Therefore, the spatial distribution should be predicted including transport and atomic processes. In the impurity transport calculations using the TASK/TI code, the following particle transport equations including atomic processes are solved for k -valent ionized impurity ions.

$$\frac{1}{V'} \frac{\partial}{\partial t} (V' n_k) = - \frac{1}{V'} \frac{\partial}{\partial \rho} \left(V' \langle |\nabla \rho| \rangle n_k V_k - V' \langle |\nabla \rho|^2 \rangle D_k \frac{\partial n_k}{\partial \rho} \right) + \gamma_{k-1} n_e n_{k-1} - \gamma_k n_e n_k - \alpha_k n_e n_k + \alpha_{k+1} n_e n_{k+1} \quad (k = 2 - Z_{max}), \quad (2.11)$$

$$\frac{1}{V'} \frac{\partial}{\partial t} (V' n_k) = - \frac{1}{V'} \frac{\partial}{\partial \rho} \left(V' \langle |\nabla \rho| \rangle n_1 V_1 - V' \langle |\nabla \rho|^2 \rangle D_1 \frac{\partial n_1}{\partial \rho} \right) - \gamma_1 n_e n_1 + \alpha_2 n_e n_2 - \alpha_1 n_e n_1 + S_0 \quad (k = 1), \quad (2.12)$$

where, γ_k is the ionization rate, α_k is the recombination rate and S_0 is the impurity source. When impurity particles are fully ionized, the valence Z_{max} is $Z_{max} = 6$ for carbon, $Z_{max} = 8$ for oxygen, $Z_{max} = 10$ for neon, $Z_{max} = 26$ for iron, and $Z_{max} = 74$ for tungsten. In the TASK/ code, the simultaneous equations of transport are solved for all of these ionization states. The impurity temperatures are assumed to be equivalent to the bulk ion temperature, i.e., $T_Z = T_i$.

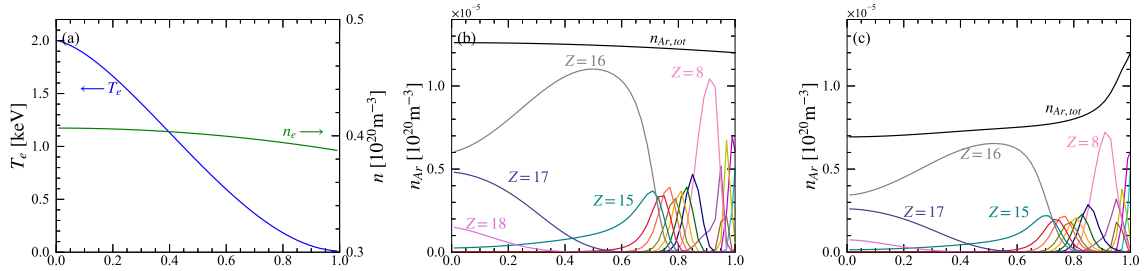


Figure 2.3: Radial profiles of (a) plasma temperature and density used in the impurity calculation, (b) Ar density assuming the coronal equilibrium and (c) Ar density at the steady state obtained by the transport calculation. The boundary density of Ar is given by $n_{Ar,edge} = 1.2 \times 10^{15} \text{ m}^{-3}$ and the diffusion coefficient by $D_{Ar} = 0.5 \text{ m}^2 \text{ s}^{-1}$. The profiles of each ionized Ar^{Z+} varies with transport process, as well as the total Ar density. In practice, the transport processes include a Z -dependence, so the deviation from coronal equilibrium can be large.

A simplified impurity-neutral particle diffusion equation is expressed as follows:

$$\frac{1}{V'} \frac{\partial}{\partial t} (n_{0s} V') = -\frac{1}{V'} \frac{\partial}{\partial \rho} \left(-V' \langle |\nabla \rho|^2 \rangle D_{0s} \frac{\partial n_{0s}}{\partial \rho} \right) - \gamma_0 n_e n_{0s} + S_{0s}, \quad (2.13)$$

where the diffusion coefficient D_{0s} is given with the thermal velocity v_{0s} , the ionization cross-section $\langle \sigma_{IZ} v \rangle$, and the charge exchange cross-section $\langle \sigma_{CX} v \rangle$ by the following formula [97]:

$$D_{0s} = \frac{v_{0s}^2}{3n_i (\langle \sigma_{IZ} v \rangle + \langle \sigma_{CX} v \rangle)}. \quad (2.14)$$

The ionization coefficient γ_0 is defined using the ionization energy E_∞^{IZ} as follows [98]:

$$\gamma_0 = 10^{-5} \frac{\left(\frac{T_e}{E_\infty^{IZ}} \right)^{0.5}}{(E_\infty^{IZ})^{1.5} \left(6.0 + \frac{T_e}{E_\infty^{IZ}} \right)} \exp \left(-\frac{E_\infty^{IZ}}{T_e} \right) \times 10^{-6}. \quad (2.15)$$

2.4.2 Heating and particle source model

In this study, the radial profiles of external heating power with the NBI and RF are given by the following Gaussian profile:

$$P_{IN} = P_0^{IN} \exp \left[-\left(\frac{\rho - \rho_0^{IN}}{\rho_w^{IN}} \right)^2 \right]. \quad (2.16)$$

Here, P_0^{IN} is the intensity of input power, ρ_0^{IN} is the radial position of input power, ρ_w^{IN} is the radial width of input power.

OH heating

The toroidal electric field by the magnetic fields generated by the induction coils increases the momentum of the electrons, resulting in a plasma current. In plasmas with finite resistivity, Joule heat is generated when Coulomb collisions occur between electrons and ions. The plasma heating method using the Joule heat is called ohmic (OH) heating, and the power P_{OH} is expressed as

$$P_{OH} = \eta j^2, \quad (2.17)$$

where, η is the plasma resistivity, j is the plasma current density. In this study, η is evaluated with the NCLASS code. Since η strongly depends on the electron temperature, the OH heating decreases as the plasma temperature increases. For Spitzer's formula, the plasma resistivity is given by $\eta = \frac{m \cdot e}{ne^2 \tau_{ei} \propto T_e^{-3/2}}$.

Equipartition power

The temperature difference in the particle species causes an exchange of energy. The equipartition power P_{EQ} based on the energy transfer due to collisions between electrons and ions is given as follows [99]:

$$P_{EQ} = \frac{3}{2} n_e \frac{T_e - T_i}{\tau_{ei}}. \quad (2.18)$$

Energy is transferred in characteristic time (electron-ion collision time $\tau_{ei} = \frac{1}{\nu_{ei}}$) when the Coulomb collisions occur with particles of different temperatures. For $T_e > T_i$, energy is transferred from electrons to ions, and for $T_e < T_i$, from ions to electrons.

Neutral beam injection (NBI)

Fast neutral beams injected into plasmas are transformed into fast ions through charge exchange and other processes. If the fast ion beam energy E_b is larger than the critical beam energy $E_c = \frac{m_b v_c^2}{2} = 14.8 T_e A_b Z_{eff}^{\frac{3}{2}}$, the beam particles are moderated by the electrons, and if $E_b < E_c$, the beam particles are moderated by the ions. The time evolution of the energy density of fast particles W_b is given as follows [100]:

$$\frac{dW_b}{dt} = P_{NB} - \frac{W_b}{\tau_b}, \quad (2.19)$$

where P_{NB} is the input power and τ_b is the slowing down time of the high energy ions given as follows:

$$\tau_b = \frac{1}{2} \frac{m_b T_e^{\frac{3}{2}}}{5 Z_i^2 n_e \ln \Lambda} \left(1 - \phi \left(\frac{E_{b0}}{E_c} \right) \right), \quad (2.20)$$

and ϕ is the fraction of the initial ion beam energy transferred the ions in plasmas given as follows:

$$\phi(x) = \frac{1}{x} \left[\frac{1}{3} \ln \left(\frac{1 - x^{\frac{1}{2}} + x}{(1 + x^{\frac{1}{2}})^2} \right) + \frac{2}{\sqrt{3}} \left(\tan^{-1} \left(\frac{2x^{\frac{1}{2}} - 1}{\sqrt{3}} \right) + \frac{\pi}{6} \right) \right]. \quad (2.21)$$

The dependence of ϕ on the ratio of beam energy to critical beam energy is shown in Fig. 2.4. As the beam energy increases, ϕ decreases, and the fraction of energy transferred to the ions decreases. As the plasma temperature increases, E_c increases, and the fraction of energy transferred to the ions increases.

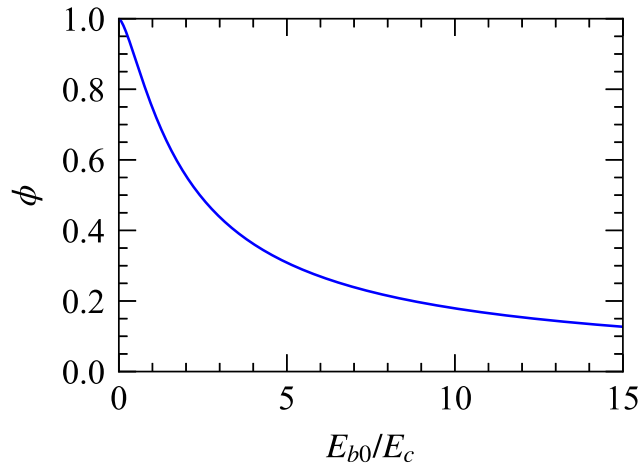


Figure 2.4: Fraction of the initial neutral beam energy transferred the ions in plasmas.

Input NBI power is distributed for electrons and ions respectively as follows:

$$P_{NB,e} = \frac{W_b}{\tau_b} \left(1 - \phi \left(\frac{E_{b0}}{E_c} \right) \right), \quad (2.22)$$

$$P_{NB,i} = \frac{W_b}{\tau_b} \phi \left(\frac{E_{b0}}{E_c} \right) \frac{\frac{n_i Z_i^2}{m_i}}{\sum \left(\frac{n_i Z_i^2}{m_i} \right)}. \quad (2.23)$$

Although the pencil beam model can be used in the TASK/TR code, in this study, a Gaussian profile gives the P_{NB} as in Eq. 2.16, and the particle source by NBI is defined as follows:

$$S_{NB} = \frac{P_{NB}}{E_{b0}}. \quad (2.24)$$

The momentum of electrons in the plasma is increased by collisions with the beam ions and decreased by collisions with thermal ions, resulting in a current drive at a steady state. The current density driven by NBI is expressed by the following equation [101]:

$$j_{NB} = \frac{W_b}{\tau_b} \frac{2eZ_b\tau_s}{m_b v_{b0} \frac{E_c}{E_{b0}}} \left(1 - \frac{Z_b}{Z_{eff}(1 - C_1)} \right) C_2, \quad (2.25)$$

where, Z_b is the charge number of beam ions, τ_s is the slowing-down time due to the electron collision, n_b is the density of beam ions and v_{b0} is the velocity of injected beam particles.

Since $\tau_s = \frac{6\pi\sqrt{2}\pi\epsilon_0^2 m_b T_e^{\frac{3}{2}}}{m_e^{\frac{1}{2}} n_e Z_b^2 e^4 \ln \Lambda}$ increases in proportion to the 3/2 power of the electron temperature, it is expected that the current drive efficiency of NBI improves with increasing electron temperature. The correction terms C_1 and C_2 are estimated by the Fokker-Planck analysis, and the formula are given as follows [102]:

$$C_1 = \sqrt{\frac{r}{R}} \left(1.55 + \frac{0.85}{Z_{eff}} \right) - \frac{r}{R} (0.2 + 1.55 Z_{eff}), \quad (2.26)$$

$$C_2 = \frac{\left(\frac{E_{b0}}{E_c} \right)^2}{4 + 3 \left(0.8 \frac{Z_{eff}}{A_b} \right) + \left(\frac{E_{b0}}{E_c} \right)^2 \left(\frac{E_{b0}}{E_c} + 1.39 + 0.61 \left(0.8 \frac{Z_{eff}}{A_b} \right) \right)}. \quad (2.27)$$

Radio frequency (RF) heating

The ray tracing method is often used to analyze ECH/LHH to evaluate heating power distribution. Evaluation using TASK/WR [103] for ECH is available in the TASK code. For the evaluation of the heating power distribution of ICRH, the wave propagation code TASK/WM [104] has been developed because of the problem that mode conversion cannot be adequately represented in the ray tracing method. Currently, TASK/WR and transport codes are linked, but TASK/WM is not yet available in transport simulation scheme. For simplicity, P_{RF} is given in this study by the Gaussian profile as in Eq. 2.16. The current density driven by the RF is used as follows [105]:

$$j_{RF} = \eta_{CD} P_{RF}. \quad (2.28)$$

The current drive efficiency η_{CD} is used as follows:

$$\eta_{CD} = \frac{0.384}{\ln \Lambda} \frac{T_e}{n_e} \tilde{\eta}_{CD}, \quad (2.29)$$

where $\tilde{\eta}_{CD}$ is a dimensionless quantity determined numerically and is given as follows:

$$\tilde{\eta}_{CD} = CMR\eta_0. \quad (2.30)$$

Here,

$$C = 1 - \exp[-(cy_t)^m], \quad (2.31)$$

$$M = 1 + a\left(\frac{\lambda_t}{w}\right)^3, \quad (2.32)$$

$$R = 1 - \epsilon^{0.77} \frac{\sqrt{3.5^2 + w^2}}{3.5\epsilon^{0.77} + w}, \quad (2.33)$$

where $\epsilon = \frac{r}{R}$, $y_t = \frac{1-\lambda_t^2}{\lambda_t^2} w^2$, $\lambda_t = \sqrt{\frac{\epsilon}{1+\epsilon}}$, $w = \frac{v_c}{v_{th,e}N_{\parallel}}$, and $N_{\parallel}$ is the refractive index along the magnetic fields, and η_0 is used as follows:

$$\eta_0 = \frac{K}{w} + D + \frac{4w^2}{5 + Z_{eff}}. \quad (2.34)$$

The respective coefficients for EC, LH, and IC waves are shown in Table 2.1.

Table 2.1: Numerical coefficients in Eq. 2.31-2.34.

parameter	EC	LH	IC
c	0.3890	0.3890	0.00987
m	1.38	1.38	2.48
a	0.0	0.0	12.3
K	3.00/Z	3.00/Z	11.91/(0.678+Z)
D	3.83/Z ^{0.707}	3.83/Z ^{0.707}	4.13/Z ^{0.707}

Power loss

Energy losses caused by radiation in the core plasma are mainly due to the bulk radiation losses, the cyclotron radiation losses, and the impurity radiation losses. Electrons, which have much lower mass than ions, are accelerated by an electric field and then emit radiation. The power of bremsstrahlung produced by the collision of charged particles is given by the following:

$$P_{brm} = 5.35 \times 10^{-37} n_e n_i Z_i T_e^{\frac{1}{2}}. \quad (2.35)$$

The cyclotron radiation tends to be more significant in the plasma center, where the temperature is higher, but the radiation reflected by the walls is reabsorbed near the plasma boundary. The CYTRAN [106] code evaluates the loss and absorption due to cyclotron radiation. The radiation losses due to the impurity ions are discussed in Sec. 2.4.3.

Gas puffing and pellet injection

The particle source assuming the gas puffing method is given by the following Gaussian expression:

$$S_p = S_0^{S_p} \exp \left[- \left(\frac{\rho - \rho_0^{S_p}}{\rho_w^{S_p}} \right)^2 \right], \quad (2.36)$$

where $S_0^{S_p}$ is the intensity of the particle feeding, $\rho_0^{S_p}$ is the radial position of particle source, $\rho_w^{S_p}$ is the radial width of particle source. The pellet absorption distribution can be evaluated using ablation models based on neutral-cloud gas shielding (NGS), widely used for the pellet particle source analysis [107, 108]. Simulations have been performed using the TASK/TR code in cooperation with the pellet melting and drift code HPI2 [109, 110].

Impurity source

In the current TASK/TI code, the impurity inflow condition is set by the following boundary conditions:

1. Reflection model : $\left. \frac{\partial n_Z}{\partial \rho} \right|_{\rho=1.0} = 0$
2. Fixed value model : $n_Z|_{\rho=1.0} = n_{Z,edge}$
3. Fixed influx model : $\Gamma_Z|_{\rho=1.0} = \Gamma_{Z,in}$
4. Fixed decay length model : $\left. \frac{\frac{\partial n_Z}{\partial \rho}}{n_Z} \right|_{\rho=1.0} = -\lambda$

In this study, the fixed value model is used basically, and the valence distribution at the plasma boundary is determined by the coronal equilibrium. Evaluation of impurity inflow conditions using a scrape-off layer plasma code is needed in future studies.

2.4.3 Ionization, recombination and radiation coefficients

The OPEN-ADAS (Atomic Data & Analysis Structure) [111] is used to evaluate the coefficients of ionization, recombination, and radiation. The ADAS is an atomic data library useful for the analysis of fusion plasmas and astronomical plasmas. The ADPAK (Atomic Data PacKage) [68] is also used as an atomic process library. Compared to the ADPAK, the ADAS has data for a broader range of electron temperatures and can be adapted to the plasma periphery. The ADAS also includes dependence on electron density, making it applicable to low-density regions, where the temperature is lower. The ADAS classifies atomic data into atomic processes, such as charge exchange cross sections (ADF01), resolved specific ion data collections (ADF04), electron impact ionization rate coefficients (ADF07), radiative recombination rate coefficients (ADF08), resolved dielectronic recombination coefficients (ADF09), iso-nuclear master files (ADF11), etc. In this study, we use the ADF11 data sets of effective recombination coefficients (ACD), effective ionization coefficients (SCD), charge exchange effective recombination coefficients (CCD), line power driven by excitation of dominant ions (PLT), continuum and line power driven by recombination and bremsstrahlung of dominant ions (PRB). Data files are available according to the year of preparation. This

study will use the following data files: carbon, 96s; neon, 96s; argon, 89s; iron, 89s; nickel, 89s; and tungsten, 89s. By using the OPEN-ADAS data, the ionization state of multiply charged ions and radiative loss power can be obtained, as shown in Fig. 2.5. Depending on the electron temperature for the radial direction, the ionization and recombination coefficients are evaluated to determine the valence distribution of impurity ions.

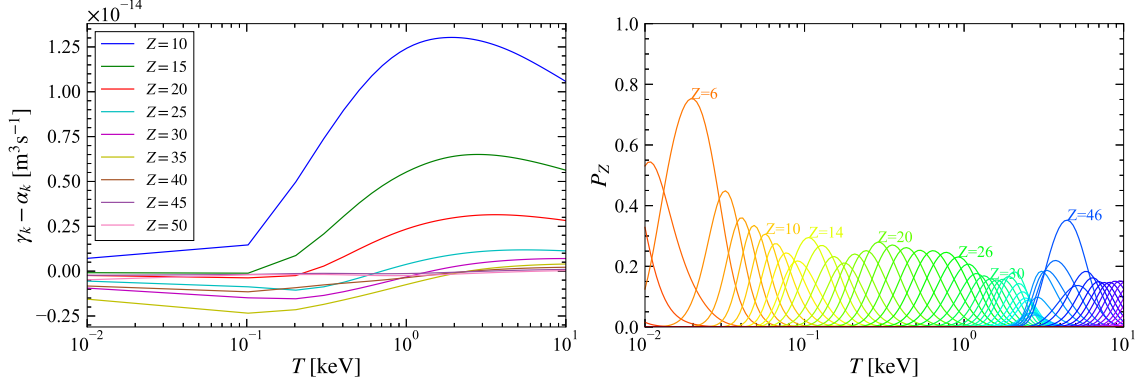


Figure 2.5: Dependence of tungsten evaluated with OPEN-ADAS. Figure (a) shows the dependence of the differences between the ionization and recombination coefficients on the electron temperature, and figure (b) shows the dependence of the fraction of each valence on electron temperature.

2.4.4 Neoclassical transport model

NCLASS

NCLASS [21] is a fluid code to calculate the neoclassical transport properties for multiple particle species of axisymmetric plasmas with arbitrary aspect ratio, geometry, and collisionality. The coefficients of the neoclassical particle fluxes are calculated by using the plasma and magnetic field profiles as parameters. For impurity ion species with the same ion temperature as the bulk ions, the impurity particle flux is given as follows:

$$\begin{aligned} \Gamma_k^{NCLASS} &= -D_k \frac{\partial n_k}{\partial \rho} + D_k n_k \left(\sum_{l \neq k} \left(g_{n_l} \frac{\frac{\partial p_l}{\partial \rho}}{p_l} \right) + g_{T_i} \frac{\frac{\partial T_i}{\partial \rho}}{T_i} + g_{T_e} \frac{\frac{\partial T_e}{\partial \rho}}{T_e} \right) \\ &\sim -D_k \frac{\partial n_k}{\partial \rho} + D_k n_k \left(g_{n_i} \frac{\frac{\partial n_i}{\partial \rho}}{n_i} + (g_{n_i} + g_{T_i}) \frac{\frac{\partial T_i}{\partial \rho}}{T_i} \right), \end{aligned} \quad (2.37)$$

where the parameter g is determined by the momentum exchange between the various particle species and the plasma state, and the g_{T_e} term can be neglected because of the small momentum exchange between electrons and impurity ions [112]. The first term is the diffusion term, the second and third terms are convection terms, and the density and temperature gradients of the bulk ions drive the flow. As shown in Fig. 2.6, the neoclassical transport of the impurity is strongly dependent on the background plasma profiles, and negative ion density gradients possibly enhance the impurity accumulation in the core plasma region. Figure 2.7 shows the dependence of convection velocity in case 2 in Fig. 2.6(c)

on the bulk ion density and temperature and their gradients. The inward flow is increased for increased density gradients and the outward flow is increased for increased temperature gradients. In a typical tokamak device, the ion density and temperature decrease from the plasma center to the edge, so the g_{n_i} term drives the inward flow, and the g_{T_i} term drives the outward flow. The $g_{T_i} < 0$ then, giving a temperature screening effect. In the neoclassical transport, the net flow direction is determined by the balance between the velocity driven by the bulk ion density gradient and the velocity driven by the bulk ion temperature gradient. In the NCLASS code, convection components include the $\langle E \cdot B \rangle$ (Ware pinch) and the external parallel forces in the banana-plateau regime. The Ware pinch flux is $\Gamma_{Ware} = Zn \langle E \cdot B \rangle$ and is driven to the plasma center (inward direction) by parallel electric fields. The evaluation of the collision operator is based on the Hirshman-Sigmar approach [113]. In addition, the NCLASS code can evaluate the Hirshman and Sigmar moment equation for poloidal rotation velocity in any collisional regime, without being limited by trace concentrations, and taking into account collision effects between other particle species [114]. However, since this collision operator does not include energy scattering, the evaluation of the temperature screening coefficient in the Pfirsch-Schlüter regime, which has high collisionality, is not valid [115]. Furthermore, the NCLASS code does not include effects related to the poloidal asymmetry, limiting its applicability to the transport calculations of high-Z impurity. Note that it is very reliable for calculating the thermal diffusion coefficients of electrons and main ion species, bootstrap currents, and plasma resistivity. The bootstrap current is evaluated by the following equation:

$$\langle j_{BS} \cdot B \rangle = \sum_s \left(C_{T_s}^{BS} \frac{\partial T_s}{\partial \rho} + C_{p_s}^{BS} \frac{\partial p_s}{\partial \rho} \right), \quad (2.38)$$

where the coefficients of $C_{T_s}^{BS}$ and $C_{p_s}^{BS}$ are obtained from NCLASS.

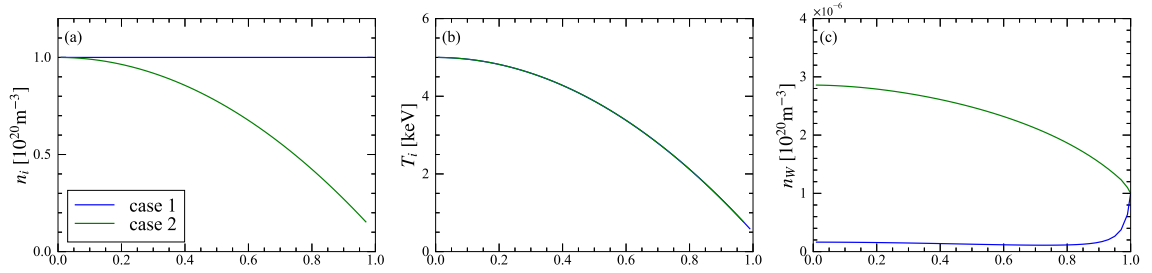


Figure 2.6: Radial profiles of (a) main ion density and (b) main ion temperature, which are fixed profiles in the impurity transport calculations. In both case 1 and case 2, the main ion temperature profiles are set in parabolic, while the ion density profiles are set with flat in case 1 and parabolic in case 2. Figure (c) shows the radial profiles of W density evaluated with NCLASS.

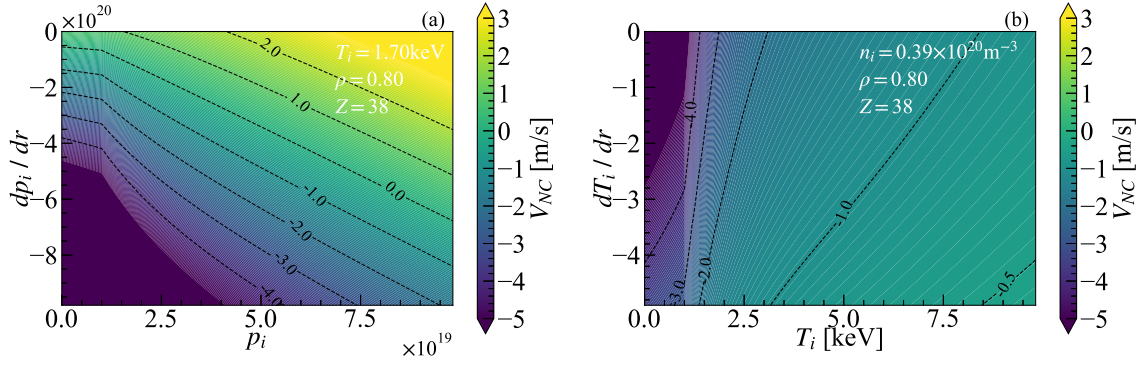


Figure 2.7: Dependence of the neoclassical convection velocity of tungsten on (a) main ion pressure and its gradient, and (b) main ion temperature and its gradient. Note that the ion temperature T_i is constant in figure (a), and the ion density n_i is constant in figure (b). The direction of convection is inward for negative value and outward for positive value.

NCLASS with poloidal asymmetry

Neoclassical transport coefficients for impurities are calculated by introducing geometrical factors (P_A and P_B) into NCLASS to consider the poloidal asymmetry due to the toroidal rotation. The neoclassical impurity flux in the Pfirsch-Schlüter regime can be written in its general form as [116]:

$$\langle \Gamma_Z^{Neo} \cdot \nabla r \rangle = \frac{q^2 n_i \langle n_Z \rangle T_i v_{iZ}}{\epsilon^2 m_i \langle \Omega_i^2 \rangle n_Z Z} \left[\left(\left\langle \frac{n}{b^2} \right\rangle - \left\langle \frac{b^2}{n} \right\rangle^{-1} \right) \tilde{g} - \left(1 - \left\langle \frac{b^2}{n} \right\rangle^{-1} \right) \tilde{u} \right], \quad (2.39)$$

where $n = n_Z / \langle n_Z \rangle$ and $b^2 = B^2 / \langle B^2 \rangle$. The factor $\tilde{g}$ and $\tilde{u}$ can be written as [117–119]

$$\tilde{g} = -\frac{d \ln p_Z}{dr} + \frac{d \ln p_i}{dr} - C_0^Z \frac{d \ln T_i}{dr}, \quad (2.40)$$

$$\tilde{u} = (C_0^Z + k_i) \frac{d \ln T_i}{dr}. \quad (2.41)$$

Here, the parallel ion heat flux term in $\tilde{u}$ is neglected (i.e., $\langle u_{2||i} \rangle = 0$). Since the momentum exchange between the impurity and the electron is small, the contribution of the electron temperature gradient is also neglected. With $D_c = v_{th,i}^2 / \langle \Omega_i^2 \rangle v_{iZ} \approx v_{th,i}^2 / \langle \Omega_i^2 \rangle v_{ii} Z^2 m_i / m_Z$, Eq. 2.39 can be rewritten as follows:

$$\begin{aligned} \langle \Gamma_Z^{Neo} \cdot \nabla r \rangle = q^2 D_c Z \frac{\langle B^2 \rangle}{2 \epsilon^2} \times \\ \left[\left(\left\langle \frac{n_Z}{B^2} \right\rangle - \left\langle \frac{B^2}{n_Z} \right\rangle^{-1} \right) \left(-\frac{1}{Z} \frac{d \ln n_Z}{dr} + \frac{d \ln n_i}{dr} - (C_0^Z - 1) \frac{d \ln T_i}{dr} \right) \right. \\ \left. - \left(1 - \left\langle \frac{B^2}{n_Z} \right\rangle^{-1} \right) (C_0^Z + k_i) \frac{d \ln T_i}{dr} \right], \end{aligned} \quad (2.42)$$

where the introduction of geometrical factors P_A and P_B leads to the following impurity flux [70, 115, 120–122]:

$$\langle \Gamma_Z^{Neo} \cdot \nabla r \rangle = \langle n_Z \rangle q^2 D_c Z \left[P_A \left(\frac{1}{Z} \frac{1}{L_{n_Z}} - \frac{1}{L_{n_i}} + (C_0^Z - 1) \frac{1}{L_{T_i}} \right) + P_B (C_0^Z + k_i) \frac{1}{L_{T_i}} \right], \quad (2.43)$$

with

$$P_A = \frac{1}{2\epsilon^2} \frac{\langle B^2 \rangle}{\langle n_Z \rangle} \left[\left\langle \frac{n_Z}{B^2} \right\rangle - \left\langle \frac{B^2}{n_Z} \right\rangle^{-1} \right], \quad (2.44)$$

$$P_B = \frac{1}{2\epsilon^2} \frac{\langle B^2 \rangle}{\langle n_Z \rangle} \left[\left\langle \frac{n_Z}{B^2} \right\rangle - \left\langle \frac{B^2}{n_Z} \right\rangle^{-1} \right], \quad (2.45)$$

and the gradient length of each plasma quantity $1/L_X = -\nabla X/X$. Absence of the poloidal asymmetry results in $P_A = 1$ and $P_B = 0$. The coefficients k_i and C_0^Z are related to the collisionality, the neoclassical ion flow coefficient and the friction coefficient, respectively. For the case where the impurity is in the Pfirsch-Schlüter regime and the main ion is in the banana regime, $-(C_0^Z - 1)$ becomes close to -0.5 , and a numerical formula is often used [113, 123]. More recently, extended formulas have been proposed with ion-electron collisional heat exchange for impurities, which is given as follows [124, 125]:

$$C_0^Z = \frac{\frac{1.5}{1+f_1(\frac{A_i}{A_A})} - \frac{0.29+0.68\alpha_{trace}}{0.59+\alpha_{trace}+\frac{1.34+f_2}{g^2}}}{1+f_3\mu_{ie}g^2}, \quad (2.46)$$

where,

$$\alpha_{trace} = \frac{n_Z Z_Z^2}{n_i Z_i^2}, \quad (2.47)$$

$$g = v_i^* \epsilon^{\frac{3}{2}}, \quad (2.48)$$

$$\mu_{ie} = \frac{96\sqrt{2}}{125Z_i^2} \sqrt{\frac{m_e}{m_i}} \left(\frac{T_i}{T_e} \right)^{\frac{3}{2}}, \quad (2.49)$$

$$f_1 = \frac{-6.8 \times 10^5 + 2.5 \times 10^6 Z}{1 + 6.0 \times 10^5 Z^{1.6}}, \quad (2.50)$$

$$f_2 = \frac{88.3 + 10.5Z}{1 + 0.2Z^{2.6}}, \quad (2.51)$$

$$f_3 = \frac{-4.4 \times 10^6 + 2.7 \times 10^6 Z}{1 + 5.3 \times 10^6 Z^{0.8}}. \quad (2.52)$$

On the other hand, $C_0^Z + k_i$ is 0.33 takes the limit of $\epsilon \ll 1$, and the expression for k_i is given by [126–128], which give the following expression including the effect of rotation [125]

$$k_i = \left(\frac{k_{i,0} + l_1 (f_i v_i^*)^{0.5} + l_2 v_i^{*0.25}}{1 + l_3 v_i^{*0.5}} - l_4 l_5 v_i^{*2} f_i^6 + l_6 v_i^{*0.25} \right) \left(\frac{1}{1 + l_4 v_i^{*2} f_i^6} \right), \quad (2.53)$$

with

$$k_{i,0} = -\frac{0.53(1 + 0.65M_i^{1.5})}{0.53} \frac{1 - f_t}{1 + 1.16(1 - 0.97M_i^{1.56})f_t - 0.98(1 - 1.23M_i^{1.7})f_t^2}, \quad (2.54)$$

and

$$l_1 = 0.38 + 5.7(1 - f_t)^{6.7}, \quad (2.55)$$

$$l_2 = -1.52 + 38.4(1 - f_t)^3 f_t^{2.1} + (-1 + 2.6f_t) M_i^{2.5} (1 - 0.6f_t), \quad (2.56)$$

$$l_3 = 0.25 + 1.2(1 - f_t)^{3.65}, \quad (2.57)$$

$$l_4 = 0.1(1 - f_t)^{1.46} f_t^{4.33} + (0.05 - 0.04f_t) M_i^{2.5}, \quad (2.58)$$

$$l_5 = 0.8 + (0.585 + 1.25f_t) M_i, \quad (2.59)$$

$$l_6 = \frac{-0.05 + 1.95f_t^{2.5}}{1 + 2.55f_t^{1.7}} + \frac{-0.22 - 14.6f_t^{6.3}}{1 + 5.62f_t^{5.72}} M_i^{1.5}, \quad (2.60)$$

where M_i is the Mach number of main ions and f_t is the trapped particle fraction evaluated with Lin-Liu and Miller model as follows [129]:

$$f_t = 0.75(1.5\sqrt{\epsilon}) + 0.25\left(\frac{3\sqrt{2}}{\pi}\sqrt{\epsilon}\right). \quad (2.61)$$

Using Eqs. 2.46 and 2.53, the dependences of C_0^Z and k_i on ν_i^* are evaluated in Fig. 2.8. In the low collisionality regime, $C_0^Z \sim 1.5$ and $k_i \sim -0.5$, while in the high collisionality regime, $C_0^Z \sim 0$ and $k_i \sim -1.0$.

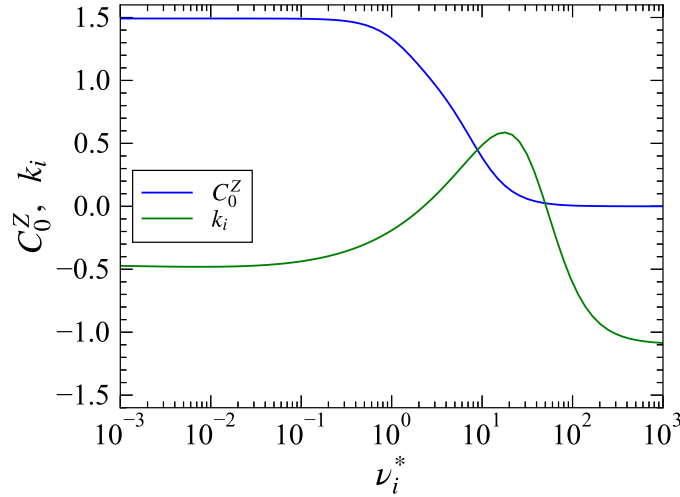


Figure 2.8: Dependences of C_0^Z and k_i on the bulk ion collisionality for sufficiently decreased impurity strength parameter ($\alpha_{trace} \ll 1$). Here, tungsten ion ($A = 183.3$, $Z = 35$) and $\epsilon = 0.06$ are used.

Equation 2.43 assumes a convection coefficient $K = 1$ multiplied to the impurity density gradient and main ion density gradient terms. This assumption is known to be valid for the

case where the impurity is in the Pfirsch-Schlüter regime and the main ion is in the banana regime. Therefore, rewriting Eq.2.43 to Eq.(7) in Ref. 115 using the temperature screening coefficient H , the neoclassical flux of the impurity is given as

$$\frac{\langle \Gamma_Z^{Neo} \cdot \nabla r \rangle}{\langle n_Z \rangle} = q^2 D_c Z \left[P_A \left(K \frac{1}{Z} \frac{1}{L_{n_Z}} - K \frac{1}{L_{n_i}} - H \frac{1}{L_{T_i}} \right) - P_B (C_0^Z + k_i) \frac{1}{L_{T_i}} \right]. \quad (2.62)$$

As shown in Ref. 115, the above analytical flux formula combined with the NCLASS code describes impurity transport including the poloidal asymmetry. Using the diffusion coefficient D_{asym}^{NCLASS} and convection velocity V_{asym}^{NCLASS} with the poloidal asymmetry (labeled with 'asym'), Eq. 2.62 can be simply expressed as

$$\frac{\langle \Gamma_Z^{Neo} \cdot \nabla r \rangle}{\langle n_Z \rangle} = D_{asym}^{NCLASS} \frac{1}{L_{n_Z}} + V_{asym}^{NCLASS}, \quad (2.63)$$

where V_{asym}^{NCLASS} is defined as

$$V_{asym}^{NCLASS} = V_{N,asym}^{NCLASS} + V_{T,asym}^{NCLASS}, \quad (2.64)$$

with the poloidally asymmetric convection velocity $V_{N,asym}^{NCLASS}$ and $V_{T,asym}^{NCLASS}$. Here the poloidally asymmetric transport coefficients can be written as follows using the symmetric (labeled with 'sym') transport coefficients D_{asym}^{NCLASS} , $V_{N,asym}^{NCLASS}$ and $V_{T,asym}^{NCLASS}$.

$$\begin{aligned} D_{asym}^{NCLASS} &= q^2 D_c K P_A \\ &= D_{sym}^{NCLASS} P_A, \end{aligned} \quad (2.65)$$

$$\begin{aligned} V_{N,asym}^{NCLASS} &= Z D_{sym}^{NCLASS} P_A \\ &= V_{N,sym}^{NCLASS} P_A \frac{1}{L_{n_i}}, \end{aligned} \quad (2.66)$$

$$\begin{aligned} V_{T,asym}^{NCLASS} &= -Z D_{sym}^{NCLASS} \left(\frac{H}{K} + \frac{C_0^Z + k_i}{K} \frac{P_B}{P_A} \right) \frac{1}{L_{T_i}} \\ &= V_{T,sym}^{NCLASS} \left(P_A + \frac{C_0^Z + k_i}{H} P_B \right). \end{aligned} \quad (2.67)$$

Therefore, the geometric factor that quantifies the poloidal asymmetry is $P_A = D_{asym}^{NCLASS} / D_{sym}^{NCLASS} = V_{N,asym}^{NCLASS} / V_{N,sym}^{NCLASS}$, which increases the inward pinch driven by the main ion density gradient in the presence of strong asymmetry. Since H is $-(C_0^Z - 1)$ from Eq. 2.43 and H is negative value except when the main ions are highly collisional regime, the screening effect driven by the main ion temperature gradient is reduced by the P_B term in Eq. 2.67 [120]. To evaluate the geometric factor P_A and P_B , a two-dimensional impurity distribution considering the poloidal asymmetry is given from the parallel force balance equation as follows:

$$\frac{1}{2} m_Z \nabla_{\parallel} V_Z^2 = -Z e n_Z \nabla_{\parallel} \Phi - T_i \nabla_{\parallel} n_Z + R_{Zi\parallel}, \quad (2.68)$$

where the V_Z is impurity velocity, Φ is the poloidal varying electrostatic potential, and $R_{Z\parallel}$ is the impurity-ion friction force. Integrating Eq. 2.68 gives

$$n_Z = n_Z(\theta = 0) \exp \left[-\frac{eZ\Phi(\theta)}{T_i} + \frac{m_Z\omega^2}{2T_i} (R(\theta)^2 - R(\theta = 0)^2) \right], \quad (2.69)$$

where the θ is poloidal angle and ω is the plasma angular rotation frequency. A two-dimensional impurity density distribution using flux surface averaged (FSA) density is given as follows [130]:

$$n_Z(r, \theta) = \langle n_Z(r, \theta) \rangle \frac{\exp \left[M_Z^2 \frac{R^2(r, \theta)}{R_0^2} \right]}{\left\langle \exp \left[M_Z^2 \frac{R^2(r, \theta)}{R_0^2} \right] \right\rangle}. \quad (2.70)$$

Here, only the centrifugal effects due to toroidal rotation with NBI is considered, and the temperature anisotropy due to the ICRH is not considered. Assuming that the temperatures of the impurity species and bulk ions are comparable, the Mach number of the impurity is given by

$$M_Z = M_i \sqrt{\frac{m_Z}{m_i}}, \quad (2.71)$$

and

$$M_i = \frac{V_{tor}}{v_{th,i}} = V_{tor} \sqrt{\frac{m_i}{2T_i}}. \quad (2.72)$$

Since the asymmetry of impurity density scales with the square of the Mach number from Eq. 2.70, the poloidal asymmetry increases with increasing Mach number as shown in Fig. 2.9, and the value of the geometric factor P_A increases, enhancing the neoclassical transport. In the presence of the strong NBI heating, P_A can increase to about 90 [131].

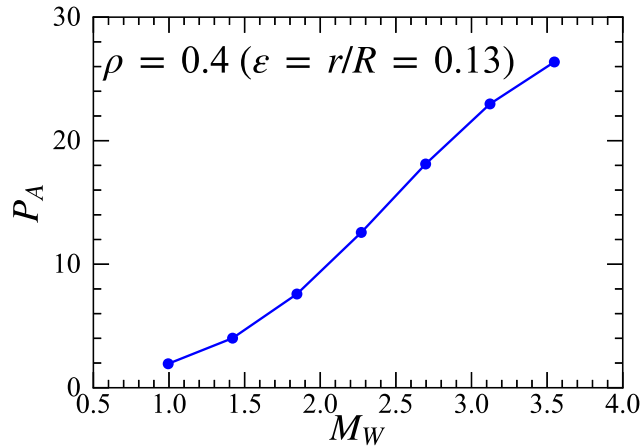


Figure 2.9: Dependence of the geometric coefficient P_A on the Mach number of tungsten M_W . The plasma profiles are used in Ref. [115]. Although there are differences in temperature screening coefficients and simulation conditions, the values are similar to the Mach number dependence of P_A evaluated in Ref. [115], verifying that the introduction of the model is consistent.

PHZ and E_r pinch model

In order to investigate the effect of toroidal rotation on the high-Z impurity transport, a analytical model based on a different physical process approach from NCLASS, which is introduced the geometric factors P_A and P_B , is also introduced for comparison. The PHZ/ E_r pinch [132] is a model proposed by Hoshino *et al.*, and is an additional pinch to the transport process of the high-Z impurity in the toroidally rotating plasmas. The PHZ (Pinch of High-Z impurity due to the atomic process) pinch model is based on the following physical processes.

1. The high-M and high-Z impurity is accelerated to a toroidal direction due to the friction force.
2. As the impurity energy increases, the drift velocity increases, and the shift of the drift orbit from the magnetic surface becomes large.
3. Changes in the ionization and recombination processes result in a variation of the charge state at the poloidal cross section.
4. Magnetic drift velocity v_d changes at the poloidal surface in proportion to $1/Z$.
5. The variation of the v_d drives inward particle pinch.

The PHZ pinch velocity is defined as

$$V_{PHZ} = \frac{v_{d0}^2}{2Z_0 \frac{C_T C_{\nabla T}}{C_Z^2 + \omega^2}}, \quad (2.73)$$

with

$$C_T = \frac{\partial v}{\partial T} = \frac{\partial n_e (\langle \sigma v \rangle_{ion} - \langle \sigma v \rangle_{recomb})}{\partial T} = \frac{n_e \partial (\gamma_k - \alpha_k)}{\partial T}, \quad (2.74)$$

$$C_Z = \frac{\partial v}{\partial Z} = \frac{\partial n_e (\langle \sigma v \rangle_{ion} - \langle \sigma v \rangle_{recomb})}{\partial Z} = \frac{n_e \partial (\gamma_k - \alpha_k)}{\partial Z}, \quad (2.75)$$

$$C_{\nabla T} = \frac{\partial T}{\partial r}, \quad (2.76)$$

and $v_{d0} = \frac{mV_{tor}^2}{Z_0 e R B}$ is the magnetic drift velocity of the impurity, $\omega = \frac{V_{tor} - \frac{E_r}{B_p}}{qR_0}$ is the angular frequency of the poloidal motion of the impurity ion, Z_0 is the charge number of the impurity ion in the ionization equilibrium at the temperature of the magnetic flux surface. The ionization coefficient γ_k and recombination coefficient α_k are evaluated with OPEN-ADAS. From the radial force balance, the radial electric field is expressed by the following equation:

$$E_r = \frac{\nabla p_i}{Z_i e n_i} + V_{tor} B_p - V_{pol} B_t, \quad (2.77)$$

where, the poloidal velocity V_{pol} is evaluated with NCLASS. Figure 2.10 shows the dependence of the PHZ pinch velocity on the radial electric field, where E_r increases as the toroidal rotation increases. The PHZ pinch is always driven in an inward direction and the value increases with increasing toroidal rotation velocity. Simulations of the JT-60

experiment show that the effect of the PHZ pinch increases the tungsten accumulation [133]. In addition, the model has been improved to include the effect of trapped particles and to consider their effect on the particle orbit [134].

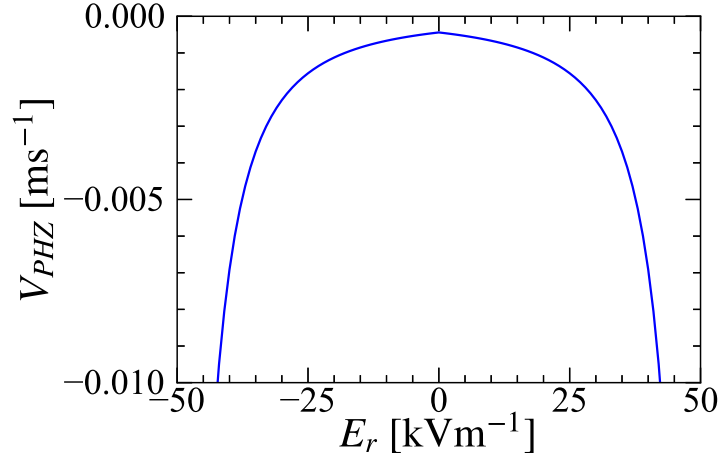


Figure 2.10: Dependence of the PHZ pinch velocity of tungsten on the radial electric field. Here, the toroidal velocity is scanned with Eq. 2.77, and the other terms are assumed to be constant. The parameters of $R = 2.9$ m, $r/a = 0.27/0.89 = 0.3$, $B_t = 2.5$ T, $T_e = 3.5$ keV, $V_{tor} = 170$ km/s, $q = 0.9$, $B_p = 0.35$ T, $n_e = 0.93 \times 10^{20}$ m $^{-3}$, $Z_0 = 41$, $C_T = 1.76$ eV $^{-1}$ s $^{-1}$, $C_Z = 1.89 \times 10^3$ s $^{-1}$, $C_{\nabla T} = 3.15 \times 10^3$ eV m $^{-1}$ are used.

On the other hand, the E_r pinch model is based on the following physical processes.

1. The high-M and high-Z impurity is accelerated to a toroidal direction due to the friction force.
2. As the impurity energy increases, the drift velocity increases, and the shift of the drift orbit from the magnetic surface becomes large.
3. The electric potential Φ changes along the drift orbit.
4. The parallel velocity changes with the variation in Φ .
5. In the co-rotation (positive E_r) case, the drift orbit expands to the radial direction.
6. The particle pinch is caused by Coulomb collisions, and the impurity moves by the distance of the expanded orbit.

The E_r pinch velocity is defined as

$$V_{E_r} = \Delta r \frac{v_c}{1 + \left(\frac{v_c}{\omega}\right)^2}, \quad (2.78)$$

with

$$\Delta r = \frac{(1 - 2\alpha_{E_r})\Delta_0}{2(1 - \alpha_{E_r})^2}, \quad (2.79)$$

$$k = \frac{ZeE_r}{mV_{tor}^2}, \quad (2.80)$$

$$\alpha_{E_r} = \frac{E_r}{B_p V_{tor}}, \quad (2.81)$$

$$\Delta_0 = \frac{v_{d0}}{\omega_0}. \quad (2.82)$$

Dependence of the E_r pinch velocity on the radial electric field is shown in Fig. 2.11. As E_r increases, the E_r pinch velocity initially increases and then decreases. This is because it is cancelled out by the increase in the poloidal angular frequency in response to the change in the drift orbit. In the regime of large values of E_r , the absolute pinch velocity is increased with increasing E_r . This model shows that the pinch velocity is considerably larger at increasing E_r , and Eq. 2.78 is valid for $\alpha_{E_r} < 1$. The direction of the E_r pinch changes with the direction of the toroidal rotation, i.e., the inward pinch is driven for CO-rotation (same direction as the plasma current) and the outward pinch is driven for CTR-rotation (opposite direction to the plasma current).

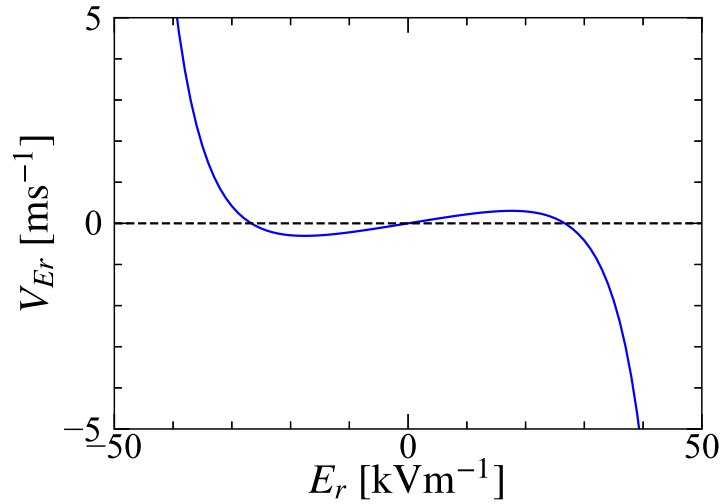


Figure 2.11: Dependence of the E_r pinch velocity of the tungsten ion on the radial electric field. Parameters are the same as in Fig. 2.10.

2.4.5 Turbulent transport mode

CDBM

Instabilities possibly occur due to reconnection of magnetic field lines with the finite plasma resistivity. The plasma resistivity is lower in high-temperature plasmas because of the lower collision frequency. However, the instability of current-diffusive ballooning mode caused by the current diffusion effect due to electron viscosity is important in high-temperature plasmas. The current-diffusive ballooning mode is stabilized by the thermal conductivity

of the electrons and the viscosity of the ions. The balance of these destabilization and stabilization effects sustains a steady turbulent state, which induces turbulent transport [135]. This state is called self-sustained turbulence, which is a stable state, and the following transport coefficient is given by the following [136]

$$\chi_{CDBM} = F(s, \alpha) F_{\kappa} F_{E \times B} \alpha^{\frac{3}{2}} \frac{c^2}{\omega_{pe}^2} \frac{v_A}{qR}, \quad (2.83)$$

where, c is the speed of light, $\omega_{pe} = \frac{n_e e^2}{m_e \epsilon_0}$ is the electron plasma frequency, $v_A = \sqrt{\frac{B^2}{\mu_0 n_e \sum m_i n_i}}$ is the Alfvén velocity, q is the safety factor, $s = \frac{r}{q} \frac{dq}{dr}$ is the magnetic shear, $\alpha = -q^2 R \frac{d\beta}{dr}$ is the normalized pressure gradient, and $\kappa = -\frac{r}{R} \left(1 - \frac{1}{q^2}\right)$ is the magnetic curvature. The normalized pressure gradient includes the contribution of a fast ion [137]. The fitting parameter $C = 12.0$, which is correction between the CDBM confinement time and the $\tau_E^{ITER89-P}$ confinement scaling [138]. The each fitting factors $F(s, \alpha)$, F_{κ} [139] and $F_{E \times B}$ [140] is defined as

$$F(s, \alpha) = \frac{1}{\sqrt{2(1-2s')(1-2s'+3s')}} \quad (s' = s - \alpha < 0),$$

$$F(s, \alpha) = \frac{1 + 9\sqrt{2}s'^{\frac{5}{2}}}{\sqrt{2(1-2s'+3s'^2+2s'^3)}} \quad (s' = s - \alpha > 0), \quad (2.84)$$

$$F_{\kappa} = \left(\frac{2\sqrt{\kappa}}{1 + \kappa^2} \right)^{\frac{3}{2}}, \quad (2.85)$$

$$F_{E \times B} = \frac{1}{1 + \frac{\omega_{E \times B}^2}{\gamma_{CDBM}^2}}, \quad (2.86)$$

where $\omega_{E \times B} = \frac{RB_p}{B_t} \frac{d}{dr} \left(\frac{E_r}{RB_p} \right)$ is the $E \times B$ flow shearing rate and $\gamma_{CDBM} = |\alpha|^{\frac{1}{2}} \frac{v_A}{qR} F(s, \alpha)$ is the growth rate of CDBM [141]. Dependence of $F(s, \alpha) \alpha^{\frac{3}{2}}$ on s and α is shown in Fig. 2.12. $F(s, \alpha)$ decreases rapidly when $s - \alpha$ becomes smaller than 0.75. The particle diffusion coefficient is set by $D_e = D_i = C_{CDBM}^D \chi_{CDBM}$, where $C_{CDBM}^D = 1/2$ is used basically in this study.

mBgB

The mBgB model is an empirical turbulent transport model that has been developed for prediction of JET tokamaks and has been widely used for predictions of various devices [142]. This model is based on the turbulent transport caused by the mechanisms related to the presence of small-scale instabilities. Electrostatic oscillations in the plasma cause $E \times B$ drift motion due to electric field fluctuations, resulting in an oscillating component perpendicular to the magnetic field. The coupling of the density fluctuations and the electric field fluctuations generates the particle flux. Then, with a phase difference between each fluctuation and a positive growth rate, the density fluctuations can grow and eventually saturate due to non-linear effects. The diffusion coefficient at maximum saturation of the fluctuations is called the Bohm diffusion coefficient. In the case where the fluctuations

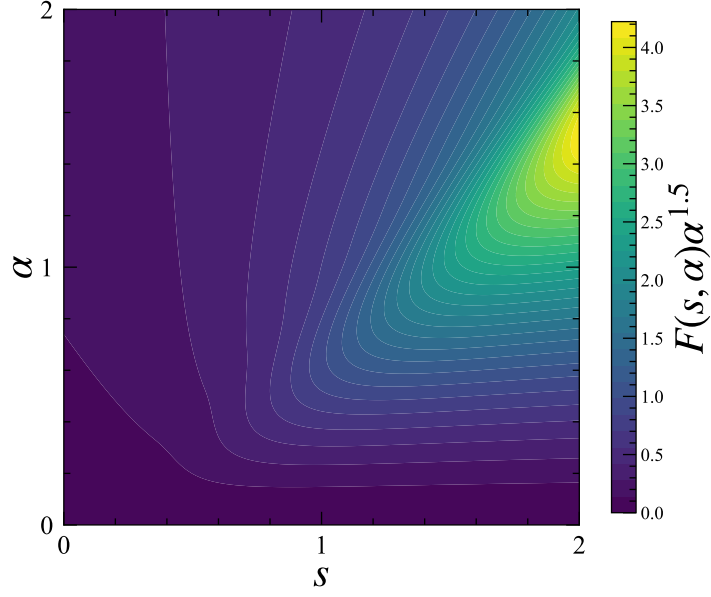


Figure 2.12: Dependence of $F(s, \alpha)\alpha^{\frac{3}{2}}$ on the magnetic shear s and the normalized pressure gradient α .

driven by the ion temperature gradient (ITG) instability have a Larmor radius scale, the diffusion coefficient is called the gyro-Bohmian diffusion coefficient [143]. When the radial spread of the oscillation is a semi-macro scale, the transport is Bohm type, and is a micro scale characterized by the Larmor radius, the transport is gyro-Bohm type. The mBGB model is a linear summation of Bohm and gyro-Bohm diffusion, and can be expressed by the following [144–146].

$$\chi_e = 8.0 \times 10^{-5} \chi_B + 3.50 \times 10^{-2} \chi_{gB} , \quad (2.87)$$

$$\chi_i = 1.6 \times 10^{-4} \chi_B + 1.75 \times 10^{-2} \chi_{gB} , \quad (2.88)$$

$$D_e = D_i = (0.3 + 0.7\rho) \frac{\chi_e \chi_i}{\chi_e + \chi_i} . \quad (2.89)$$

The impurity particle diffusion coefficient is $D_Z = D_i$. The Bohm term χ_B and gyro-Bohm term χ_{gB} are given by

$$\chi_B = \rho_s c_s q^2 \frac{a(dp_e/dr)}{p_e} \frac{T_e^{\rho=0.8} - T_e^{\rho=1.0}}{T_e^{\rho=1.0}} \Theta \left(-0.14 + s - \left| \frac{1.47 \omega_{E \times B}}{\gamma_{ITG}} \right| \right) , \quad (2.90)$$

$$\chi_{gB} = \frac{\rho_s^2 c_s}{a} \frac{a(dT_e/dr)}{T_e} , \quad (2.91)$$

where, $\rho_s = \frac{c_s m_i}{e B_t}$ is the gyro radius, $c_s = \sqrt{\frac{kT_e}{m_i}}$ is the sound speed, p_e is the electron pressure, Θ is the Heaviside function, $\omega_{E \times B}$ is the $E \times B$ flow shear rate and $\gamma_{ITG} = \frac{v_{th,i}}{qR}$ is the ITG growth rate.

MMM95

The MMM95 model is a theory-based model represented by a linear summation of the ITG mode and TEM of the Weiland model, the resistive ballooning (RB) mode of the Guzdar-Drake model, and the kinetic ballooning (KB) mode of the semi-empirical model [147]. In the Weiland model, the eigenvalue problem for the linearized fluid equations is solved and the growth rate and frequency of the unstable modes are evaluated as shown in Fig. 2.13. The mode frequency is TEM for positive value and ITG mode for negative value. The transport coefficients evaluated by the MMM95 model generally has a significant contribution from the transport by the Weiland model. The Weiland model is a stiff transport model, and the diffusion coefficient increases sharply as the normalized temperature gradient increases. The resistive and kinetic ballooning modes increase with pressure gradient and collisions. The resistive ballooning mode has a more significant contribution at the plasma edge. In contrast, the kinetic ballooning mode has a more negligible contribution over the entire plasma. The diffusion coefficients of the MMM95 model are given by the following:

$$\chi_e = 0.80\chi_e^{ITG,TEM} + 1.00\chi_e^{RB} + 1.00\chi_e^{KB}, \quad (2.92)$$

$$\chi_i = 0.80\chi_i^{ITG,TEM} + 0.65\chi_i^{RB} + 1.00\chi_i^{KB}, \quad (2.93)$$

$$D_H = 0.80\chi_H^{ITG,TEM} + 0.65\chi_H^{RB} + 1.00\chi_H^{KB}, \quad (2.94)$$

$$D_Z = 0.80\chi_Z^{ITG,TEM} + 1.00\chi_Z^{RB} + 1.00\chi_Z^{KB}. \quad (2.95)$$

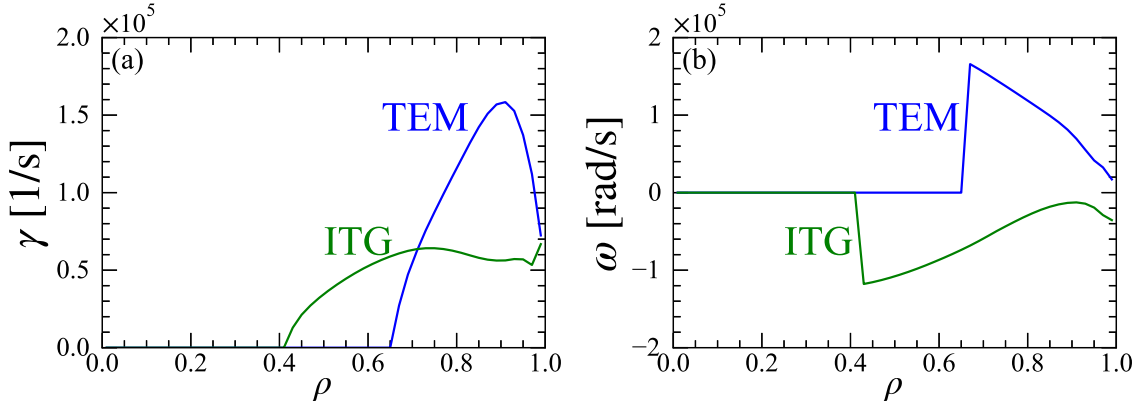


Figure 2.13: Radial profiles of (a) linear growth rates and (b) mode frequencies for the ITG and TEM instabilities. The results evaluated by the Weiland model in MMM95 model are shown for the PLATO plasma parameters.

2.4.6 Pedestal profile model

A H-mode pedestal profile is given by a modified tanh function expressed as [148]

$$X(\rho) = \frac{X_{ped} - X_{sol}}{2} \left(\text{mtanh} \left(\frac{\rho_{mid} - \rho}{2\rho_w}, \alpha_{ped} \right) \right) + X_{sol}, \quad (2.96)$$

$$\text{mtanh} = \frac{(1 + \alpha_{ped}\rho)e^\rho - e^{-\rho}}{e^\rho + e^{-\rho}}. \quad (2.97)$$

Here, X represents each plasma quantity, and X_{ped} , X_{sol} , ρ_{mid} , ρ_w and α_{ped} are the height of pedestal, the height of scrape-off layer, the middle position of pedestal, the radial width of pedestal, the radial profile fitting parameter, respectively. The definition of each variable is shown in Fig. 2.14.

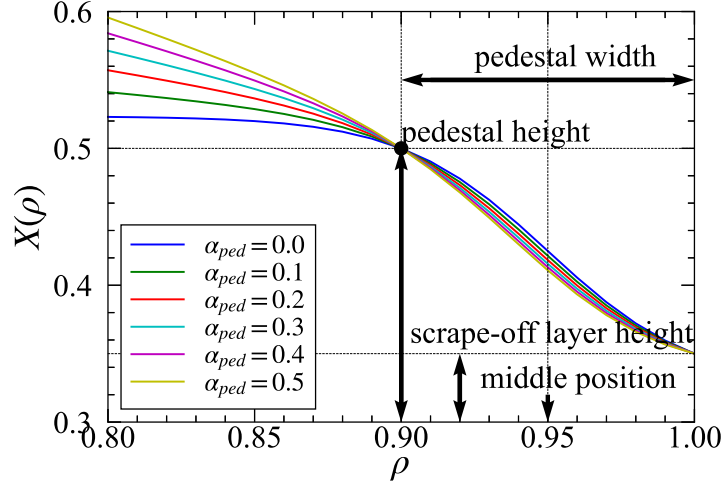


Figure 2.14: Pedestal profiles and parameter variable definitions of the modified tanh function.

2.4.7 Modeling of time evolution

As shown in Sec. 2.2, the several modules and the physical models contained within modules are combined to describe the time evolution of the plasma and impurity. To show that the modeling of the time evolution is appropriately performed, the dynamic response of the impurity and the plasma is shown in Fig. 2.15. At $t = 0 - 0.5$ s, tungsten accumulates in the plasma center region due to the inward neoclassical particle pinch driven by the main ion density gradient, resulting in increased radiation losses. On-axis heating of electrons and ions at $t = 0.5$ s increases the plasma temperature and decreases the plasma density, as shown in Figs. 2.15(f) and (g). These variations of the plasma profiles increase the outward neoclassical pinch, and the W accumulation and radiation losses are decreased at the plasma center. Thus, a simulation scheme has been developed to determine the profiles of the plasma, impurity and other quantities, including time evolution.

2.5 Summary of the chapter

A plasma transport simulation scheme including the impurity is being developed in the integrated code TASK. The 1.5D transport simulation combines the two-dimensional MHD equilibrium analysis code, the one-dimensional plasma transport code, and the one-dimensional impurity transport code. For the MHD equilibrium analysis, two methods are available: one is fixed the plasma shape and the other determines the plasma shape with the external coil current conditions. For the plasma transport analysis, the radial profiles are determined by combining transport coefficients, particle sources, and heating sources. For the impurity transport analysis, the radial profiles are determined considering

the ionization state using the atomic process database. In this study, geometric coefficients P_A/P_B are introduced into the neoclassical transport code NCLASS to reproduce the poloidal asymmetry of the high-Z impurity caused by the toroidally rotating plasma due to the NBI heating. The model considering the toroidal rotation effect will be validated by using the PHZ/ E_r pinch model, an analytical representation of the transport by the toroidal rotation of the high-Z impurity. Using these modules and models, the integrated transport simulation scheme can be used to predict the time evolution of the plasma and impurity.

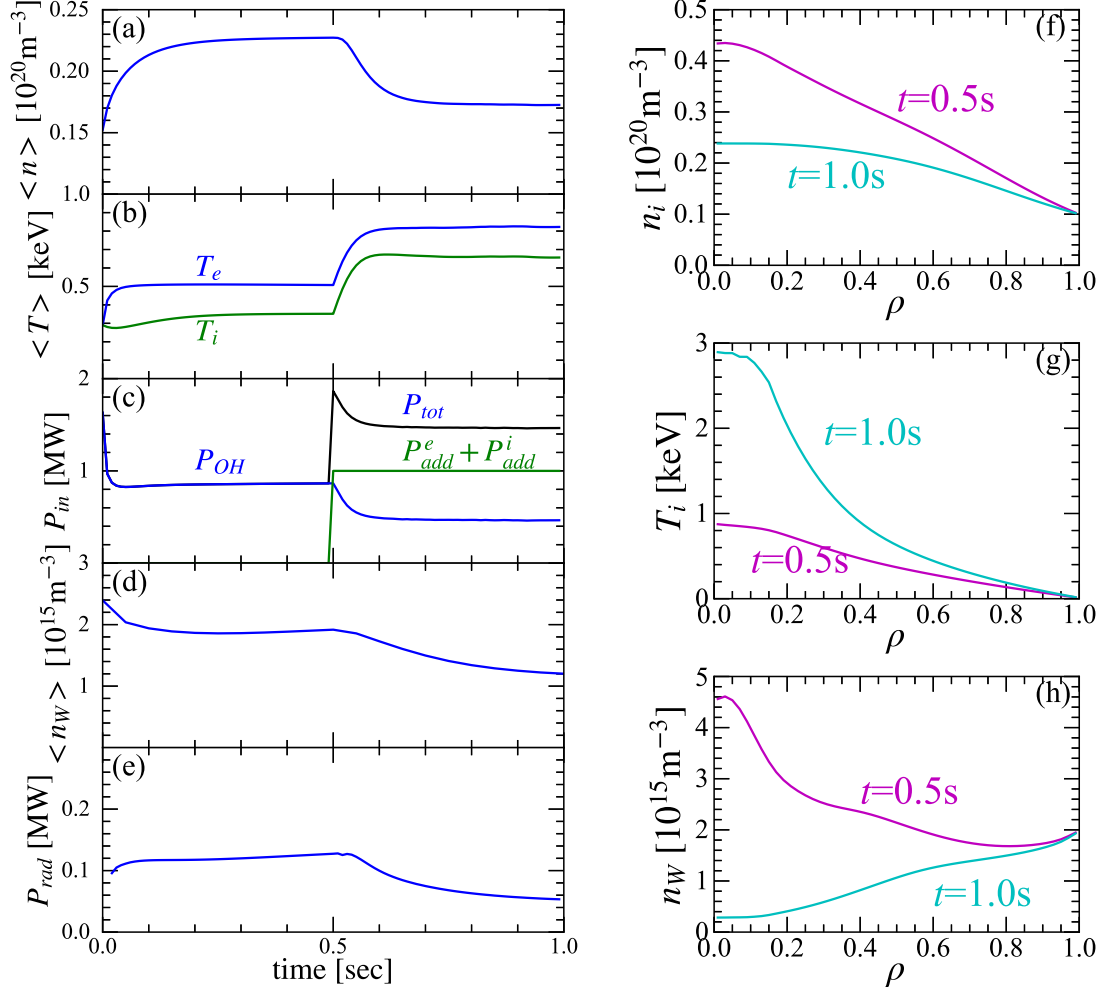


Figure 2.15: Simulation of plasma dynamics with the TASK code. The left column shows the time evolutions of the plasma and impurity, from top to bottom: (a) volume-averaged density, (b) volume-averaged temperature, (c) input power, (d) volume-averaged tungsten density, and (e) radiation loss power. The right column shows the radial profiles of (a) ion density, (b) ion temperature, and (c) tungsten density. The ASDEX-UPGRADE-like plasma parameters ($R = 1.65$ m, $a = 0.5$ m, $I_p = 0.8$ MA, $B_t = 2.0$ T, $T_e = 3.5$ keV) and the W density at the boundary $n_{W,edge} = 2.0 \times 10^{15} \text{ m}^{-3}$ are used. Additional on-axis heating $P_{add}^e = P_{add}^i = 0.5$ MW at $t = 0.5$ s are applied for impurity control.

Chapter 3

Transport simulation in L-mode plasma

This chapter presents the results of applying the transport simulation scheme to L-mode plasmas. The target device here is the plasma experimental device PLATO (PLAsma Turbulence Observatory), and the experiment is currently in progress at the Research Institute for Applied Mechanics, Kyushu University [149]. Figure 3.1(a) shows a photograph of the device. PLATO is equipped with a large number of instruments to measure the plasma locally and globally at the same time for exploring the unresolved problems of turbulence. The vacuum vessel is designed to have a high proximity for accurate turbulence measurement. There are two types of vacuum vessels, short-type and long-type, each with four pieces, and the vacuum vessel is modularly divided into eight parts in total. The vacuum vessel is equipped with observation ports for tomography and the heavy ion probe (HIBP). A global observations will be performed by using tomography systems with more than 1500 detectors to observe microscopic turbulent structures. The system has successfully visualized the structure of turbulent oscillations using a tomography system in the linear plasma experiment PANTA at Kyushu University [150, 151]. For profile measurement, HIBP is used to simultaneously observe the fluctuations of the density, potential, and magnetic field inside the plasma [152]. The HIBP is an excellent fluctuation measurement system, as it has succeeded in identifying zonal flows for the first time in the world [153]. In PLATO, multiple HIBPs will be installed to enable more detailed local observations and simultaneous measurement of the overall structure. In PLATO, the gas puffing method is used to supply particles. A gas valve for puffing is installed between the gas cylinder and the vacuum vessel. Hydrogen gas is injected into the vacuum vessel, and the ionized gas becomes the fuel for plasma discharge. As shown in Fig. 3.1(b), the PLATO device has three main pairs of vertical field coils: D (dipole), Q (quadropole), and H (hexapole) coils, and two pairs of feedback coils, F1 and F2 coils. The ohmic heating coil in the center solenoid is divided into two pairs: OH and OHC coils [154]. The PLATO project is a collaboration of experimental, theoretical, and simulation research groups, and numerical simulation research is performed earlier in the project [155, 156].

In this chapter, transport analysis using the integrated simulation code TASK is performed on PLATO tokamak to predict plasma performance. The transport analysis is performed under the assumption that plasma is in the MHD equilibrium state. For transport analysis considering the experimental conditions, MHD equilibria in the simulations are calculated by TASK/EQU, which considers the current conditions of the external coils. The transport simulation scheme in the TASK code is used to predict the plasma performance and the effects of impurities on the main plasma.

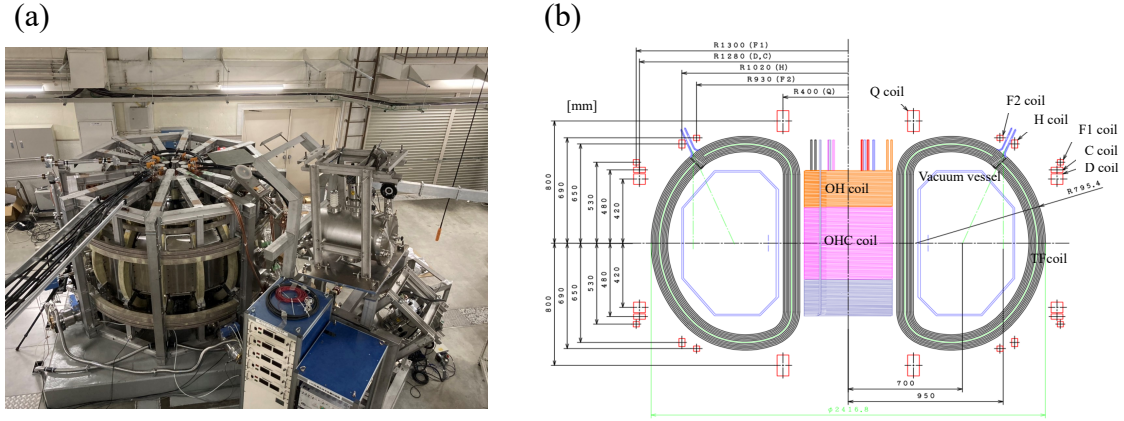


Figure 3.1: (a) Photograph and (b) sectional view of PLATO device.

3.1 Prediction of plasma performance with PLATO parameter

The transport simulation scheme (EQU-TR combination) is used to evaluate the plasma performance and transport property with PLATO plasma parameters. The toroidal magnetic field $B_t = 0.3$ T and plasma current $I_p = 40$ kA are kept to be constant. In this calculation, the following coil currents in table 3.1 are used with non-zero D, Q and H coil currents.

Table 3.1: Vertical coil currents used in the simulation. The value is positive when the direction of the current is the same as that of the plasma current. Max represents the maximum values of the coil currents.

Coil name	Max coil current [kA·turn]	Coil current [kA·turn]
OH	-134.0	-67.2
OHC	-134.0	-67.2
D	38.0	-23.3
Q	76.0	-10.8
H	14.0	0.0
C	-	0.0
F1	9.0	0.0
F2	9.0	0.0

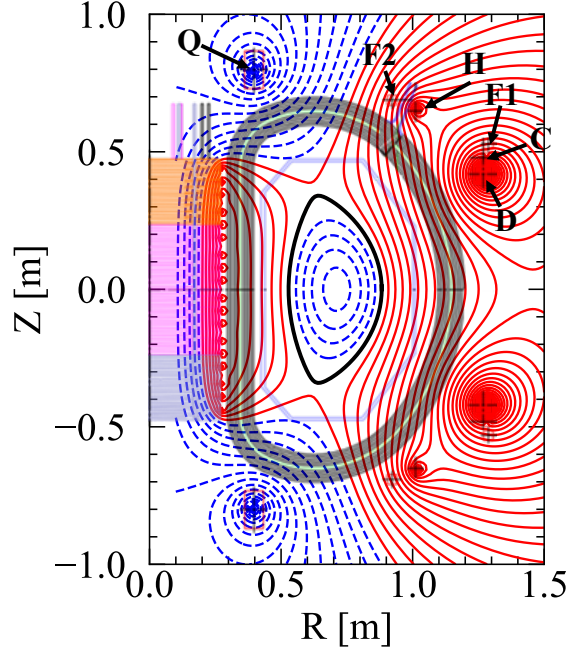


Figure 3.2: Equilibrium magnetic configuration of PLATO tokamak obtained by TASK/EQU module. The position of each vertical field coils and the ohmic coils are also shown.

Figure 3.2 shows the poloidal cross-section of the flux contours of the equilibrium with EQU calculation. The obtained equilibrium has $R = 0.7$ m, $a = 0.18$ m, ellipticity $\kappa = 1.9$, and triangularity $\delta = 0.4$. In the plasma transport simulations, the CDBM, mBgB, and MMM95 models are used as turbulent transport models in the equilibrium configuration shown in Fig. 3.2. The particle diffusion coefficient in the CDBM model is assumed to be 1/10 times the thermal diffusion coefficient. The density at the plasma center is set to be $n_0 = 1.0 \times 10^{19} \text{ m}^{-3}$ by supplying electrons and hydrogen ions as the particle source as shown in Fig. 3.3(a). The gas-puffing method is the main particle supply in PLATO, and this source profile peaks near the edge. The particle source profile is given by Eq. 2.36, with $\rho_0^{S_p} = 0.7$ and $\rho_w^{S_p} = 0.3$. The plasma profiles in the steady-state are obtained with TR calculation as shown in Fig. 3.3. The obtained temperature has $T_e \sim 200 - 300$ eV and $T_i \sim 45 - 100$ eV at the plasma center. The heating mechanism of PLATO plasmas are only OH heating and collisional heat exchange, both of which are collision-dependent mechanisms. Therefore, the collision and plasma resistivity determine the heating profiles, and strongly depend on the plasma density and temperature profiles. As shown in Fig. 1.7, the collisions of the ions are in the banana-plateau regime over a broad region, and in the PS regime at the plasma periphery. Transport coefficients also strongly depend on the plasma profiles and vary with the transport model. The steady-state profiles are determined from the balance between the transport and source terms based on the transport equation. As shown in Fig. 3.4, each turbulent model has the dependency on pressure gradients, so the diffusion coefficient decreases at the plasma central region where the density and temperature gradients are small. In this simulation, the neoclassical convection directs inward, but its contribution to the transport process is small, and the neoclassical diffusion also contributes less than the turbulent diffusion, except in the plasma center. In the case of the mBgB model, the Bohm diffusion term is dominant over the gyro-Bohm diffusion term in the entire plasma. As shown in Eq. 2.90, the Bohm term in

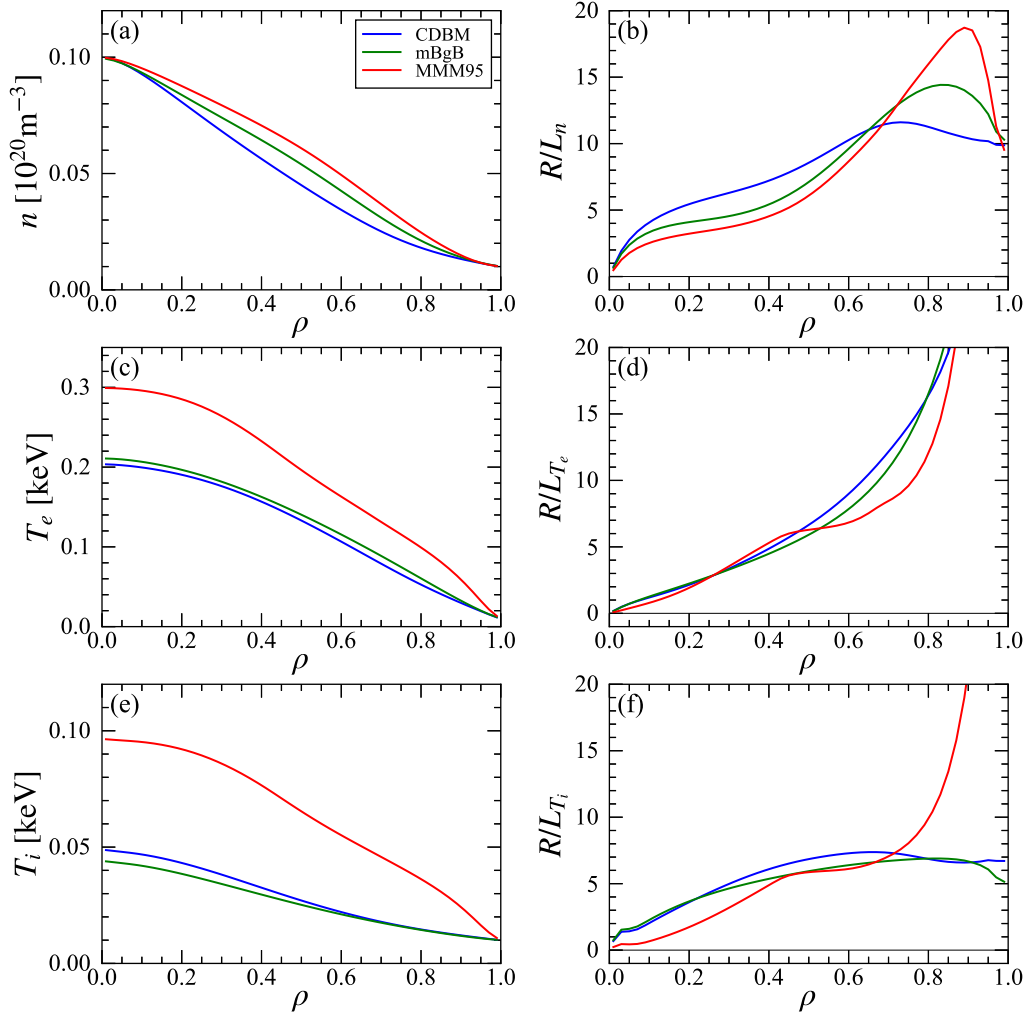


Figure 3.3: Radial profiles of (a) density, (b) normalized density gradient, (c) electron temperature, (d) normalized electron temperature gradient, (e) ion temperature, (f) normalized ion temperature gradient. The cases with CDBM, mBgB and MMM95 are shown.

the mBgB model depends on the electron temperature at the plasma edge to determine the magnitude of the overall plasma transport, and thus the contribution of Bohm scaling is significant. In the case of the MMM95 model, the resistive ballooning mode is dominant in the plasma periphery and decreases in the plasma center region. The magnitude of kinetic ballooning mode is small because of the low- β value, and even in the plasma center, the contribution is about the same as that of the resistive ballooning mode. The ITG/TEM mode exists for $\rho \sim 0.6 - 0.9$, where the plasma gradient is large. The thermal diffusion coefficient of the ITG/TEM mode increases to the same order of magnitude as that of resistive ballooning mode, which is different from the particle diffusion coefficient. Similar to this simulation, the ITG/TEM mode component does not increase the particle diffusion coefficient in the calculation using MMM95 for the similar plasma parameters as in PLATO [157]. It is noted that the MMM95 transport model depends fully on local plasma parameters, so unlike the mBgB model, the gyro-Bohm scaling is effective [158]. For the

entire region, the thermal diffusion coefficient of MMM95 is small for both electrons and ions, resulting in a higher plasma temperature than other models. Especially, as shown in Fig. 3.3(f), the ion temperature profiles are flattened at the plasma periphery for the CDBM and mBgB model, where the thermal diffusion coefficients are increased. These significant differences in the ion temperature profiles are due to the absence of an external heating mechanism for the ions, which results in the dominant transport process. The global plasma performance values for the CDBM, mBgB, and MMM95 cases, in that order, have normalized β value $\beta_N = 0.35, 0.40, 0.71$, and the confinement enhancement factor of ITER89-P scaling $H_{89L} = \tau_E / \tau_E^{ITER89-P} = 0.43, 0.53, 1.29$. The magnitude of the transport coefficients significantly affects the plasma performance. In this prediction by simulation, the CDBM and mBgB models predict similar plasma performance, while the MMM95 model overestimates the temperature. Compared to experiments, it is noted that the Weiland model, which is included in the MMM95 model, tends to overestimate the electron temperature in L-mode plasmas, while the CDBM model reproduces the temperature profile well [139]. In addition, since the plasma performance in actual experiments can be degraded more than predicted here, the CDBM is used as the turbulent transport model in the following simulations. The PLATO aims to observe turbulence, and a method to increase the ion temperature is necessary to observe various types of turbulence. The dependences of plasma performance on the externally controllable quantities are evaluated in appendix A.1 and appendix A.2.

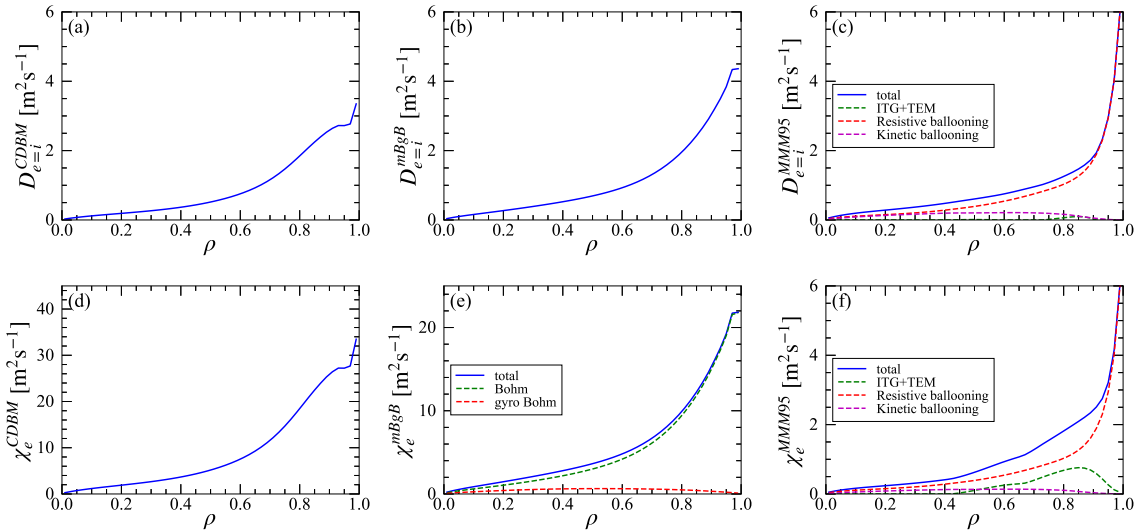


Figure 3.4: Radial profiles of particle diffusion coefficient of turbulence in the case with (a) CDBM, (b) mBgB, (c) MMM95, and radial profiles of thermal diffusion coefficient of turbulence in the case with (d) CDBM, (e) mBgB, (f) MMM95.

3.2 Prediction of impurity impacts

In this section, the effects of impurity contamination are investigated using the same calculation conditions of plasma transport as in Sec. 3.1. Transport simulations are performed for carbon (C), oxygen (O), and iron (Fe) as impurity species. Neoclassical

transport coefficients of the impurity species are evaluated using NCLASS without considering toroidal rotation effects. For turbulent transport of impurities, only diffusion coefficients are given and convective velocities are neglected. The turbulent diffusion coefficient profiles in this calculation are given by $D_Z^{turb} = D_Z^{\rho=0} + (D_Z^{\rho=1} - D_Z^{\rho=0})\rho^6$, where $D_Z^{\rho=0} = 0.05$ and $D_Z^{\rho=1} = 10$ set with reference to the analysis of the Aditya tokamak, which has the similar plasma parameters as PLATO [159]. The impurity densities at the plasma boundary are given to be fixed.

The plasma performance is evaluated for the amount of impurity contamination by scanning the impurity density at $\rho = 1.0$ (boundary condition), as shown in Fig. 3.5. With increasing impurity concentration, the electron temperature is decreased, and the electron density is increased. Heavy impurities cannot exist in large quantities in the plasma. This is because the electron temperature becomes negative due to increased radiation loss caused by impurity accumulation, and the calculation breaks down. While Fe affects the main plasma by impurity radiation loss at $c_Z = \langle n_Z \rangle / \langle n_e \rangle \sim 10^{-4}$, O and C affect for $c_Z \geq 10^{-2}$. In addition, the calculation breaks down at a higher volume-averaged electron temperature for Fe than for the other particle species. This is a result of impurity profile effects.

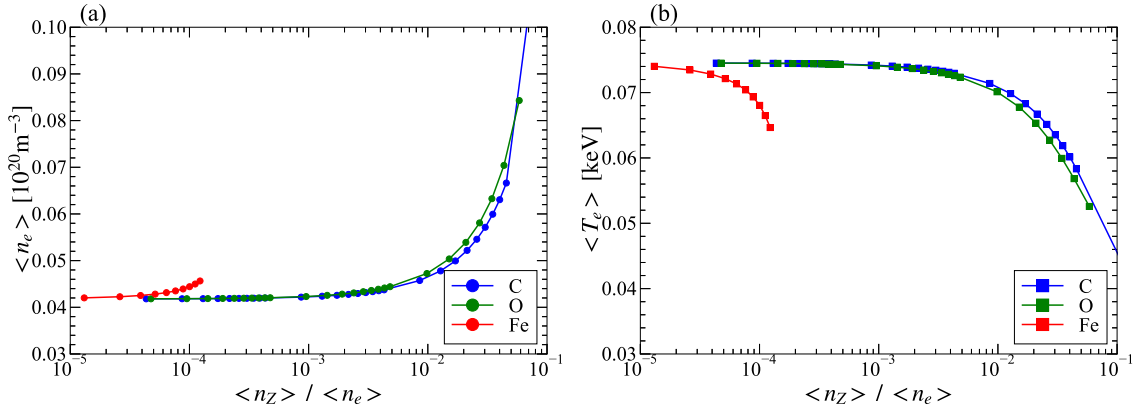


Figure 3.5: Dependence of volume-averaged (a) electron density and (b) electron temperature when impurities exist with concentration $c_Z = \langle n_Z \rangle / \langle n_e \rangle$ in the case of impurity C, O and Fe.

Figure 3.6 shows the radial profile of the density and radiation loss power for each impurity species with the boundary density $n_{Z,edge} = 1.0 \times 10^{15} \text{ m}^{-3}$. The carbon ion is dominated by C^{5+} in the plasma center region, and its contribution to radiation loss is also large. In the entire plasma region, the contribution to radiation loss is dominated by C^{4+} , which has a peak in density profile at $\rho \sim 0.5$. The oxygen ion is dominated by O^{6+} in the plasma center region, and its contribution to radiation loss is also large over a wide region. At the plasma edge, the radiation loss due to O^{3+} increases and a strong peak is formed. The iron ion is dominated by Fe^{13+} at the plasma center region, thus the iron is ionized to about half ($Z_{max} = 26$). The radiation loss power increases with increasing the impurity accumulation in the plasma center for the heavier impurity species. Due to the high radiation loss power in the center of the plasma, it exceeds the heating power locally. Even if the impurity accumulation is small in the entire plasma, the electron temperature can become negative in the case of Fe, which can cause the calculation to break down. The mechanism driving the impurity accumulation is the inward neoclassical particle pinch. As shown in Fig. 3.7(a), the impurity

diffusion coefficients are the same value for all species for $\rho = 0.6 - 1.0$. However, in the central region of the plasma, the diffusion coefficient of Fe is about twice larger than that of other particle species due to the increased collision frequency with increasing the density. On the other hand, as shown in Fig. 3.7(b), the Fe convection velocity increases inwardly for the entire plasma region. Therefore, particle species with higher masses and higher charge number easily accumulate in the plasma center, and damage to materials such as Fe should be avoided as much as possible. The Ware pinch included in the V^{NC} shown in Fig. 3.7 is constantly inward direction, with $V_C^{Ware} \sim -0.6$, $V_O^{Ware} \sim -0.9$, and $V_{Fe}^{Ware} \sim -1.3$ at large values of $\rho = 0.8$, contributing less than 5% to the overall convective velocity.

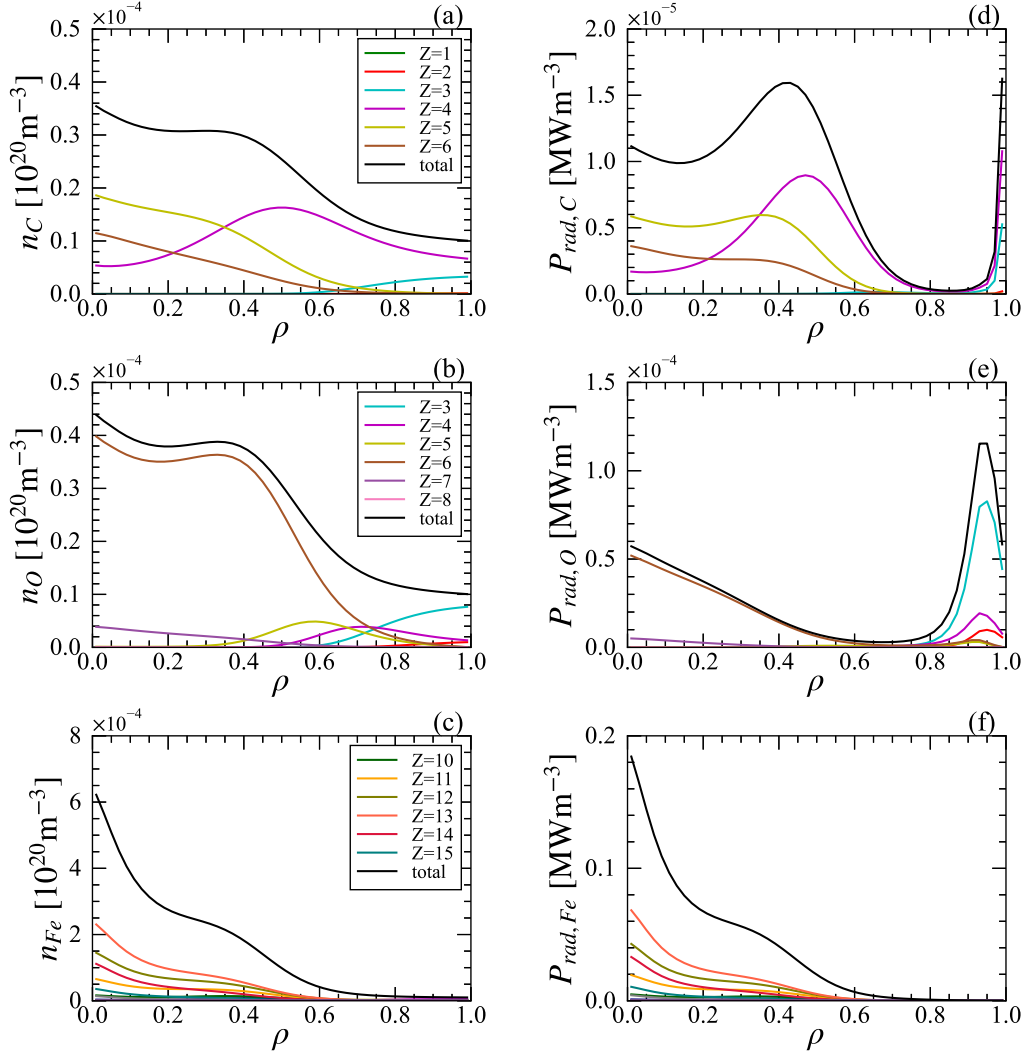


Figure 3.6: Radial profiles of density of (a) C, (b) O, (c) Fe, and radiation loss power of (d) C, (e) O, (f) Fe. The densities at boundary are $n_{Z,edge} = 1.0 \times 10^{15} \text{ m}^{-3}$ for all impurity species.

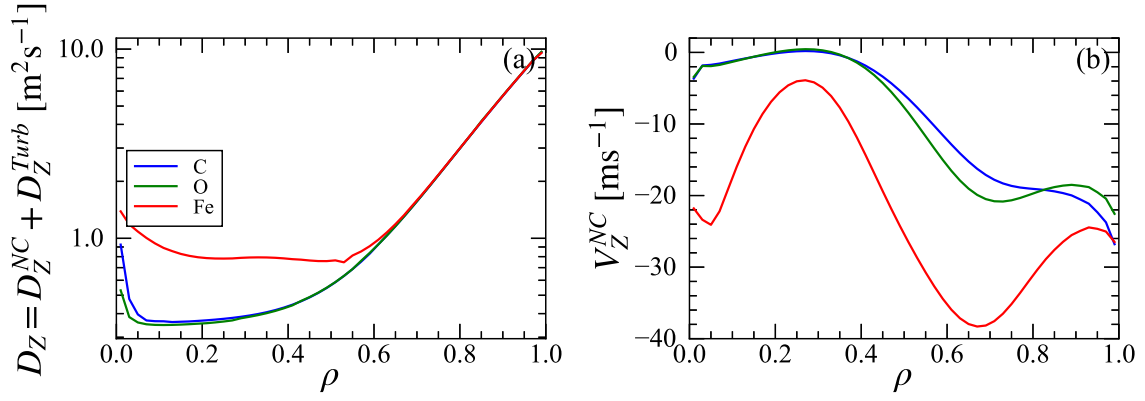


Figure 3.7: Radial profiles of (a) diffusion coefficient and (b) neoclassical convection velocity for C, O, and Fe, where the label 'NC' denotes the neoclassical transport coefficient.

3.3 Summary of the chapter

In this chapter, simulations are performed for low confinement plasmas as a first step in predicting improved confinement mode plasmas. The transport analyses are performed with calculation of free-boundary condition equilibria, considering external vertical field coils in PLATO experiments. Plasma performance is predicted in PLATO using the CDBM, mBgB, and MMM95 turbulence models. The electron temperature is obtained $T_e \sim 200 - 300\text{eV}$ at the plasma center, and the ion temperature is obtained $T_i \sim 45 - 100\text{eV}$ at the plasma center. Heat transport in L-mode discharges is typically characterized by Bohm transport, which is a type of turbulent transport occurring on a larger scale than gyro-Bohm transport. The current simulation using the mBgB model have showed a significant contribution from Bohm diffusion. Additionally, the CDBM model has predicted diffusion coefficients that are comparable in magnitude to the Bohm diffusion coefficients. On the other hand, the Weiland model included in the MMM95 model evaluates transport coefficients based on microscopic instabilities, which are influenced by local plasma conditions. As a result, the Weiland model is considered to have evaluated to be smaller diffusion coefficients compared to the Bohm diffusion coefficient. The effect of including impurities in the predicted PLATO plasmas is also investigated. The PLATO plasmas with small ion temperature do not drive temperature screening in the impurity neoclassical transport, resulting in the inward particle pinch due to ion density gradients and the accumulation in the plasma center region. Even if the densities of carbon and oxygen increase in the plasma, the effect of the radiation loss is small. In the case of iron, the electron temperature decreased sharply in the plasma center when the iron density is increased above 10^{-4} with respect to the electron density. In order to suppress the effects of impurities, it is necessary to avoid damaging the material, to perform complete vacuum evacuation, and to control the plasma profiles to decrease impurity transport. As shown in the appendix A.1 and A.2, the plasma performance and the possibility of controlling the plasma profiles are investigated. There no direct ion heating method in the initial phase of PLATO experiments, so the indirect methods are proposed to increase the ion temperature. The method of increasing ion temperature could contribute to the understanding of turbulence and could contribute to the enhancement of the neoclassical temperature screening effect of impurities.

Chapter 4

Transport simulation in improved confinement mode plasma

In this chapter, transport simulations including the impurity are performed for improved confinement mode plasmas. In the previous chapter, the simulations are performed for L-mode plasmas. In contrast, in this chapter, simulations are performed using several MW of plasma heating and plasma current control to form ITBs in the plasma profile. Simulation results are shown to reproduce the improved confinement mode as observed in experiments by using the neoclassical and turbulent transport models in the TASK code. Simulations of tungsten transport as impurity species are performed simultaneously with plasma transport. Impurity accumulation in the ITB forming plasma has been experimentally confirmed [160–165]. Here, the impurity transport simulations focus on neoclassical transport, which is considered to accelerate the impurity accumulation in the ITB region. In addition, the plasma with ITBs forming transiently is investigated, and the mechanism of impurity profile formation is clarified.

Simulations of high- β_p mode and negative magnetic shear mode are performed with typical plasma parameters for a JET-like tokamak; the major radius $R = 3.0$ m, minor radius $a = 1.2$ m, ellipticity $\kappa = 1.5$, magnetic field $B_t = 3.0$ T. As described below, different values of plasma current are used in case of the high- β_p mode and negative magnetic shear mode. The on-axis heating power $P_H = 5$ MW for electrons and ions are used Eq. 2.16, the radial position $\rho_0^H = 0.0$ and the width $\rho_w^H = 0.2$. For simplicity, the additional heating is pure heating without external current drive. The particle sources for electrons and ions are used Eq. 2.36, with the radial position $\rho_0^{S_p} = 1.0$ and the width $\rho_w^{S_p} = 0.7$. The same profiles of heating and particle sources are used in case of high- β_p mode and negative magnetic shear mode. For the transport model, the NCLASS is used for neoclassical transport, and the CDBM model is used for turbulent transport in these calculations. The CDBM model is known to reproduce well the ITBs of the plasma temperature profile [138]. The turbulent diffusion coefficient for particles is half of the CDBM thermal diffusion coefficient, and the particle diffusion coefficient for impurity ions is the same value as for bulk ions. Poloidal asymmetries are not considered in impurity transport calculations. The boundary condition for tungsten impurity is used $\Gamma_{w,in} = 1 \times 10^{13} \text{ m}^{-2}\text{s}^{-1}$, which is the inflow condition where the effects of radiation loss are negligible due to the low level of tungsten density.

4.1 High- β_p mode

Figure 4.1 shows the time evolution of the high- β_p mode plasma. The plasma current is steady to be 0.6 MA, which is a relatively small value. In this calculation, the additional heating of 5 MW is applied at $t = 1.0$ s for electrons and ions. As shown in Fig. 4.1(a), the bootstrap current increases with plasma heating, and peaks at $t \sim 1.0$ s, and then decreases to a steady state. The bootstrap current increases in the plasma center region, where the external heating is stronger, as shown in Fig. 4.2(a). The safety factor profiles have $q_0 > 1$ at the plasma center due to the contribution of the bootstrap current (Fig. 4.2(b)). Note that the plasma current is decreased significantly in this calculation, so the safety factor increases in the entire radial direction. The bootstrap current drive ratio $f_{BS} \sim I_{BS}/I_p = 0.47$ is obtained at $t = 1.1$ s, and $f_{BS} = 0.35$ is obtained at $t = 2.0$ s, so the bootstrap current is maintained at the steady state. The OH heating is sufficiently small compared to the additional heating, and the radiation loss of tungsten is also sufficiently small.

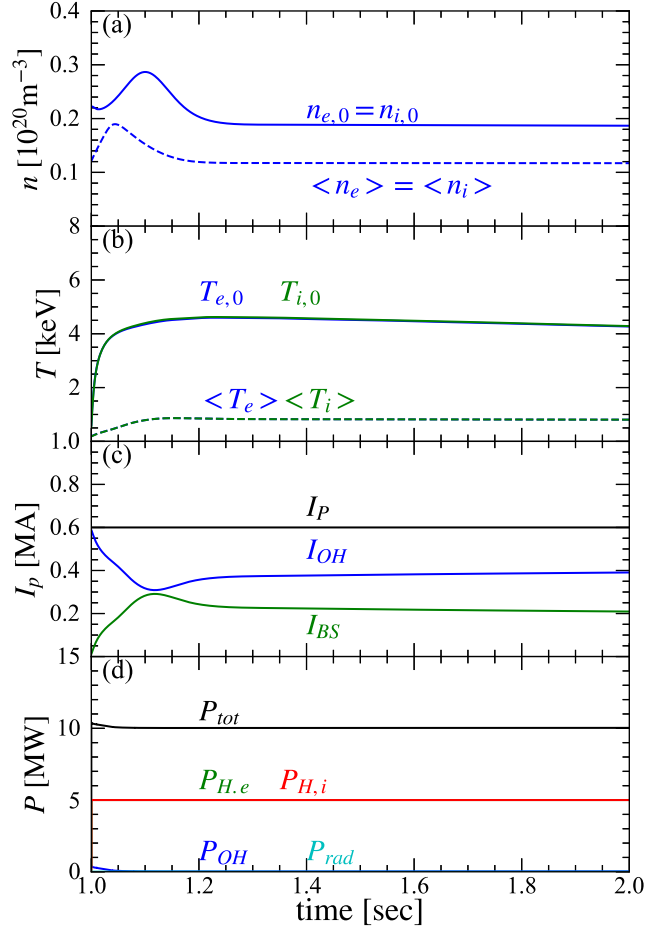


Figure 4.1: Time evolution of (a) volume-averaged and central plasma density, (b) volume-averaged and central temperatures of electron and ion, (c) plasma current, (d) heating power.

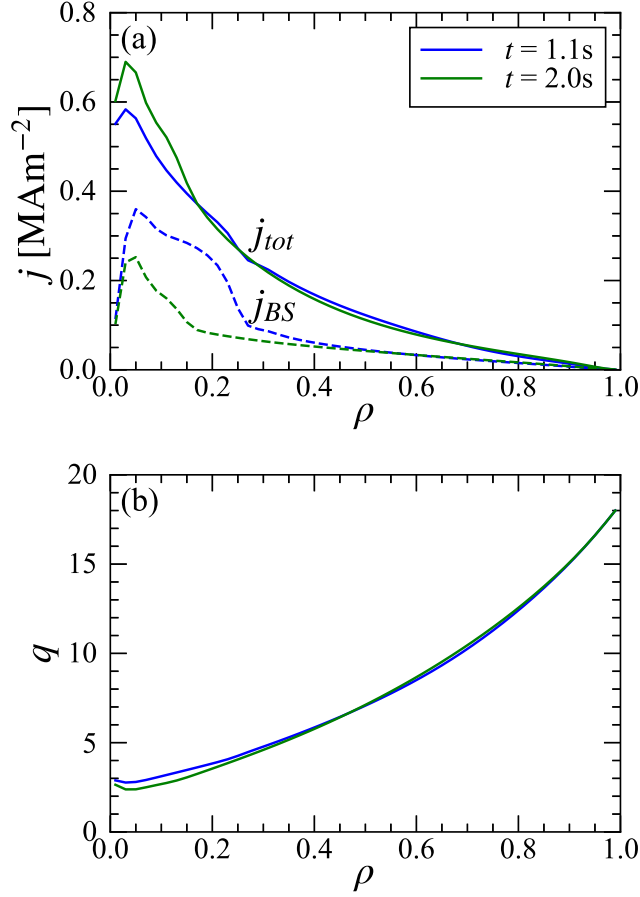


Figure 4.2: Time evolution of radial profiles of (a) current density and (b) safety factor at $t = 1.1$ s and 2.0 s. The bootstrap current density profiles are plotted with dashed lines in figure (a).

Figure 4.3 shows the time evolution of the plasma profiles. At $t = 1.1$ s, the plasma density increases in the entire radial region, and the ion temperature increases in the plasma center region. In particular, the ion temperature increases sharply inside $\rho \sim 0.25$, which is called the ITB-foot. In this ITB region, the turbulent diffusion coefficients are locally decreased as shown in Fig. 4.3(c). At $t = 2.0$ s, the ion temperature gradient increases inside $\rho \sim 0.15$, which is the improved transport region as well as at $t = 1.1$ s. The turbulent transport is suppressed by the increased normalized pressure gradient α due to the strong on-axis heating, and by the decreased magnetic shear s at the plasma center region. In the CDBM model, the transport coefficient decreases as $F(s, \alpha)$ decreases significantly for $0.5 < s - \alpha < 0.75$. The poloidal β value is $\beta_p = 1.14$ at $t = 1.1$ s, when the ITB is enhanced, while $\beta_p = 1.45$ at $t = 2.0$ s, when the ITB is relaxed. The β value and f_{BS} reach a maximum at $t = 1.1$ s and then decrease, corresponding to the evolution of the plasma density. Until $t = 1.1$ s, fuel particles flow into the center of the plasma by the source set at the plasma periphery, but as the plasma temperature increases, the particles are ejected outside the plasma, and then the plasma density decreases. Unlike the temperature profile, since the particle source does not exist in the center of the plasma, the density profile does not form a strong gradient and the central density is not maintained.

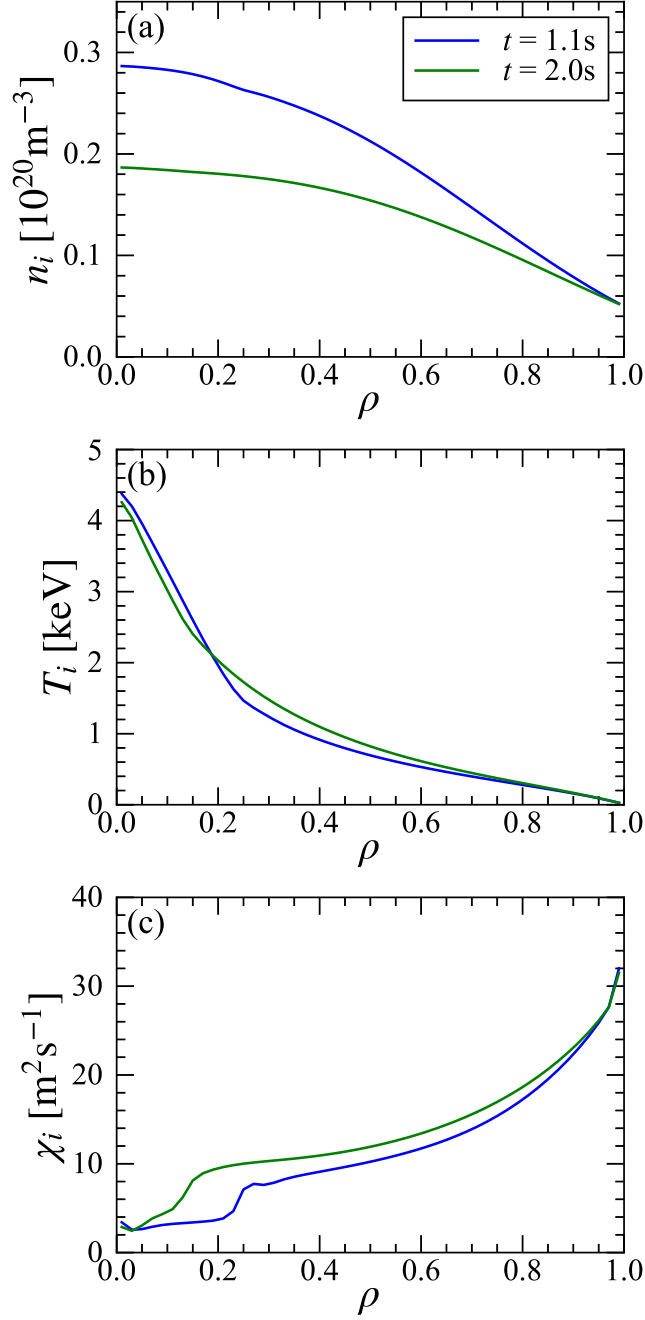


Figure 4.3: Time evolution of radial profiles of (a) ion densities, (b) ion temperatures, (c) ion thermal diffusion coefficients at $t = 1.1$ s and 2.0 s.

Figure 4.4 shows the radial profiles of the W density and neoclassical convection velocity. The W density is constant over a broad region and decreases in the plasma center region. As mentioned above, the ion temperature gradient is increasing in the ITB region, but the ion density gradient is decreased due to the absence of particle supply. Especially at $t = 1.1$ s, the normalized ion density gradient is $R/L_{n_i} \sim 1.5$ for $\rho = 0.2$, while the normalized ion temperature gradient is $R/L_{T_i} \sim 15.0$. Therefore, as shown in Fig. 4.4(b), the outward neoclassical convection is driven, and turbulent diffusion is decreased, resulting in decreased W density. The outward neoclassical convection is driven at $t = 2.0$ s outside the ITB,

whereas the inward neoclassical convection is driven at $t = 1.1$ s at $0.30 < \rho < 0.75$. This results from the relatively increased ion density at $t = 1.1$ s, and the ion density gradient increase at the plasma periphery. At $\rho = 0.6$, the normalized gradient is $R/L_{n_i} \sim 4.6$, whereas $R/L_{T_i} \sim 6.9$. Note, however, that outside the ITB region, the W density has a small radial dependence due to the increased turbulent diffusion.

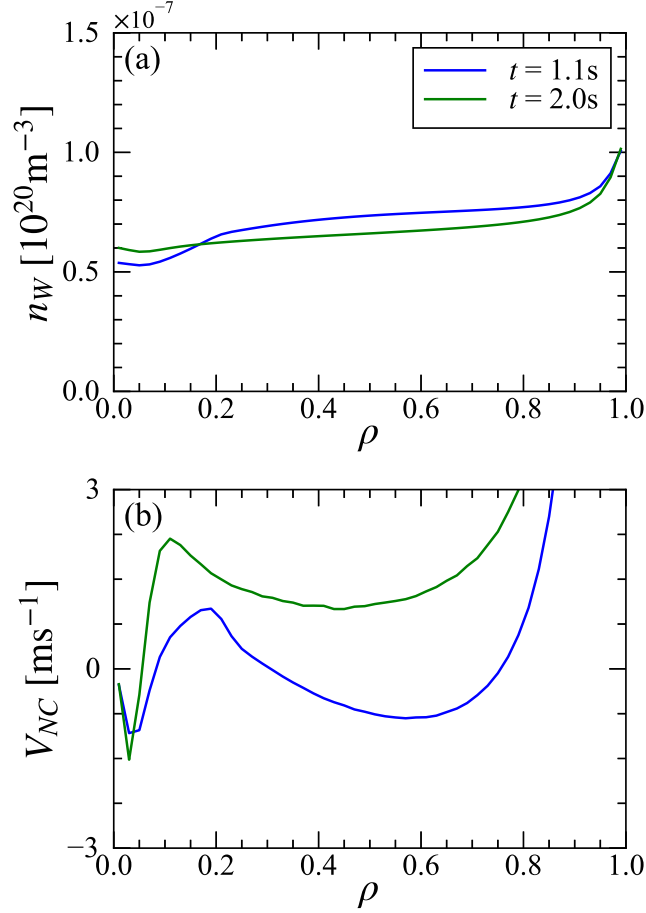


Figure 4.4: Time evolution of radial profiles of (a) W densities and (b) W neoclassical convective velocities at $t = 1.1$ s and 2.0 s.

4.2 Negative shear mode

Figure 4.5 shows the time evolution of the negative shear mode plasma. The plasma current is ramped up from 1 MA to 3 MA in 1 second as shown in Fig. 4.5(c). Note that a rapid plasma current ramp-up is performed to form ITBs in the simulation. The additional heating of 5 MW is applied simultaneously with the current ramp-up. The bootstrap current increases as the plasma current ramp-up, and peaks at $t \sim 2.4$ s, and then decreases. The bootstrap current drive ratio is $f_{BS} = 0.20$ at the peak of the bootstrap current, and is decreased compared to the high- β_p mode case due to the increased plasma current. As shown in Fig. 4.6(a), the current ramp-up operation causes the OH current to increase in the plasma periphery, and current density profiles have a peak. After $t = 2.0$ s, when the plasma current ramp-up is finished, the peaking of the plasma current density moves toward the plasma center, where the plasma

resistivity is small by the current diffusion. Corresponding to the current profile, a negative shear region is formed (Fig. 4.6(b)). At the position of the peak in the current density profile, the magnetic shear is $s \sim -2.35$ ($\rho = 0.73$) at $t = 2.0$ s and $s \sim -3.26$ ($\rho = 0.61$) at $t = 3.0$ s, which are considerably lower than 0. The safety factor profiles have $q_0 > 1$ at the plasma center, and the q is decreased in the entire radial direction compared to the high- β_p mode case. In addition, compared to the high- β_p mode case, OH heating reaches up to 1.5 MW, especially during the current ramp-up phase, due to the higher OH current.

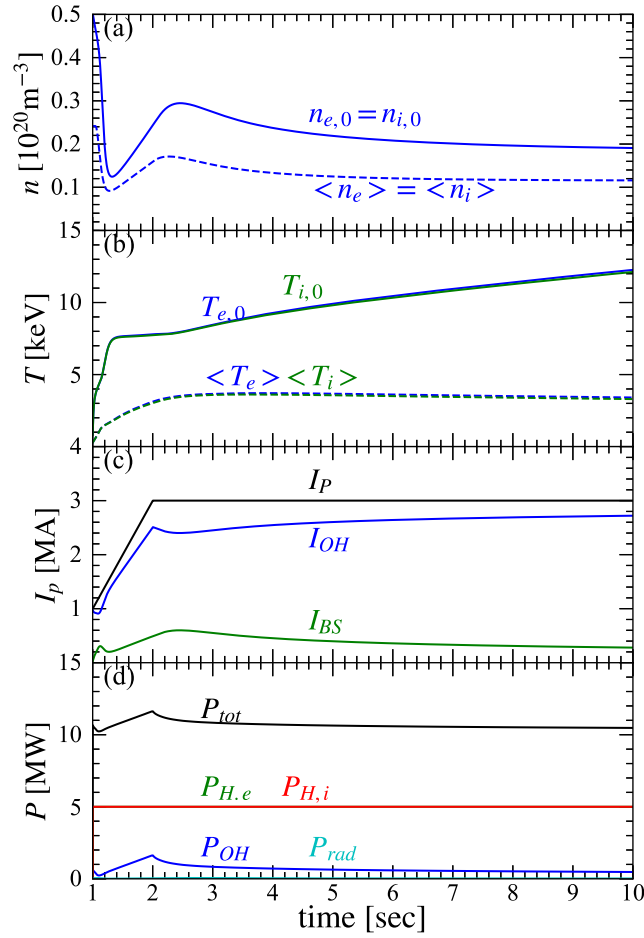


Figure 4.5: Time evolution of (a) volume-averaged and central plasma densities and (b) volume-averaged and central temperatures of electron and ion, (c) plasma current, and (d) heating power.

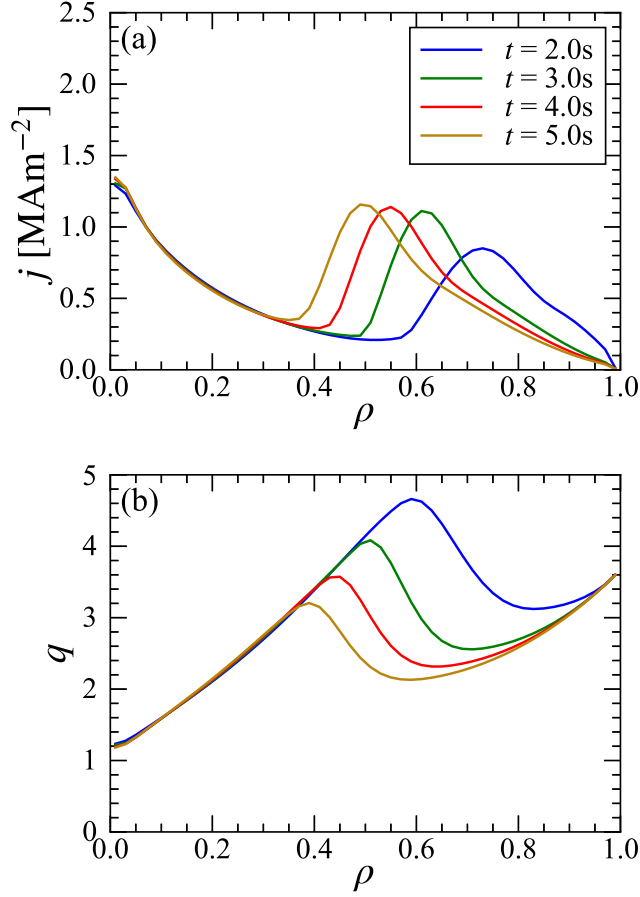


Figure 4.6: Time evolution of radial profiles of (a) current density and (b) safety factor at $t = 1.0\text{s} - 5.0\text{s}$.

Figure 4.7 shows the time evolution of the negative shear mode plasma profiles. The strong gradient regions are formed in the ion temperature and density profiles in the negative shear mode. This is because the transport is improved by a locally reduced diffusion coefficient in the negative shear region. As shown in Fig. 2.12, the CDBM model has a dependence on $F(s, \alpha)\alpha^{1.5}$, and in regions where the magnetic shear is small ($s \lesssim 1.0$), the increase in diffusion coefficient is suppressed even as the normalized plasma pressure gradient increases. Furthermore, since $\chi_{CDBM} \propto \frac{c^2}{\omega_{pe}^2} \frac{v_A}{qR} \propto \frac{1}{n_e^{1.5}} \frac{1}{q}$ in Eq. 2.83, the increase of safety factor and plasma density in the negative magnetic shear region suppress the turbulent diffusion. The formation of the negative magnetic shear is transient, and the ITB moves toward the plasma center according to the current diffusion. The magnetic diffusion equation Eq. 2.10 can be simplified and expressed as

$$\frac{\partial B_p}{\partial t} = \frac{\eta}{\mu_0} \frac{\partial^2 B_p}{\partial \rho^2} \sim -\frac{\eta}{\mu_0 r_s^2} B_p, \quad (4.1)$$

therefore, the magnetic diffusion time (penetration time for current) is given by

$$\tau_\eta = \frac{\mu_0 r_s^2}{\eta}, \quad (4.2)$$

where r_s is the characteristic radial scale length. At $t = 3.0\text{s}$, $\tau_\eta \sim 454\text{s}$ for $r_s = a = 1.2\text{m}$

and the plasma resistivity $\eta \sim 5 \times 10^{-8} \Omega\text{m}$ at $\rho = 0.6$. Since the energy confinement time $\tau_E = 0.31$ s and the particle confinement time $\tau_p = \frac{a^2}{2.4^2 D} \sim \frac{1.2^2}{2.4^2 \times 0.42} = 0.59$ s at $t = 3.0$ s, the time scale of transport is much smaller than the time scale of magnetic diffusion, so the ITB moves according to the negative magnetic shear region.

As shown in Fig. 4.7(b), the ion temperature at the plasma center increases as the ITB moves toward the center of plasma. On the other hand, as shown in Fig. 4.7(a), the plasma density at the center of plasma reaches a maximum at $t = 3.0$ s and decreases subsequently. From Eq. 2.8 and Eq. 2.9, neglecting the convection term, the simplified transport equation is given by

$$\frac{\partial}{\partial t}(n_i V') = \frac{\partial}{\partial \rho} \left(V' \langle |\nabla \rho|^2 \rangle D_i \frac{\partial n_i}{\partial \rho} \right) + S_i V', \quad (4.3)$$

$$\frac{\partial}{\partial t} \left(\frac{3}{2} n_i T_i V'^{\frac{5}{3}} \right) = V'^{\frac{2}{3}} \frac{\partial}{\partial \rho} \left(V' \langle |\nabla \rho|^2 \rangle \chi_i n_i \frac{\partial T_i}{\partial \rho} \right) + P_i V'^{\frac{5}{3}}. \quad (4.4)$$

Considering the steady state for simplicity, the source and diffusion terms are balanced, so the source profiles determine the curvature (second radial derivative) of the density and temperature profiles, and the gradients are expressed as

$$\frac{\partial n_i}{\partial \rho} = - \frac{1}{V' \langle |\nabla \rho|^2 \rangle} \frac{1}{D_i} \int_0^\rho S_i V' d\rho', \quad (4.5)$$

$$\frac{\partial T_i}{\partial \rho} = - \frac{1}{V' \langle |\nabla \rho|^2 \rangle} \frac{1}{n_i \chi_i} \int_0^\rho P_i V' d\rho'. \quad (4.6)$$

Inside the ITB, the turbulent diffusion is dominant, so the plasma density profiles are flattened due to the absence of particle sources. In the ITB region, the turbulent transport is suppressed, so the density and temperature gradients strongly depend on the source profiles. At $t = 3.0$ s, the curvature is $\frac{\partial^2 n_i}{\partial \rho^2} \Big|_{\rho=0.71} = 4.8$ and $\frac{\partial^2 n_i}{\partial \rho^2} \Big|_{\rho=0.53} = -7.2$, and the normalized density gradient at the midpoint of the ITB is $\frac{R}{L_{n_i}} \Big|_{\rho=0.62} = 12.7$. At $t = 5.0$ s, $\frac{\partial^2 n_i}{\partial \rho^2} \Big|_{\rho=0.57} = 2.2$ and $\frac{\partial^2 n_i}{\partial \rho^2} \Big|_{\rho=0.41} = -3.2$, and the normalized density gradient at the ITB midpoint is $\frac{R}{L_{n_i}} \Big|_{\rho=0.49} = 6.8$. On the other hand, for the ion temperature profile with a strong source in the plasma center region, at $t = 3.0$ s, the curvature is $\frac{\partial^2 T_i}{\partial \rho^2} \Big|_{\rho=0.71} = 101.6$ and $\frac{\partial^2 T_i}{\partial \rho^2} \Big|_{\rho=0.53} = -137.2$, and the normalized density gradient at the midpoint of the ITB is $\frac{R}{L_{T_i}} \Big|_{\rho=0.62} = 15.1$. At $t = 5.0$ s, $\frac{\partial^2 T_i}{\partial \rho^2} \Big|_{\rho=0.57} = 111.6$ and $\frac{\partial^2 T_i}{\partial \rho^2} \Big|_{\rho=0.41} = -176.4$, and the normalized density gradient at the midpoint of the ITB is $\frac{R}{L_{n_i}} \Big|_{\rho=0.49} = 13.8$. Thus, from $t = 3.0$ s to 5.0 s, the normalized density gradient decreases by about 46%, while the normalized temperature gradient decreases by only about 9%. When the ITBs are in the small source region, a small curvature and weak gradient are formed. Maintaining strong ITBs also affects the density and temperature at the plasma center, and the density decreases while the ion temperature increases in this simulation. At the end of the discharge, the ITB is no longer formed in the plasma, and the plasma profiles become comparable to an L-mode plasma.

In order to form the ITB constantly, it is necessary to maintain a negative magnetic shear with bootstrap current or external current drive after the plasma current ramp-up phase. In addition, a particle source in the plasma center is required to maintain the ITB in the plasma density profile. The poloidal β value $\beta_p = 0.28$ at $t = 2.4$ s, when the bootstrap current fraction f_{BS} peaks, is decreased compared to the high- β_p mode case. The plasma density and temperature are higher in the negative magnetic shear mode case than in the high- β_p mode case. Therefore, as in the high- β_p mode case, the plasma current is reduced to increase the bootstrap current fraction, which is also necessary in terms of maintaining the ITB.

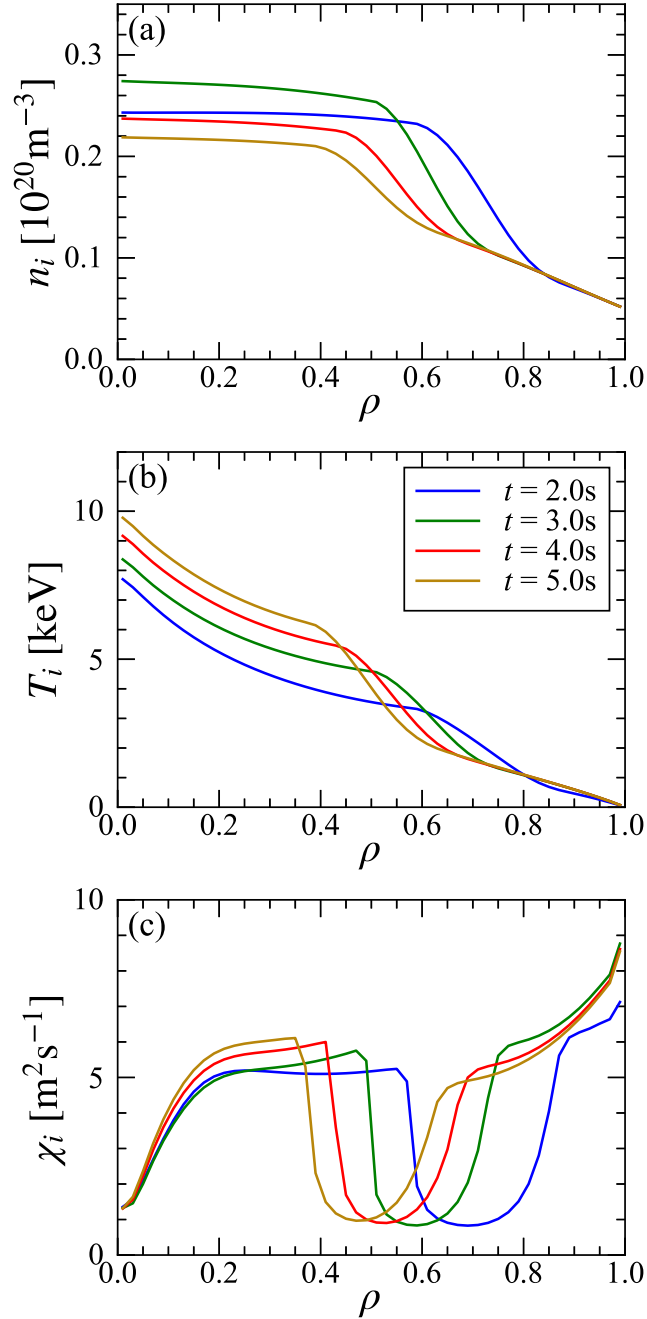


Figure 4.7: Time evolution of radial profiles of (a) plasma density, (b) ion temperature, (c) ion thermal diffusion coefficient at $t = 1.0$ s – 5.0 s.

Figure 4.8 shows the radial profiles of tungsten density and neoclassical convection velocity. The tungsten density profiles are similar to the bulk density profiles. Step-like tungsten profiles are formed in the phases when the strong ITBs are formed. The normalized density gradient is $\left. \frac{R}{L_{nW}} \right|_{\rho=0.62} = 23.2$ at $t = 3.0$ s, forming a stronger gradient than the bulk. The strong gradients accelerate the accumulation of impurities in the core plasma region. In this simulation, the tungsten turbulence model is the same as the bulk one, so the turbulent transport is suppressed in the ITB region. As shown in Fig. 4.8(b), the tungsten density profiles change in accordance with the neoclassical convection velocity profiles. As described above, due to the large bulk density gradient, the tungsten density increases from $t = 2$ s to 3 s. For $t = 4$ s, the bulk density gradient decreases while the temperature gradient is maintained, so the tungsten convection directs outward to reduce the tungsten density.

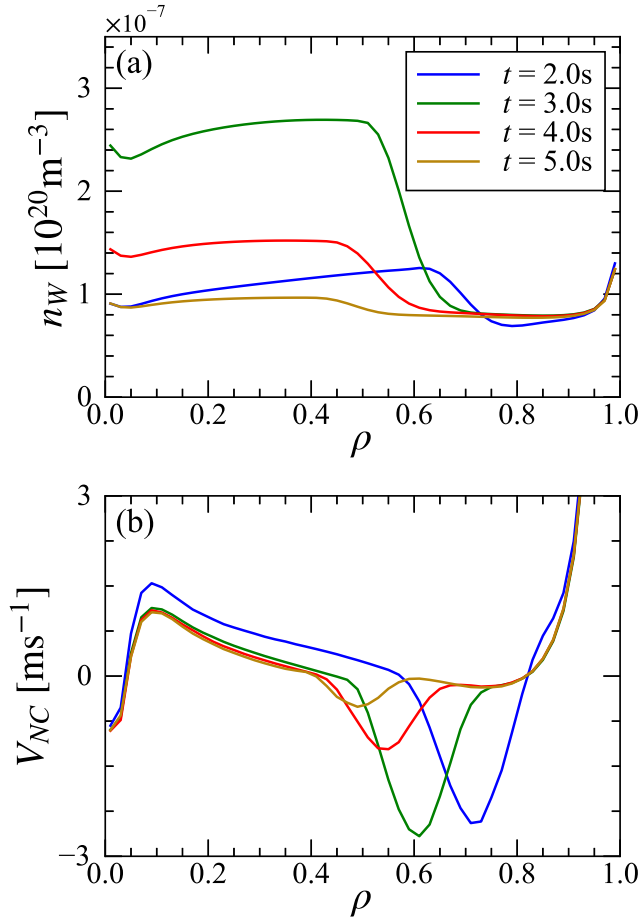


Figure 4.8: Time evolution of radial profiles of (a) W density and (b) W neoclassical convective velocity at $t = 1.0$ s – 5.0 s

4.3 Summary of the chapter

In this chapter, results are shown for high- β_p mode and negative magnetic shear mode plasmas using a plasma transport simulation scheme including impurities. As observed in the experiments, the plasma profiles with the ITBs are obtained. In the high- β_p mode, plasma

heating and bootstrap current play an important role in improved confinement, and ITBs are observed to form in the plasma center region. In the weak magnetic shear region, the increased pressure gradient suppresses turbulent transport and improves the confinement. In the negative shear mode, the plasma heating and the current density peak formed by the plasma current ramp-up play an important role, and the transient ITB formation is confirmed. Negative magnetic shear regions are formed near the peak of the current density profiles, which suppress turbulent transport and improves the confinement. Compared to the negative magnetic shear mode, the high- β_p mode results in a higher bootstrap current fraction and higher poloidal β value. In the present simulation of negative magnetic shear mode, to increase the beta value, decreasing the plasma current is effective as in the high β_p mode case. As the bootstrap current fraction increases with decreasing plasma current, the magnetic shear decreases at the end of the negative magnetic shear mode case, and thus the ITB can be maintained. Furthermore, the strong pressure gradient in the ITB region has a positive feedback effect increasing the bootstrap current. It is known that external current drive using NBI or RF heating is important for steady ITB formation, and results of improved confinement mode using NBI and LH are presented in the next chapter.

The results for tungsten transport in the ITB region differ for the high- β_p and negative magnetic shear modes. In case of the high- β_p mode, the tungsten accumulation does not occur in the plasma center. In case of the negative shear mode, on the other hand, the tungsten density increased in the ITB region, resulting in tungsten accumulation in the center of the plasma. In the improved confinement plasma, the turbulent transport is locally improved, and the dominant mechanism for impurity transport in that region is the non-diffusive terms (neoclassical particle pinch and neoclassical temperature screening), which strongly depend on the background ion profiles. The simulation results show that the balance between the temperature and density gradients of the bulk ions either enhances the accumulation or suppresses the accumulation. An indication of the ITB strength is the gradient length, and it is found that the heating source and the particle source of the background plasma play an important role. In regions where the turbulent transport is suppressed, the sources determine the curvature of the density and temperature profiles, and the sources vary in the radial direction. Therefore, depending on the position of the ITBs, the magnitude of the ion density gradient and the ion temperature gradient varies with radial direction, resulting in a change in the neoclassical transport as well. In the present high- β_p mode calculation, the ITBs are formed in a narrow region of the center of the plasma, and the outward neoclassical convection of the tungsten is driven by the ion temperature gradient due to the strong heating source. The calculation of the negative shear mode show that the tungsten density is significantly increased due to the presence of particle sources when ITBs are formed in the outer plasma region. On the other hand, when the ITB moves to the central region of the plasma, the increase in tungsten density is suppressed by the presence of the heating source. Although the plasma density in the plasma center needs to be increased to achieve high plasma performance, it should be noted that impurity accumulation can be accelerated. In the next chapter, simulations of the improved confinement mode, in which NBI is used to supply the particles to the plasma center, are also shown.

Chapter 5

Impurity transport including toroidal rotation

The purpose of this chapter is to investigate impurity transport behavior in improved confinement tokamak plasmas forming ITBs, and centrifugal effects of impurity due to toroidal rotation are considered. Variation of the neoclassical transport is paid attention, and plays an important role in the plasma central regions. Simulations focused on the W neoclassical transport are performed for a tokamak experimental device, corresponding to the three confinement modes considered in the ITER operation scenario. Simulations are performed in an improved confinement plasma that is more realistic than in the Chapter 4. For the neoclassical transport calculations, the NCLASS code is used, and since original one does not include the effect of impurity poloidal asymmetry, the model is extended to include the toroidal rotation effect. Comparisons will be performed for the cases with the effect of poloidal asymmetry and for the case with different pinch models. Comparisons are performed with tungsten profiles in experiments and transport coefficients obtained by first principles calculations. Then, impurity transport simulations are performed to investigate the effect of the toroidal rotation on impurity distribution formation in improved confinement modes. The variation of the neoclassical transport and the tungsten distribution due to the toroidal rotation effect of tungsten are shown. Magnitudes of impurity accumulation are evaluated for the temperature screening effect in the improvement transport region.

5.1 Simulation of improved confinement plasma

In this section, the results of transport simulations in the absence of impurities are presented using various JET plasma parameters to clarify the effects of improved confinement. Three cases are simulated: H-mode, hybrid scenario and advanced scenario plasmas.

For the H-mode case, the plasma parameters of $I_p = 2.75$ MA and $B_t = 2.5$ T are used Ref. 120. Only NBI is applied as external heating and fueling with $P_{NB} = 17.5$ MW at the radial position $\rho_0^{NB} = 0.0$ with the width $\rho_w^{NB} = 0.6$. The toroidal rotation velocity and pedestal profiles are given by the experimental profiles shown in Ref. 120 (JET#83351 at $t = 12.37$ s). The pedestal region is set for $0.9 \leq \rho \leq 1.0$, and the radial profile is set with Eq. 2.96 for $0.8 \leq \rho < 1.0$. The plasma density and temperature profiles are calculated by transport analysis for $\rho < 0.8$ without the pedestal region.

In the hybrid scenario case, the plasma parameters of $I_p = 1.7$ MA and $B_t = 2.0$ T are used Ref. 166. The profile of the NBI heating power is set to be $P_{NB} = 16$ MW with $\rho_0^{NB} = 0.0$ and

$\rho_w^{NB} = 0.6$, and the toroidal rotation velocity and the pedestal profiles are used in Ref. 166. In the advanced scenario case, the same plasma parameters as in the hybrid scenario cases are used for the NBI heating power, toroidal rotation velocity, and pedestal profiles, but a different value is used for the radial width $\rho_w^{NB} = 0.8$. For a current profile control, the additional LH heating with $P_{LH} = 2$ MW is applied with $\rho_0^{LH} = 0.4$ and $\rho_w^{LH} = 0.1$, and parallel refractive index $N_{\parallel} = 2.1$ [167].

Radial profiles of main plasma in steady state for three cases are obtained. Figures 5.1 and 5.2 show the radial profiles of plasma quantities. As shown in Fig. 5.1, the safety factors of the hybrid scenario and advanced scenario plasmas have higher values in the entire radial region, compared to the H-mode plasma. At the plasma edge, $q \sim 3$ for the H-mode plasma, while $q > 5$ for the hybrid scenario and advanced scenario plasma. In the H-mode plasma, there is a region with $q = 1$, possibly causing sawtooth crashes, which are not considered in this simulation.

In the advanced scenario plasma, the safety factor profile is flattened depending on the position of the LH heating, and negative magnetic shear regions are also formed. The magnetic shear is $s = 0.54$ (H-mode), $s = 0.38$ (Hybrid scenario), and $s = -0.50$ (Advanced scenario) at $\rho = 0.3$. The bootstrap current fraction is $f_{BS} = 0.18$ (H-mode), $f_{BS} = 0.21$ (Hybrid scenario), $f_{BS} = 0.26$ (Advanced scenario). In the advanced scenario plasma, the fraction of ohmic current is less than 10 % and the non-inductive current components are dominant. In the H-mode case, the plasma current density profile is monotonic, and is peaked at the magnetic axis where the plasma resistivity is small. In the hybrid scenario and advanced scenario cases, on the other hand, the plasma current is smaller than in the H-mode case, and the bootstrap current fraction and externally-driven current can be relatively large, resulting in a flattened or hollow current density profile. This gives a flattened or reversed q profile in the plasma confinement region, reducing the turbulent transport level to improve confinement.

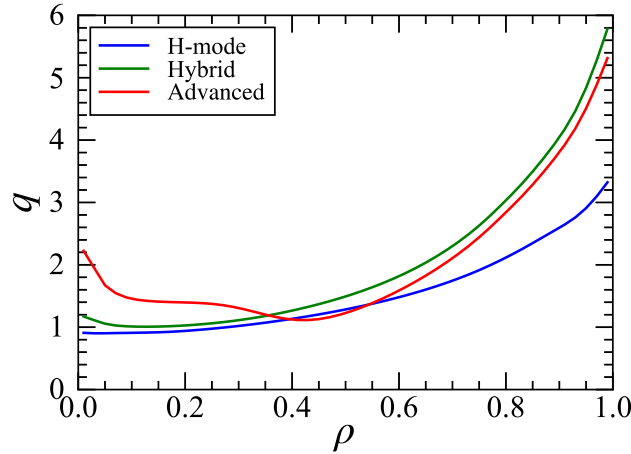


Figure 5.1: Comparison of the radial profiles of safety factor q at the steady states of the H-mode, hybrid scenario and advanced scenario cases.

In the H-mode case, the plasma density is higher than in the other conditions, and the profile is flattened for $\rho < 0.8$. The density gradient is relatively large in the plasma center

due to the lower transport level and the stronger NBI particle source. The Mach number is $M_i = 0.26$ at the plasma center, and a flat radial profile is given for $\rho < 0.6$. The H -factor and normalized- β value are $H_{98y2} = 0.87$ and $\beta_N = 1.35$, respectively. In the hybrid scenario case, the plasma density is lower than in the H-mode case, but the plasma rotation is larger for $0.0 \leq \rho \leq 0.6$, with $M_i = 0.36$ at the plasma center and a maximum value of $M_i = 0.45$ at $\rho = 0.35$. In the weak shear region inside $\rho \sim 0.4$, turbulent transport is suppressed and the plasma gradient becomes stronger. Although a ITB is formed at $\rho \sim 0.4$ in Ref. 166, the $F_{E \times B}$ in Eq. 2.86 to be less than a few percent of the CDBM thermal diffusion coefficient, so that the effect of transport suppression due to toroidal rotation is small. A model for $F_{E \times B}$ that can be adjusted based on experiments is proposed in Ref. 168. The H -factor and normalized- β value are $H_{98y2} = 1.11$ and $\beta_N = 1.97$, respectively.

In the advanced scenario case, the plasma gradient increases steeply for $\rho < 0.4$ with the LH heating. The ITB formation results in larger temperature at the plasma center than that in the hybrid scenario plasma. This is due to the suppression of turbulent diffusion by the relaxation of magnetic shear s resulting from the LH current drive and by the increase of pressure gradient α resulting from the heating. The width of this increased region in the plasma gradient is approximately same with the width of the LH heating. The same toroidal rotation velocity is set, but the ion temperature is higher than in the hybrid case due to the formation of an ITB, giving a smaller Mach number $M_i = 0.30$ at the plasma center. The H -factor and normalized- β value are $H_{98y2} = 1.26$ and $\beta_N = 2.34$, respectively.

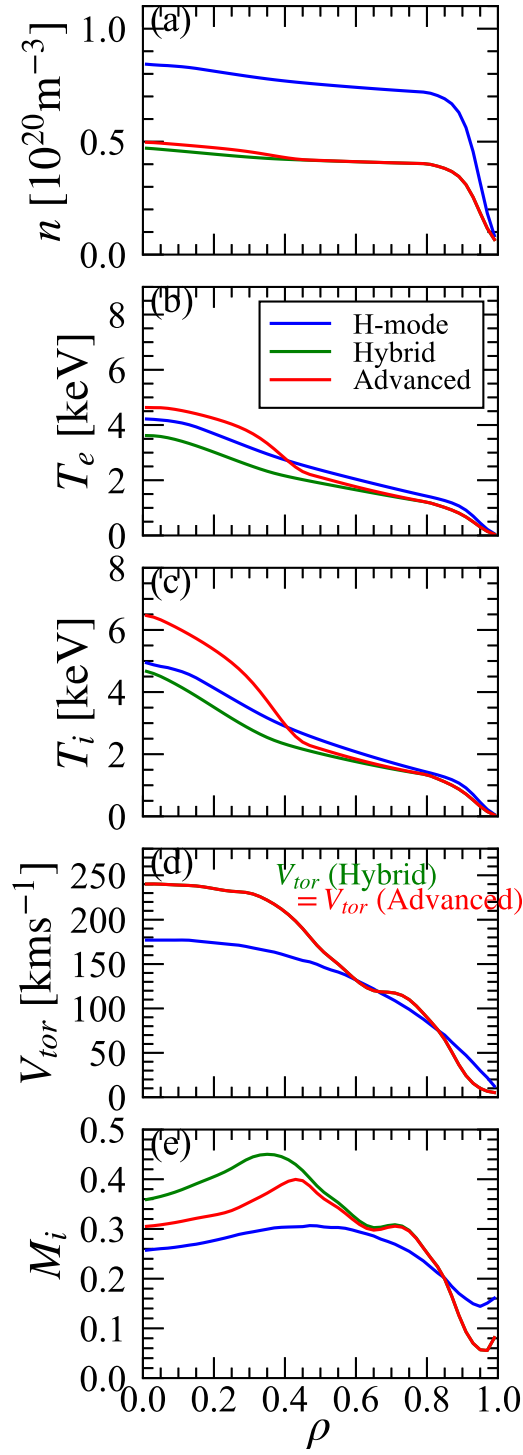


Figure 5.2: Comparison of the radial profiles of (a) plasma density, (b) electron temperature, (c) ion temperature, (d) toroidal rotation velocity, (e) Mach number of the main ion at the steady states of the H-mode, hybrid scenario and advanced scenario cases. The toroidal velocity profile in the advanced scenario case is the same as in the hybrid scenario case.

5.2 Simulation of impurity transport

5.2.1 Validation of impurity transport model

The impurity transport model described in Sec. 2.4.4 is introduced, and for the first step, the impurity transport simulation model is validated for the H-mode case in Ref. 120. For that purpose, calculations with fixed background plasma profiles are carried out. As shown in Fig. 5.3, the plasma density, temperature and toroidal rotation velocity profiles are used for two time slices at $t = 9.32$ s, when the W is not peaking in the center of the plasma (case A), and at $t = 12.37$ s, when the W is peaking in the plasma center region (case B). Here, it is assumed that the electron and ion density are same, and the electron and ion temperature are same. In case A, the plasma density is hollow at the plasma center, and has positive gradient. In case B, the plasma density is peaked at the plasma center, and has negative gradient. The toroidal rotation velocity profiles are flat at the plasma center with the Mach numbers of the main ion at the plasma center $M_i = 0.19$ in case A and $M_i = 0.30$ in case B. The W transport analysis using NEO and GKW shows that the neoclassical transport is dominant at the plasma center, while the turbulent contribution is also significant at the plasma periphery [120]. The turbulent diffusion component is nearly equal or larger than the neoclassical component for $\rho > 0.45$, and the turbulent convection component is larger and directing inward for $\rho > 0.8$. The turbulent diffusion coefficient of W is estimated from $D_{W,GKW}/\chi_{i,NEO}$ in Ref. 120 to be set as in Fig. 5.4 by using the following turbulent diffusion coefficient, a combination of constant and parabolic ones, is used

$$D_W^{turb} = \begin{cases} C_1^{W,turb} & (0.0 \leq \rho \leq \rho_1^{W,turb}) \\ \frac{C_2^{W,turb} - C_1^{W,turb}}{\rho_2^{W,turb^2} - \rho_1^{W,turb^2}} \rho^2 + \frac{C_1^{W,turb} \rho_2^{W,turb^2} - C_2^{W,turb} \rho_1^{W,turb^2}}{\rho_2^{W,turb^2} - \rho_1^{W,turb^2}} & (\rho_1^{W,turb} \leq \rho \leq \rho_2^{W,turb}) \\ C_2^{W,turb} & (\rho_2^{W,turb} < \rho \leq 1.0) \end{cases} \quad (5.1)$$

and the turbulent convection component of impurity is ignored. The radial profiles of the neoclassical thermal diffusion coefficient of the main ion obtained by the NCLASS code are used for the estimation assuming $\chi_{i,NEO} = \chi_{i,NCLASS}$. The density of W at the boundary is set as $n_{W,edge} = 2.5 \times 10^{14} \text{ m}^{-3}$ in case A and $n_{W,edge} = 1.0 \times 10^{14} \text{ m}^{-3}$ in case B.

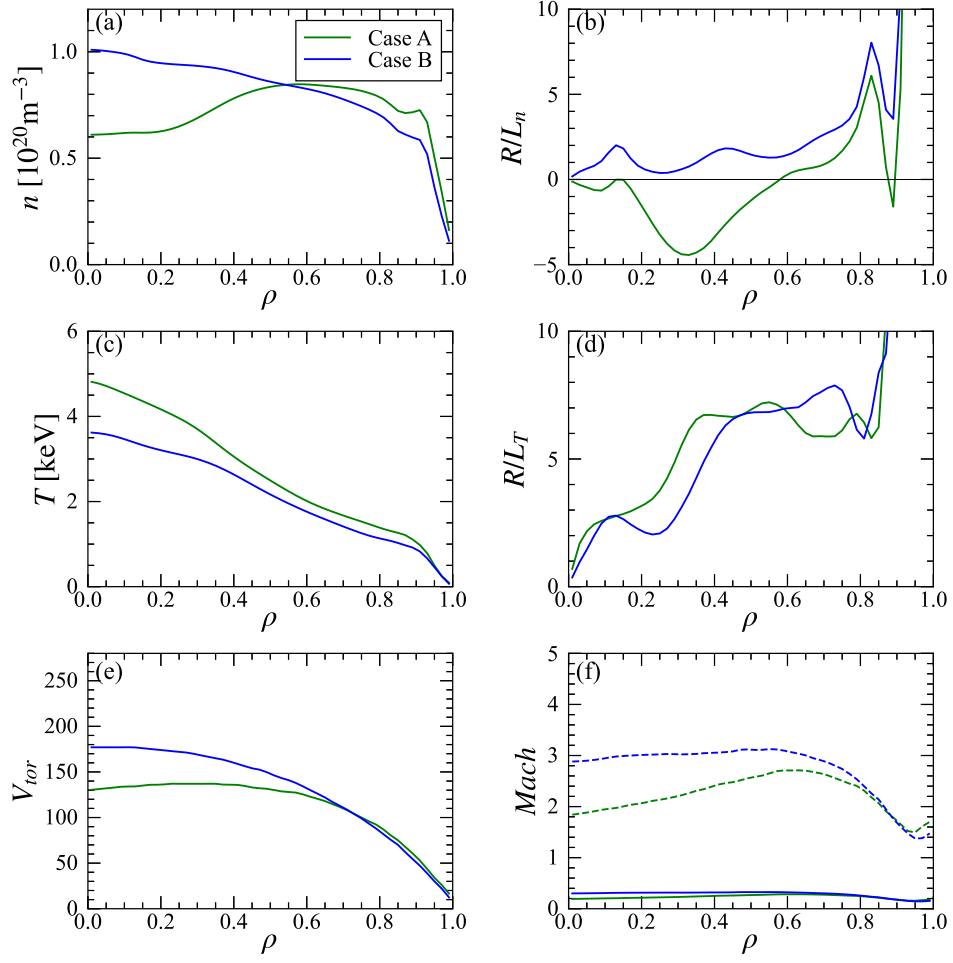


Figure 5.3: Referenced radial profiles of (a) plasma density, (b) normalized plasma density gradient, (c) ion temperature, (d) normalized ion temperature gradient, (e) toroidal velocity of the main ion, (f) Mach number of the main ion (solid line) and the W (dashed line).

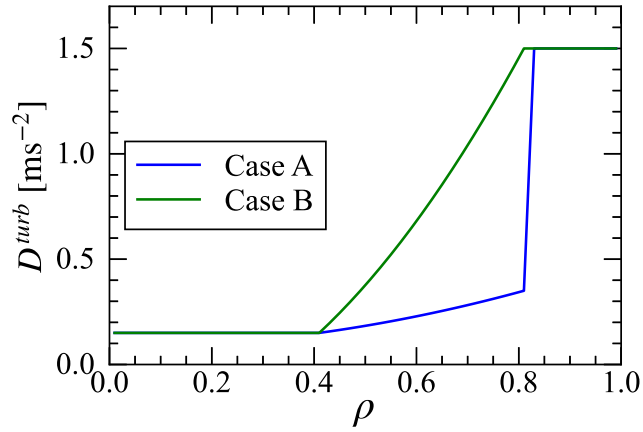


Figure 5.4: The radial W transport coefficients of turbulent diffusion in case A and case B.

5.2.1.1 Neoclassical transport without poloidal asymmetry

First, the transport coefficients and density profiles of W are evaluated using the conventional NCLASS code, which is not included the effect of poloidal asymmetry. The diffusion coefficient profiles are shown in Figs. 5.5(a) and 5.5(b). The neoclassical diffusion coefficients D^{NC} are about 1/10 of the turbulent diffusion coefficient, consistent with the results shown in Ref. 124, which compares the diffusion coefficients of GKW and NEO without toroidal rotation effect. As shown in Fig. 5.6, W does not accumulate in both case A and case B in the plasma center region. The neoclassical convection velocity is outward over a broad region in case A due to the negative ion density gradient and the positive ion temperature gradient. In case B, the neoclassical convection velocity is outward over a broad region, except for the plasma center region, where the temperature screening effect of the positive ion temperature gradient is increased. The neoclassical convection velocities shown in Fig. 5.5 also include the Ware pinch, which is driven inward for the entire radial direction. In Case A, even the maximum value of $\rho = 0.3$ is about 3% of the V_{NC} . The W density profiles obtained here significantly differ from the experimental and first-principles calculations in Ref. 124. Especially in case A shown in Fig. 5.6(b), only the sign of the W normalized density gradient for $0.00 < \rho < 0.65$ is consistent, but the others are completely different. In case B, only the sign of the normalized gradient matches for $0.00 < \rho < 0.15$, but the others are completely different. In both cases, the value of the normalized gradient is significantly underestimated.

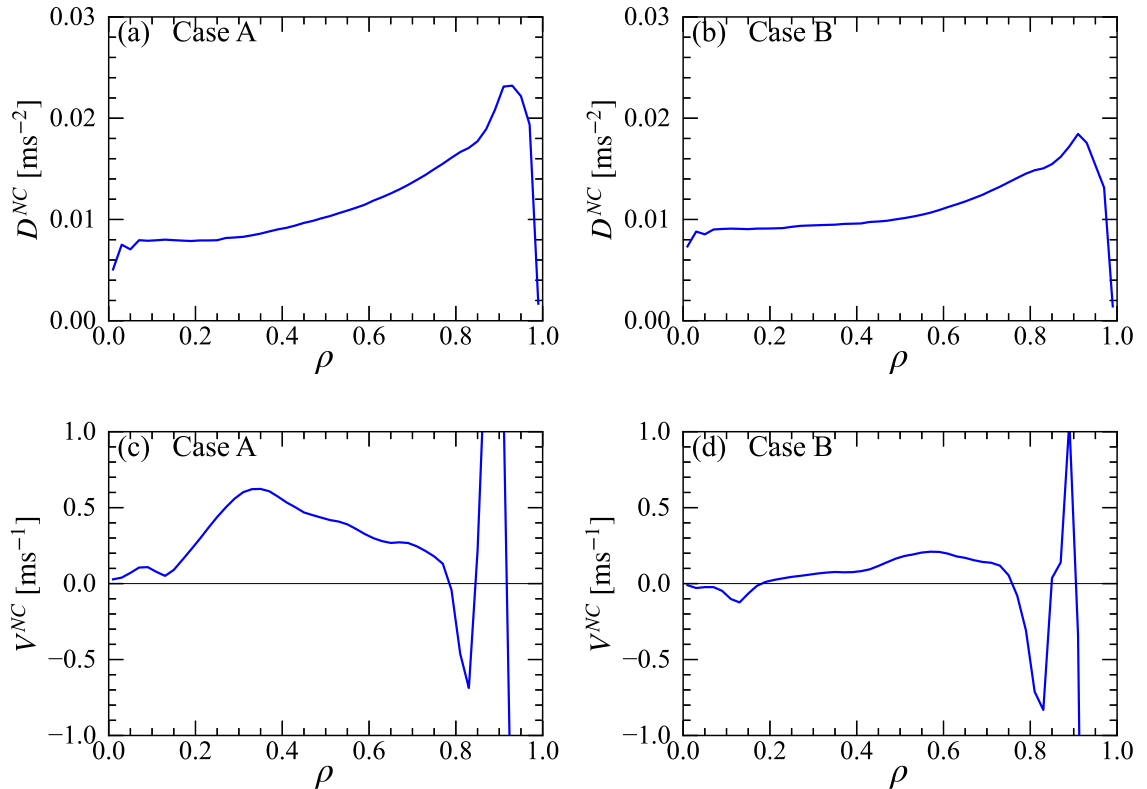


Figure 5.5: The radial W transport coefficients of neoclassical diffusion in (a) case A and (b) case B, and neoclassical convective velocity in (c) case A and (d) case B.

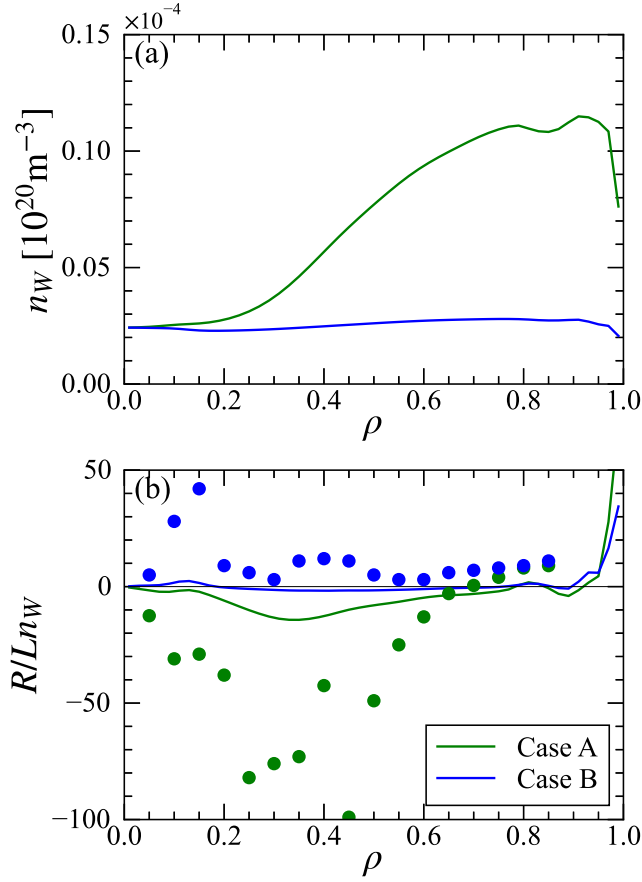


Figure 5.6: Radial profiles of (a) W density and (b) normalized W density gradient in case A and case B. The dot plots in (b) represent the LFS value obtained from the NEO+GKW analysis in Ref. 120.

5.2.1.2 Neoclassical transport with PHZ/ E_r pinch model

Here, the transport coefficients and density profiles of W are evaluated using the transport coefficients obtained by the NCLASS code and the PHZ/ E_r pinch model. As in the previous section, the effects of poloidal asymmetry are not included in the NCLASS code. The radial profiles of diffusion coefficient are shown in Figs. 5.7(a) and 5.8(b). The neoclassical diffusion coefficients D^{NC} are almost the same values as in Fig. 5.5 for both case A and case B for $0.2 < \rho < 1.0$, about 1/10 of the turbulent diffusion coefficients. For $0.0 < \rho < 0.2$, however, the D_{NC} increases sharply. This is because that W strongly accumulates in both case A and B at the central region, as shown in Fig. 5.8. At the plasma center, $\alpha_{trace} = \frac{n_Z Z^2}{n_i} \sim 0.37$ in case A and $\alpha_{trace} \sim 1.39$ in case B, thus the impurity trace limit condition ($\alpha_{trace} > 1$) is violated and the W particle cannot be treated as a test-particle (passive species). As a result, the diffusion coefficients and convection velocities evaluated by the NCLASS code are significantly increased compared to Fig. 5.5, due to the increased collisions of W. The formation of W peaking at the plasma center region is due to the inward flow driven by the E_r pinch, which is driven independently of the convection of the NCLASS, as shown in Figs. 5.7(c) and 5.8(d). Note that the magnitude of the PHZ pinch is small in the entire radial

direction. The stronger toroidal rotation in the present calculations drives the increased radial electric field E_r , which becomes $E_r > 50 \text{ kVm}^{-1}$ at $\rho \sim 0.3$, where the value of the E_r pinch is significant. In a broad radial region, $\alpha_{E_r} > 1$, and the model is no longer valid. Therefore, as shown in Fig. 5.8(b), the E_r pinch reproduces the W central peaking compared as the first-principles calculation in Ref. 124, but it does not quantitatively agree. Furthermore, in case A, in which the W central accumulation is not observed in the experimental results, the inward pinch is driven to form the peak in the W profile.

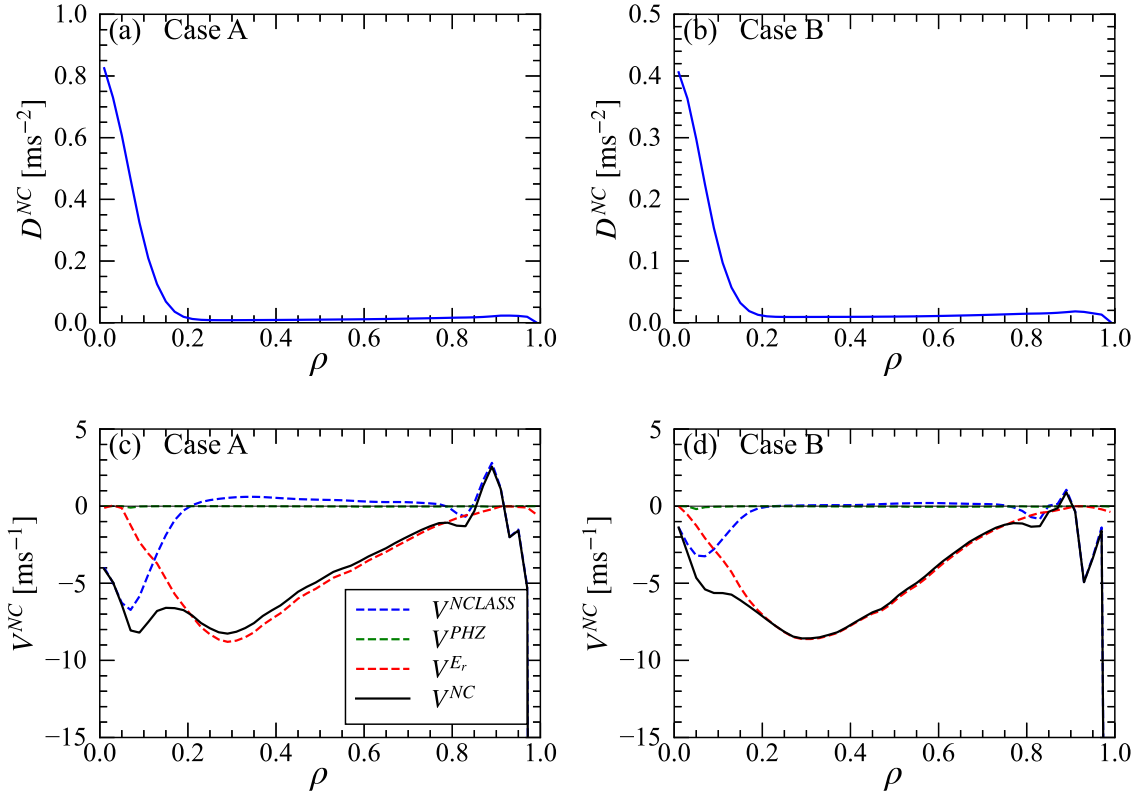


Figure 5.7: The radial W transport coefficients of neoclassical diffusion in (a) case A and (b) case B, and neoclassical convective velocity in (c) case A and (d) case B.

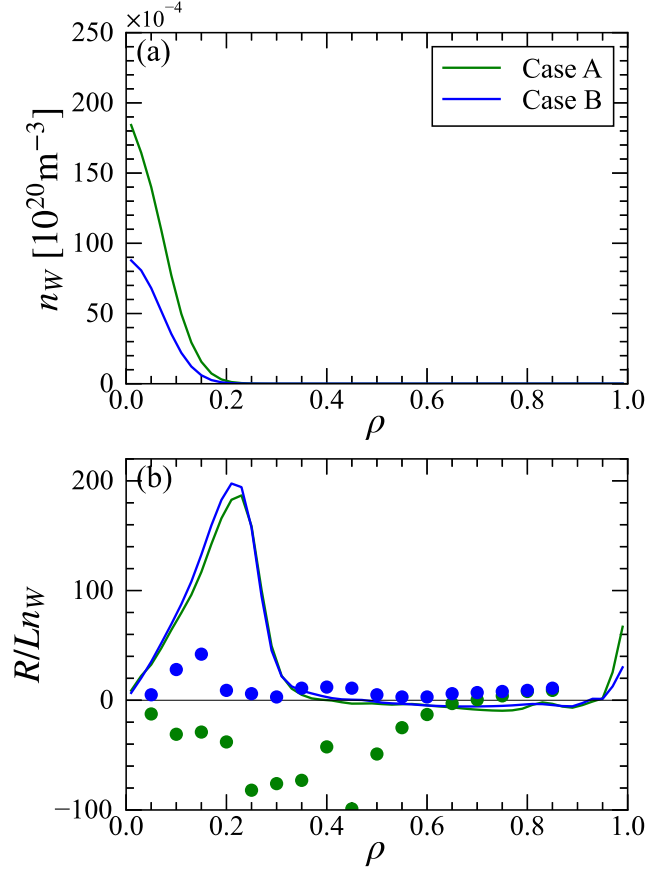


Figure 5.8: Radial profiles of (a) W density and (b) normalized W density gradient in case A and case B. The dot plots in (b) represent the LFS value obtained from the NEO+GKW analysis in Ref. 120.

5.2.1.3 Neoclassical transport with poloidal asymmetry

The diffusion coefficients profiles are shown in Figs. 5.9(a) and 5.9(b). The neoclassical diffusion coefficients D^{NC} is larger for $\rho < 0.6$ in case B because the higher density of main ion and W results in higher collisionality. In particular, the increase of D^{NC} is significant for $\rho < 0.15$, where the W density profile is strongly peaking. As shown in Fig. 5.10, W does not accumulate in case A at the central region, whereas a strong accumulation occurs in case B. In case B, $n_W Z$ at the plasma center is about 1.5 % of the electron density, so the quasi-neutrality condition is not significantly affected in fixed background plasma density. The densities at the low field side (LFS) in Figs. 5.11(c) and 5.11(d) are larger than the magnetic flux surface-averaged (FSA) W density, indicating the W density distribution shifts towards the LFS side due to centrifugal effects, as in Figs. 5.11(a) and 5.11(b). Figure 5.12 shows that the normalized density gradient profiles of W are strongly depending on the magnitude and direction of the neoclassical convection velocities shown in Figs. 5.9(c) and 5.9(d). In both A and B cases, the neoclassical convective components $V_{T,asym}^{NCLASS}$ driven by the main ion temperature gradient are positive in all the radial region. Furthermore, in case A, $V_{N,asym}^{NCLASS}$ is positive for $\rho < 0.6$ due to the positive bulk density gradient, resulting in the absence of W density peaking at the plasma center. On the other hand, $V_{T,asym}^{NCLASS}$ remains positive for $\rho > 0.6$, whereas $V_{N,asym}^{NCLASS}$ becomes negative, and the difference of these terms gives

negative V_{asym}^{NCLASS} for $\rho > 0.7$. As a result, inward convection is driven only in the outer region of the plasma, where a relatively small peak of the W density is formed. In case B, $V_{N,asym}^{NCLASS}$ is negative in all the radial region, and V_{asym}^{NCLASS} is negative. In the outer region of the plasma, since the Mach numbers have similar value in both cases, the geometrical coefficients P_A are similar, but at the plasma center, P_A is increased to be 40 in case B, while 10 in case A. Therefore, in case B, a sharp peak of the W density is formed at the plasma center. As shown in Figs. 5.11(a) and 5.11(b), the W density peak is shifted to the LFS. The poloidal asymmetry appears clearer in case A, when the central density is lower. Note that in case B, the structure is also asymmetric. From Fig. 12 of Ref. 120, the reconstructed experimental values from soft X-ray in case A are $n_{W,LFS} \sim 5.6 - 9.2 \times 10^{15} \text{ m}^{-3}$ at $\rho \sim 0.7$ and $n_{W,LFS} \sim 0.5 - 1.4 \times 10^{15} \text{ m}^{-3}$ at $\rho \sim 0.5$, and are in good agreement with the result in Figs. 5.11(c) and 5.11(d). Similarly, the experimental values in case B are $n_{W,LFS} \sim 0.3 - 1.1 \times 10^{16} \text{ m}^{-3}$ at $\rho \sim 0.7$ and $n_{W,LFS} \sim 0.6 - 1.5 \times 10^{16} \text{ m}^{-3}$ at $\rho \sim 0.15$, and are also in good agreement. In the present simulations, inward turbulent convection, which have large values for $0.8 < \rho < 1.0$, is ignored, whereas the results are valid with the experiment at two points in $\rho < 0.8$. As shown in Fig. 5.12, the normalized W density gradients are similar to the results obtained with NEO and GKW in Ref. 120; the difference around $\rho = 0.3$ in case A is due to the larger turbulent diffusion coefficient evaluated by GKW. In following calculations, impurity transport calculations are performed with the NCLASS code, which includes poloidal asymmetry.

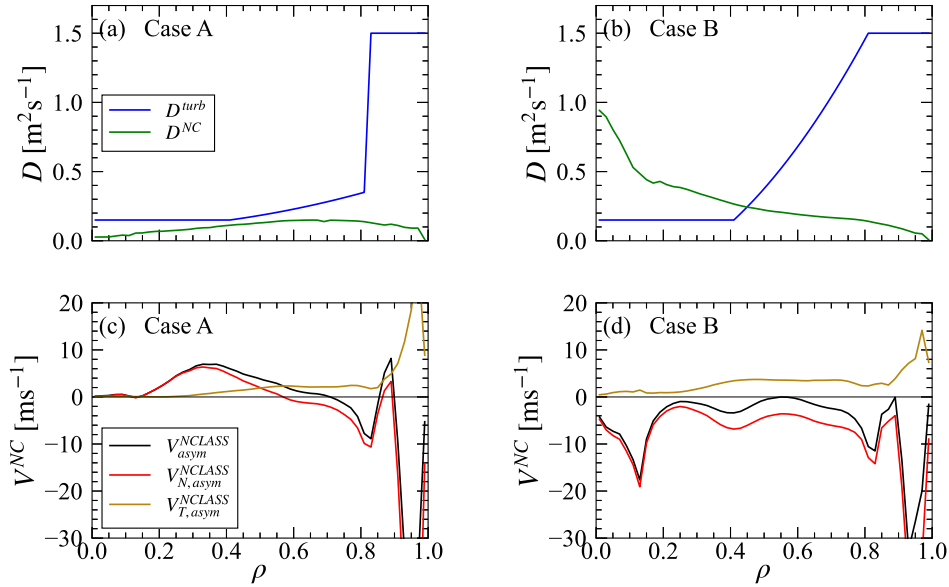


Figure 5.9: The radial W transport coefficients of neoclassical diffusion in (a) case A and (b) case B, and neoclassical convective velocity in (c) case A and (d) case B.

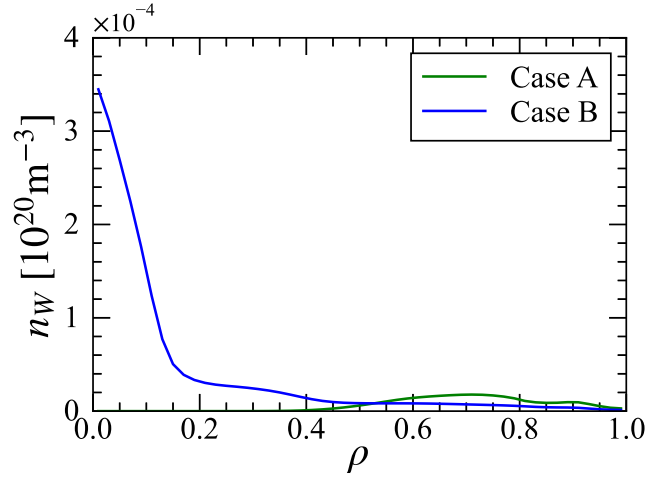


Figure 5.10: Radial profiles of W density of the FSA value.

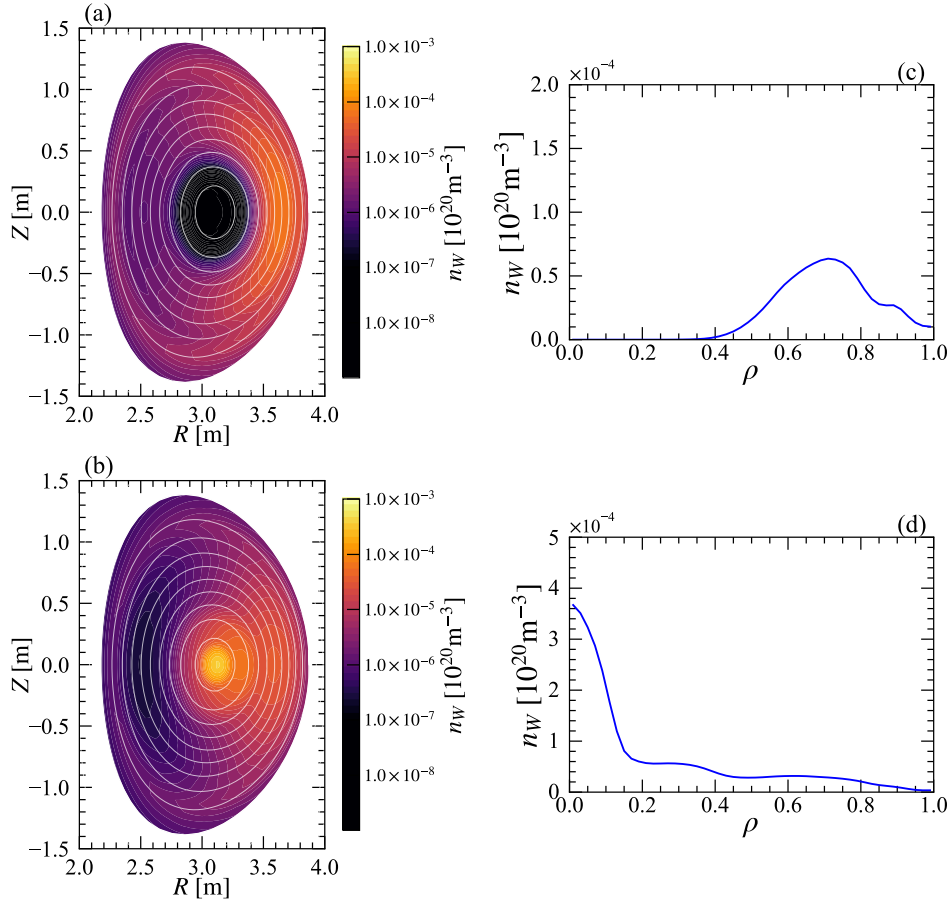


Figure 5.11: Calculated 2D plots of W density on the poloidal cross-section in (a) case A and (b) case B, and radial profiles of W density of LFS value on the midplane in (c) case A and (d) case B. The white lines in (a) and (b) indicate the equilibrium magnetic field configuration. The absolute values of the W density in (a) and (b) are significantly different, and the poloidal asymmetry is clearly visible in (a) since the central peaking is not occurring.

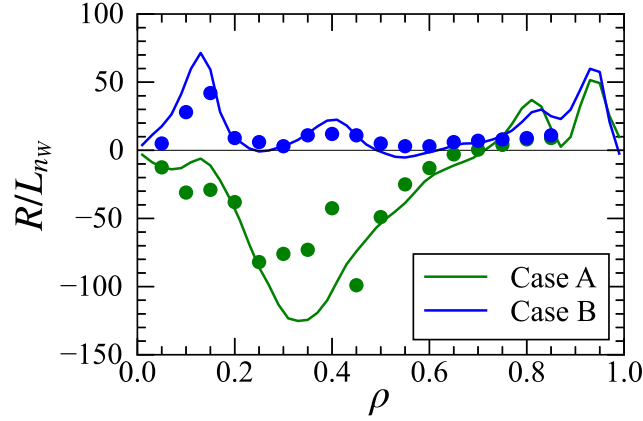


Figure 5.12: Radial profiles of normalized W density gradient of LFS value on the midplane. The dot plots represent the LFS value obtained from the NEO+GKW analysis in Ref. 120.

5.2.2 Tungsten transport in improved confinement plasma

Next, the impurity transport calculation is coupled with that of main plasma species. The cases with improve confinements in Sec. 5.1 are analyzed to include the effects of impurity. In the H-mode case, the same W turbulent diffusion coefficient as in case B in Sec. 5.1 are used. In the hybrid scenario case, the turbulent diffusion coefficient is estimated from Ref. 55 to have $\rho_1^{W,turb} = 0.4$, $\rho_2^{W,turb} = 0.8$, $C_1^{W,turb} = 0.15$, $C_2^{W,turb} = 1.5$, and $D^{W,turb} = 1.9$ at outer edge of the parabola. In the advanced scenario plasma, $\rho_1^{W,turb} = 0.6$ is set from the previous conditions. The W density at the boundary $n_{W,edge} = 1.0 \times 10^{14} \text{ m}^{-3}$, which is same with the case in Sec. 5.2.1, is used for all the cases. The toroidal rotation velocity profiles are shown in Fig. 5.2(d), which gives central Mach numbers $M_Z \sim 2.5$ (H-mode), 3.4 (Hybrid), and 2.9 (Advanced).

In the H-mode case, peaked W density is obtained as show in Fig. 5.13. The normalized W density gradient in the center region of the H-mode case is similar to $R/L_{nW} \sim 50$ comparable to the case with the fixed background plasma profile, which reproduces the result of the central accumulation. With the inclusion of W, the temperature gradient of the plasma is significantly reduced at $\rho \sim 0.1$, due to stronger radiation losses with the impurity accumulation. Compared to Fig. 5.2(b), the electron temperature at $\rho = 0.1$ is reduced by about 4 % and the normalized electron temperature gradient is reduced by about 40 % due to the W accumulation. In this state, the dominant convection component is inward one driven by the bulk density gradient, as shown in Fig. 5.14(b). Thus, in this case of the W accumulation is large enough to cause a strong change in the electron temperature profile, and interactions occur between the plasma and the impurity as follows: (1) increase radiation losses, (2) reduce plasma temperature gradient and increase plasma density gradient, (3) increase neoclassical inward flow, and (4) further W accumulation. A loop from (1) to (4) is followed until a steady state, resulting in enhanced and sustained W density peaking. Furthermore, the temperature screening effect is reduced by the term proportional to P_B in Eq. 2.67 in the presence of a strong asymmetry in the W density.

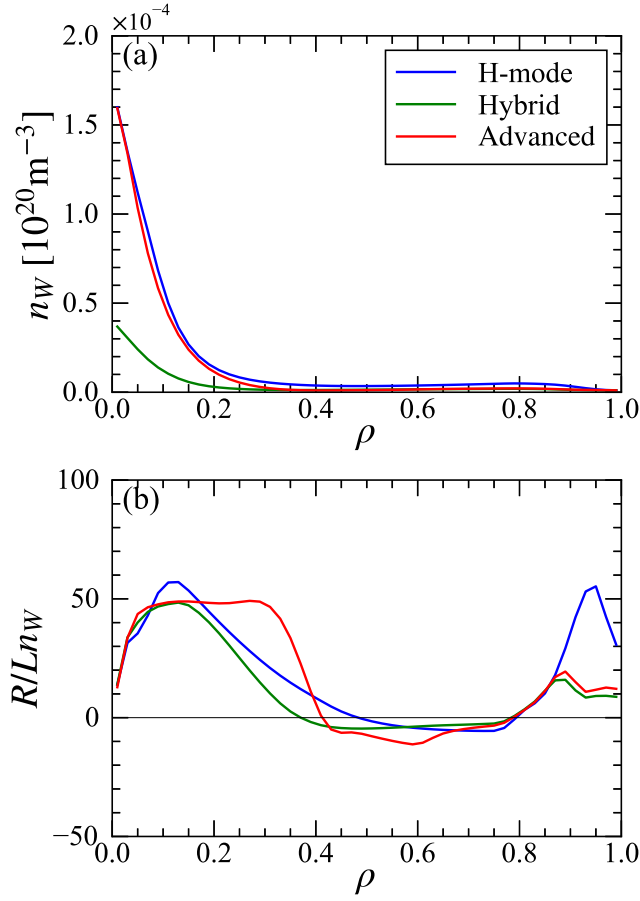


Figure 5.13: Comparison of the radial profiles of (a) W density and (b) normalized W density gradient in case of the H-mode, hybrid scenario and advanced scenario.

In the hybrid scenario case, the W density also has a peak at the center of plasma. For $\rho < 0.4$, as shown in Fig. 5.15, the value of both $V_{T,asym}^{NCLASS}$ and $V_{N,asym}^{NCLASS}$ increase toward the center of plasma due to an increase in the main ion temperature gradient by the reduced turbulent diffusion and an increase in the main ion density gradient by the NBI particle source. The convective component driven by the main ion density gradient is stronger, resulting in an inward particle pinch for $\rho < 0.35$. The temperature screening effect is reduced in the center of the plasma, where the W density peaking is large, as in the H-mode case. Compared to the H-mode case, the W density is lower, whereas the W normalised density gradient is of similar value in the plasma center region. This difference is due to the strong inward neoclassical particle pinch driven in the H-mode case in the plasma periphery $0.85 < \rho < 1.00$, which causes the inflow into the plasma confinement region with a significant amount of tungsten.

In the advanced scenario case, as shown in Fig. 5.16, the particle pinch driven by the main ion density gradient is stronger due to the effects of reduced turbulent diffusion and the NBI particle source, as in the hybrid scenario case. Note that the neoclassical convection components are larger than in the hybrid scenario case, even though the Mach number is smaller, because the plasma transport is improved, and the plasma gradient is stronger than in the hybrid scenario case. The normalized W density gradient changes sharply at $\rho \sim 0.4$, where the ITB is formed, and the W density peaks becomes wider in the radial direction than the hybrid scenario case. This result indicates the acceleration of W accumulation in the ITB region, as confirmed by previous experiments [160–165].

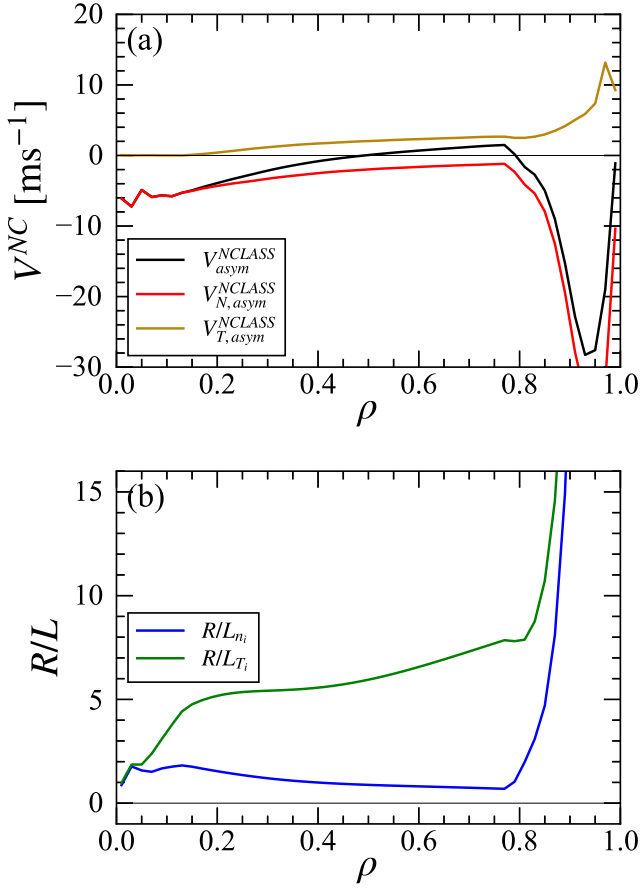


Figure 5.14: Radial profiles of (a) W neoclassical convective velocities and (b) normalized main ion density and temperature gradients of the H-mode case.

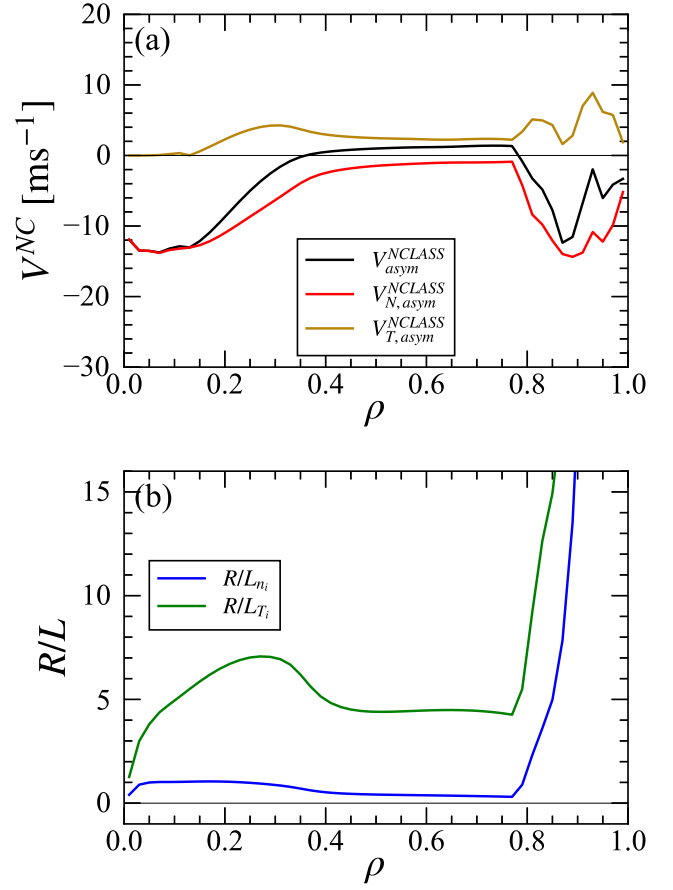


Figure 5.15: Radial profiles of (a) W neoclassical convective velocities and (b) normalized main ion density and temperature gradients of the hybrid scenario case.

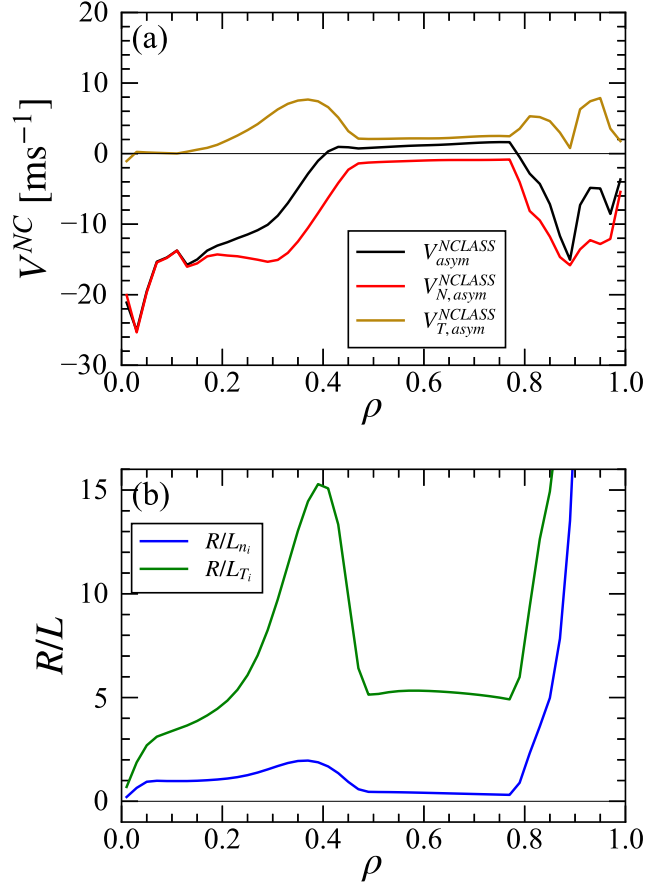


Figure 5.16: Radial profiles of (a) W neoclassical convective velocities and (b) normalized main ion density and temperature gradients of the advanced scenario case.

5.3 Parameter dependence of tungsten transport

In the previous section, it is shown that impurity accumulation occurs for improved confinement discharges. However, since impurity accumulation is not always observed in the experiments, we will investigate the conditions that lead to impurity accumulation from the dependence on the parameters. The first is the effect of rotation, which has a significant impact on impurity transport, and the second is the position of the ITB related to the performance of the transport barrier. In appendix B, the accumulation mechanisms, including toroidal rotation effects of low-Z and medium-Z impurity particle species in the improved confinement region, are investigated.

5.3.1 Toroidal rotation dependence

The dependence on Mach number of the main ion is investigated in the advanced scenario case to clarify the effect of toroidal rotation on the impurity profile formation. To avoid violation of the impurity trace limit condition $\alpha_{trace} = \frac{n_Z Z^2}{n_i} \ll 1$ and breakdown of calculation due to excessive radiation losses, the boundary condition of the impurity density is set to be smaller value $n_{W,edge} = 1.0 \times 10^{13} \text{ m}^{-3}$. Mach number M_{i0} at the plasma center

is used here as a parameter, and the central toroidal velocity is set by $V_{tor,i}^{\rho=0} = M_{i,0} \sqrt{\frac{2T_i}{m_i}}$. The profile of the toroidal velocity is given by using the function $V_{tor,i} = V_{tor,i}^{\rho=0} (1 - \rho^3)^2$ to be similar to that of the hybrid case in Fig. 5.2(d).

In the small M_{i0} case as $M_{i0} = 0.1$, the W density profile flattens due to the turbulent diffusion without the enhancement by the toroidal rotation. The case with $M_{i0} = 0.1$ gives $P_A = 3$ at the plasma center, while the case with $M_{i0} = 0.5$ gives $P_A = 150$, indicating formation of a strong asymmetric structure at the LFS. Figure 5.17 shows the dependence of the W density profile. The W density profile in the high Mach number cases have the same kind of the structure as in the previous section. The W density profile flattens due to the turbulent diffusion, decreasing the neoclassical transport in the $M_{i0} = 0.1$ case. In addition, in the region near $\rho \sim 0.6$, where outward W convection is driven, the enhancement by the toroidal rotation gives impurity ejection. The poloidal asymmetry of the W density distribution as in Eq. 2.70 scales to M_Z^2 , and the W density is localized in the central region, so the value of P_A is considerably larger at high Mach numbers. Figure 5.18 shows the dependence of the W normalized density gradient on the Mach number of W, plotting at $\rho = 0.2$, where the W density peaking is strong, $\rho = 0.4$, where it is close to the ITB-foot, and $\rho = 0.6$, where the outward particle pinch is driven. The value of the normalized W density gradient tends to decrease in accordance with decreased toroidal rotation velocity, especially at $\rho = 0.2$. Except for the higher Mach numbers, the value of $V_{N,asym}^{NCLASS}$ and $V_{T,asym}^{NCLASS}$ decrease with decreasing Mach number. At $\rho = 0.6$, the value of the W normalized density gradient decreases from $M_Z \sim 3.3$ to 1, but at $\rho = 0.4$, the value of the the W normalized density gradient does not decrease due to a delicate balance of $V_{N,asym}^{NCLASS}$ and $V_{T,asym}^{NCLASS}$. On the other hand, the normalized W density gradient is unchanged or even decreased at higher Mach numbers. This is due to the suppression of the temperature screening coefficient H/K via the M_i dependence of k_i as shown in Ref. 125. In Ref. 169, it is also theoretically shown that the W transport is suppressed in the high Mach number regime. In the low Mach number case, however, the inward neoclassical convection of the W is driven in the plasma center region, and the accumulation is enhanced in the ITB region.

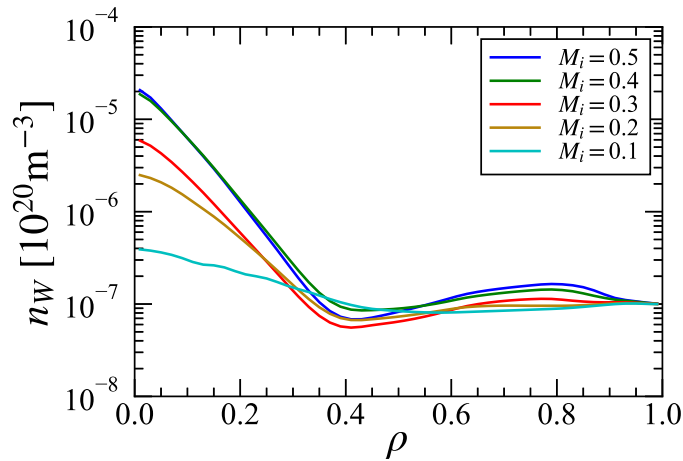


Figure 5.17: Comparison of the radial W density profiles, on the toroidal rotation velocities.

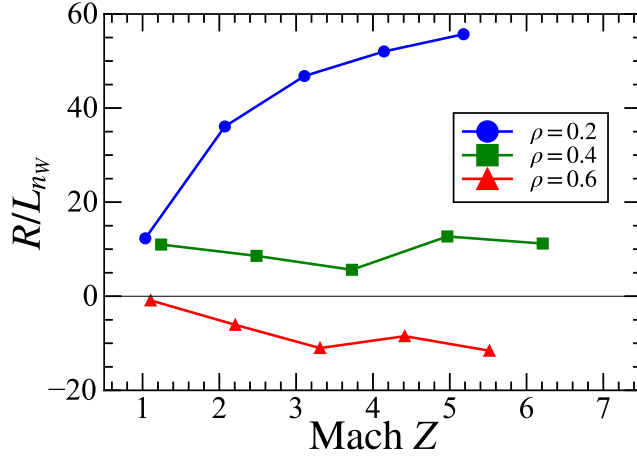


Figure 5.18: Dependence of the normalized W density gradient on Mach number at several radial positions $\rho = 0.2$ (circles), 0.4 (squares), 0.6 (triangles).

5.3.2 ITB position dependence

The dependence on ITB position is investigated in the advanced scenario case. The plasma parameter in Sec. 5.2 is used, and the LH heating profile is changed with heating position $\rho_0^{LH} = 0.4 - 0.7$. The obtained plasma profiles show formation of the ITB, and the radial position of the ITB-foot is consistent with the LH heating position ρ_0^{LH} (Fig. 5.19). In the central region of the plasma, the ion temperature profile is peaking, but the electron temperature is flat due to the effect of radiation losses caused by W accumulation. The global plasma performance is improved to have $H_{98y2} = 1.39$ and normalized- β value $\beta_N = 2.59$ for $\rho_0^{LH} = 0.7$, compared to $H_{98y2} = 1.25$ and $\beta_N = 2.39$ for $\rho_0^{LH} = 0.4$. As shown in Fig. 5.20, the electron density at the plasma center is changed little with that at the ITB position, while the volume-averaged electron density increases from $\langle n_e \rangle = 3.79 \times 10^{19} \text{ m}^{-3}$ to $\langle n_e \rangle = 3.85 \times 10^{19} \text{ m}^{-3}$ for $\rho_0^{LH} = 0.4$ to 0.7 .

The peak of the normalized gradient of the bulk ion in the ITB region is $R/L_{n_i} \sim 2.0$ and $R/L_{T_i} \sim 14.8$ in the $\rho_0^{LH} = 0.4$ case, whereas $R/L_{n_i} \sim 1.5$ and $R/L_{T_i} \sim 16.9$ in the $\rho_0^{LH} = 0.7$ case, with a variation of 24% decrease in the normalized density gradient and 14% increase in the normalized temperature gradient. As shown in Eqs. 4.5 and 4.6, the magnitude of the plasma gradient in the transport improvement region is strongly dependent on the background source. The W accumulation inside the ITB is suppressed in the case of the ITB formed in the outer region. In the case where the LH heating position ρ_0^{LH} is at the outer region, the particle source due to NBI is relatively small, so the outward $V_{T,asym}^{NCLASS}$ driven by the main ion temperature gradient suppresses the inward particle pinch. On the other hand, the case where the ρ_0^{LH} is close to the center of plasma, the inward $V_{N,asym}^{NCLASS}$ causes a sharp increase in the W density gradient at the ITB region (Fig. 5.13). At the peak position of the normalized ion temperature gradient in the ITB region, the $V_{N,asym}^{NCLASS}$ decreases by 55% and the $V_{T,asym}^{NCLASS}$ increases by 17% in the $\rho_0^{LH} = 0.7$ case, compared to the $\rho_0^{LH} = 0.4$ case. In addition, here the Mach number of tungsten decreases by 29% because the toroidal rotation is weakened in the outside of the plasma, so P_A decreases by 62%. The normalized W density gradient is $R/L_{n_W} \sim 23.3$ in the $\rho_0^{LH} = 0.4$ case, while $R/L_{n_W} \sim -10.8$ in the $\rho_0^{LH} = 0.7$ case. The concentration of W, $c_W = \langle n_W \rangle / \langle n_e \rangle$, also decreases from 1.37×10^{-5}

(for $\rho_0^{LH} = 0.4$) to 0.89×10^{-5} (for $\rho_0^{LH} = 0.7$). The W accumulation over the entire plasma can be suppressed by increasing the temperature screening effects.

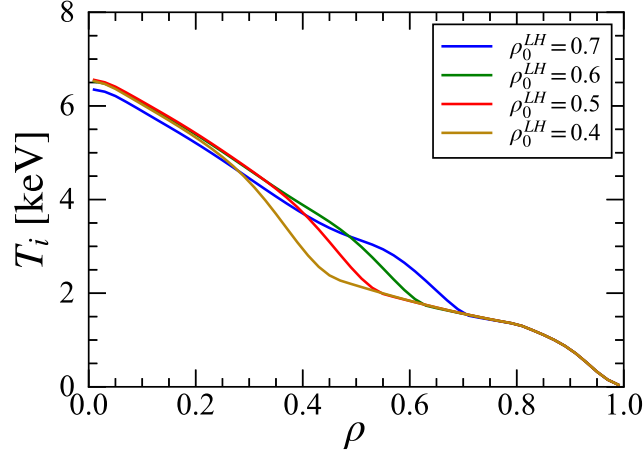


Figure 5.19: Change of the radial ion temperature profile depending on the LH heating position.

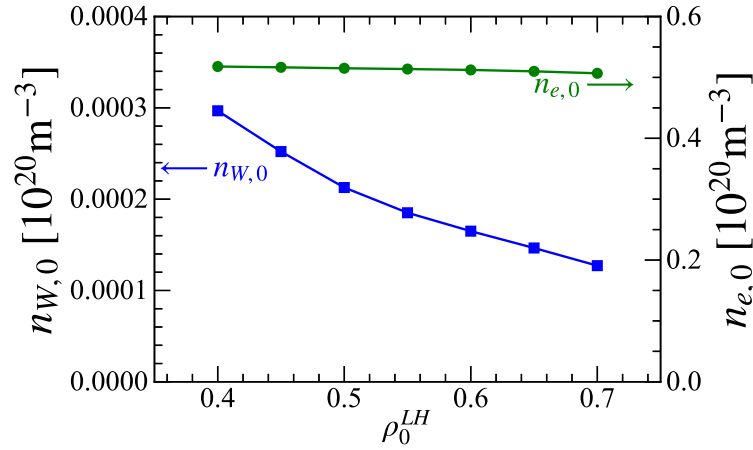


Figure 5.20: Dependence of the central W and electron density on the LH heating position.

5.4 Summary of the chapter

In this chapter, the simulations are performed for the H-mode plasma, the hybrid scenario plasma, and the advanced scenario plasma in a tokamak device. A neoclassical transport model including the effect of the toroidal rotation is applied to impurity transport analyses, because the large toroidal rotation is observed in medium-sized tokamaks. Comparisons between models including poloidal asymmetry in the NCLASS code and analytical models using the PHZ/ E_r pinch are presented. The calculated W profiles in H-mode plasmas using the NCLASS code show good agreements with experiments and first-principle simulations. Simulations with the NCLASS code and the additional PHZ/ E_r pinch model, which do not

include the poloidal asymmetry, could not successfully reproduce the experimental results. In the plasma center, W accumulation can be enhanced due to the toroidal rotation effects, especially in the improved confinement mode where turbulent transport is suppressed. In addition, the accumulation of impurities becomes larger, because radiation losses cause an increase in the bulk plasma density gradient with a decrease in the bulk plasma temperature gradient. For smaller toroidal rotation velocity, the asymmetry of W is weakened, and the neoclassical transport is also reduced to suppress the accumulation. By reducing the Mach number of the bulk ion to about 0.1, the geometry factor P_A , which enhances the neoclassical transport, is estimated to be about 3. In ITER, the central Mach number is predicted to be $M_{i,0} \sim 0.1$ [125], so the effect of enhanced neoclassical transport of W due to toroidal rotation is expected to be smaller. On the other hand, the direction of particle convection is determined by the balance between the main ion density gradient and the temperature gradient, and impurity accumulation may be enhanced in the improved confinement mode plasma where the ITB is formed. The ITB plasma simulation in ITER predicts W accumulation in the plasma center [173], and it is necessary to control the plasma profiles to reduce the impurity accumulation. The present simulations show that the ITB formed at the outer region of the plasma can achieve higher plasma performance, including the suppression of impurity accumulation. The key mechanism is that the plasma density and temperature gradients in the ITB region are strongly dependent on the background source, and since the sources vary in the radial direction, a transport analysis of plasmas and impurities in the entire radial direction is required.

In this chapter, the transport simulation is carried out with a focus on the neoclassical transport of W. However, it is known that turbulent transport and MHD activity have strong effects in the plasma periphery. In addition, it has been reported that minority heating with ICRH can enhance the neoclassical temperature screening effects [70], and a model considering the background electrostatic field is needed. In future work, we will establish a simulation scheme including a turbulence model of impurity and simulate the entire plasma discharge including impurity control with RF heating.

Chapter 6

Summary

Understanding improved confinement plasma and impurity transport is an essential issue for the future realization of fusion power in terms of improving plasma performance, maintaining steady-state burning plasma and impurity control. To control the tungsten accumulation in the core plasma, it is essential to understand the mechanism that leads to the accumulation. In high-temperature plasmas, transport processes and power balance of multiple ion species must be considered. Transport processes include collisional and turbulent transport processes, while power balance includes the interaction between the plasma main components and impurities with different charge numbers due to radiation losses. In these processes, to predict the behavior of the plasma, including impurities, the profile formation mechanism must be investigated in a self-consistent manner. Integrated modeling simulations, which combine several models to include various processes such as equilibrium, transport and heating, enable fast calculations. In this study, a transport simulation scheme including impurities has been developed in the integrated transport code TASK to establish a physical model with fast and accurate calculations and to enable reliable prediction of plasma phenomena. The simulation scheme is then applied to improved confinement mode plasmas, and the mechanism of plasma and impurity profile formation and the actuator for effectively controlling impurity accumulation are clarified. In this study, the following results can be obtained by numerical simulations.

In Chapter 3, at first, the transport simulation scheme is applied to a L-mode plasma, PLATO plasma, and predictions are presented on the plasma performance and the effects of impurities. Several turbulent transport models are used to predict the plasma profile, resulting in a monotonous plasma profile as observed in L-mode plasmas in other devices. Comparison of turbulent transport models is also performed, and the CDBM and mBgB models result in similar plasma profiles. Since there is no external heating mechanism for the ions in PLATO, the ion temperature decreases relative to the electron temperature. The electron temperature is more than three times higher than the ion temperature because of Ohmic heating. Furthermore, due to the decreased ion temperature, results of impurity accumulation in the plasma center region are observed. Increasing the ion temperature is important for PLATO devices, which are aimed for turbulence observation and may prevent the accumulation of impurities. The dependence of plasma performance on the externally controllable parameters of particle feeding rate and vertical magnetic field coil current was investigated. Through these evaluations, it was confirmed that increasing the plasma density and increasing the plasma volume increases the ion temperature. Since there is no direct ion heating mechanism, the energy increased by the collision of ions with electrons plays an

important role.

In Chapter 4, next, simulation schemes are applied to high- β_p mode and negative magnetic shear mode plasmas to investigate the mechanisms of ITB formation and tungsten profile formation. In the high- high- β_p mode, the on-axis heating and the decrease in the magnetic shear due to the bootstrap current decreased the turbulent transport, resulting in the formation of ITBs. The W accumulation in the plasma center region does not occur due to the outward neoclassical convection driven by the ion temperature gradient due to the presence of the heating source. In the negative magnetic shear mode, on the other hand, the magnetic shear is negative due to the peak in the plasma current density profiles, reproducing the ITB plasma as seen in the experiment. Here, the ITBs are formed transiently and the peak in the plasma current density is reduced by current diffusion, resulting in the disappearance of the ITBs. The W density profiles in the ITB region are determined by the neoclassical particle pinch, which strongly depends on the background plasma profile and, therefore, on the particle source and heating source profiles, which strongly influence the gradient of the background plasma profiles. The W density increases significantly when the particle source enhances the ion density gradient, while the increase in W density is suppressed when the heating source enhances the ion temperature gradient. Therefore, since the region where the improved transport occurs is strongly dependent on the source profile, it is important to note that not only the ITB region but also the pedestal region can lead to the contamination of impurities.

In Chapter 5, the impact of toroidal rotation on impurity transport is investigated using a neoclassical impurity transport model that considers the poloidal asymmetry. A comparison with experimental H-mode plasmas confirms the validity of the impurity transport model, and also the validity of the plasma transport simulation scheme including the impurity. Transport simulations in H-mode and improved confinement modes, hybrid scenario and advanced scenario plasmas, are then performed to investigate the formation mechanisms of the plasma and W profiles in each of these modes. In addition, the dependence of impurity transport is investigated to advance the understanding of the transport mechanism. First, the dependence of the W transport on toroidal rotation velocity is evaluated, and it is confirmed that as the toroidal rotation velocity decreases, the W asymmetry weakens, neoclassical transport also decreases, and accumulation is suppressed. In large fusion devices, the toroidal rotation of the core plasma is predicted to be decreased, and the enhancement of W accumulation due to the centrifugal effect is expected to be reduced, as indicated by the present results. Next, the dependence of W transport on ITB position is investigated, and it is shown that forming ITBs in the outer region of the plasma increases the outward convection component due to the main ion temperature gradient, suppressing impurity accumulation and also achieving higher plasma performance. In order to achieve the condition in which impurity accumulation is reduced, the particle source and heating source profiles of the background plasma are found to be important factors, as discussed in Chapter 4. In the improved confinement plasma here, the inward particle convection of tungsten due to the bulk ion density gradient is decreased in the plasma periphery because the particle supply by the NBI is decreased. Conversely, particle supply by pellets or other sources in the plasma periphery could increase the risk of impurity contamination. The plasma profiles should be adjusted to a state where the temperature screening effect due to the bulk ion temperature gradient is relatively stronger than the particle pinching effect due to the bulk ion density gradient. Finally, the dependence of the impurity transport on particle species is investigated. For lighter particles such as carbon and neon, the poloidal asymmetry due to the toroidal

rotation does not occur, and the neoclassical convection is decreased. In the ITB region, however, low-Z and medium-Z impurities are observed to undergo central accumulation due to the neoclassical transport. As shown in Chapter 4, high-Z impurities such as tungsten could also accumulate without the effect of toroidal rotation in ITB plasmas, so impurity control is necessary.

In order to develop an operational scenario for burning plasmas, including impurity control, several models need to be introduced or improved in future studies. One is the neoclassical transport model of impurities, including the asymmetry generated by temperature anisotropy due to ICRH. This asymmetry possibly suppresses the impurity accumulation, and ICRH is an efficient heating source of ions, so it is an important factor in the study of impurity control methods. Next, a turbulent particle transport model for main ions and impurities needs to be introduced. In this study, the turbulent diffusion coefficients of impurities were determined based on experimental and analytical results, but basically, turbulent transport is determined by the plasma and impurity conditions. Behavior of light impurity species, in particular, are determined by turbulent transport rather than neoclassical transport dominantly, similar to the bulk ions. It has been reported that inward flow driven by turbulence increases impurity accumulation, and it is known that turbulent transport cannot be neglected in the plasma periphery, even for tungsten. In addition, it is necessary to introduce a model of impurity transport with MHD events and to couple the SOL and core regions for self-consistent determination of impurity inflow.

In summary, since impurities in the core could degrade plasma performance, a transport simulation scheme in the integrated code TASK is developed, enabling to predict the profile of both plasma and impurity in a self-consistent manner with low computational cost. Then, for the first time, integrated simulations of plasma transport were performed for tungsten transport considering the poloidal asymmetry in the improved confinement mode. The mechanism is clarified to determine behavior of multi-species ions by coupling plasma profile formation with transport improvement. The integrated code developed in this study will contribute to developing fusion operation scenarios and understanding plasma and impurity transport as a tool for evaluating the evolution of plasma phenomena.

Appendix A

A.1 Particle feeding rate control with PLATO parameter

The dependence on the particle feeding rate is evaluated, as it is externally controllable in PLATO experiments. The plasma parameters and coil currents are same with those in the previous section. The particle feeding rate is adjusted to give the density at the center from $n_0 = 0.5$ to $1.5 \times 10^{19} \text{ m}^{-3}$. Figure A.1 indicates that the electron temperature decreases and the ion temperature increases, as the density increases. This dependency can be explained with the power balance. The power balance of electrons and ions are presented with the ohmic power, energy transfer from electrons to ions and heat fluxes in the steady state as

$$\text{Electron} : -\frac{1}{V'} \frac{\partial(V' Q_e)}{\partial \rho} + P_{OH} - P_{EQ} = 0, \quad (\text{A.1})$$

$$\text{Ion} : -\frac{1}{V'} \frac{\partial(V' Q_e)}{\partial \rho} + P_{EQ} = 0. \quad (\text{A.2})$$

Figure A.2 shows the radial profiles, and the values at $\rho = 0.3$, where the temperature gradient is large, are picked up in Fig. A.3 for the explanation. A larger density leads increase of electron flux and P_{EQ} , which makes the electron temperature lower. Equation 2.4.2 gives

$$P_{EQ} \propto n^2 T_e^{-1.5} (T_e - T_i). \quad (\text{A.3})$$

Electron-ion collision is responsible to the energy exchange between electrons and ions, which tends to make difference of the temperature between electrons and ions smaller, as in Fig. A.4. Increases of n^2 and $T_e^{-1.5}$ and decrease of $T_e - T_i$ determines the magnitude of P_{EQ} , as is indicated with the red dashed line in Fig. A.3. Therefore, the ion temperature increases, as P_{EQ} increases. In addition, the larger plasma resistivity gives P_{OH} increase. Equations 2.17, and the spitzer's formula for the resistivity

$$\eta = \frac{m_e}{ne^2 \tau_{ei}}, \quad (\text{A.4})$$

give the following relation

$$P_{OH} \propto T_e^{-1.5} j^2. \quad (\text{A.5})$$

The deviation from this relation indicated by the green dashed line in Fig. A.3 comes from density dependence of neoclassical resistivity by NCLASS routine.

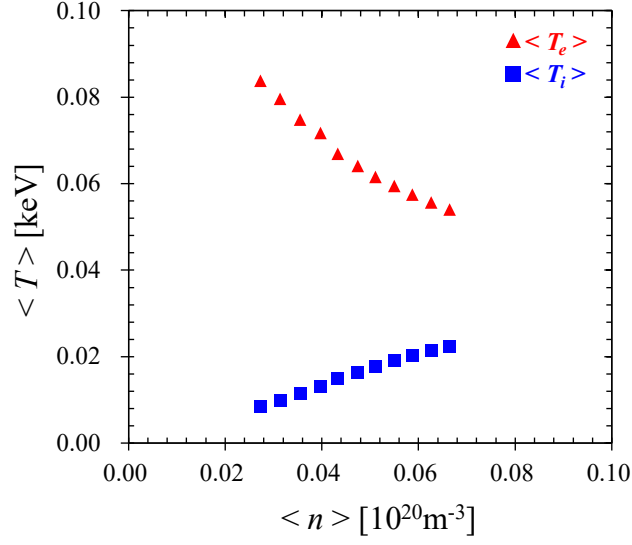


Figure A.1: Relationship between volume-averaged temperature and density, when the particle feeding rate is changed.

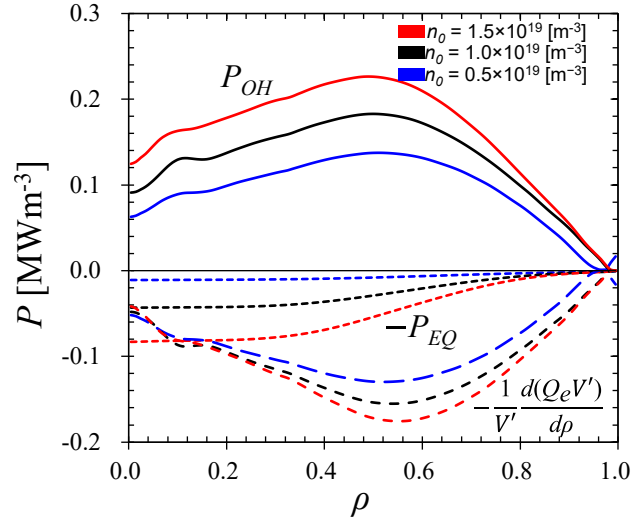


Figure A.2: Radial profile of ohmic power, equipartition power and electron heat transport term. The cases with $n_0 = 0.5$ to $1.5 \times 10^{19} \text{ m}^{-3}$ are shown.

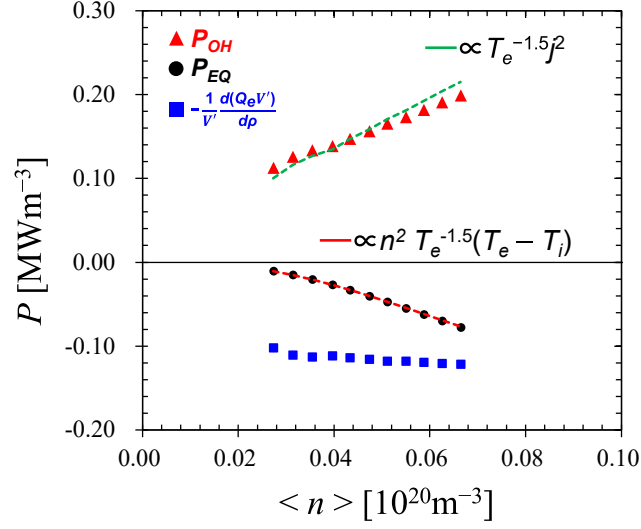


Figure A.3: Volume averaged density dependence of terms at $\rho = 0.3$ in the power balance.

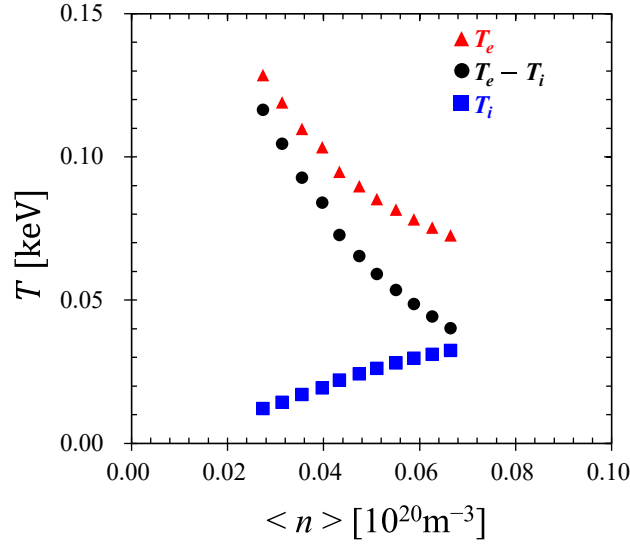


Figure A.4: Volume averaged density dependence of electron and ion temperature at $\rho = 0.3$, and temperature difference.

A.2 Coil current control with PLATO parameter

The other transport analysis is performed with variation of an external coil current. The plasma parameters are same with the previous sections excepts D-coil currents, which are changed in the range from -18.8 to -23.3 kA turn, while the other coil currents are same. The particle source is set to be peaked at $\rho = 0.7$. The plasma size V_p is changed as in Fig. A.5, in accordance with the D-coil current. In the transport model of the TR code, the effect of plasma shape is introduced by the plasma volume, metric $\langle |\rho| \rangle$ and $\langle |\rho|^2 \rangle$. As shown in Fig. A.6, the density slightly decreases with increase of plasma volume. Particle source

$S(r)$ decreases as the total particle feeding rate is constant, but the density variation is slight, due to decrease of the particle fluxes.

On the other hand, the larger the plasma shape is, the lower the electron temperature and the higher the ion temperature become. Figure A.7 shows the radial profiles of the terms in the power balance of electrons. The values at $\rho = 0.3$ are picked up in Fig. A.8 for explanation of the dependency as well. Since the plasma current is kept constant, the current density decreases in accordance with the increase of the plasma volume. Power P_{OH} is represented as

$$P_{OH} \propto T_e^{-1.5} A_p^{-2}. \quad (\text{A.6})$$

where A_p is the size of the poloidal cross-section of the plasma. Therefore, P_{OH} is decreased depending on A_p as in Fig. A.8, which gives the decrease of the electron temperature. P_{EQ} decreases slightly as A_p increases, but the energy confinement time becomes longer as a result of the small change in the thermal diffusivity and the increase of the plasma minor radius as in Fig. A.9. Therefore, the ion temperature increases due to the improvement of confinement.

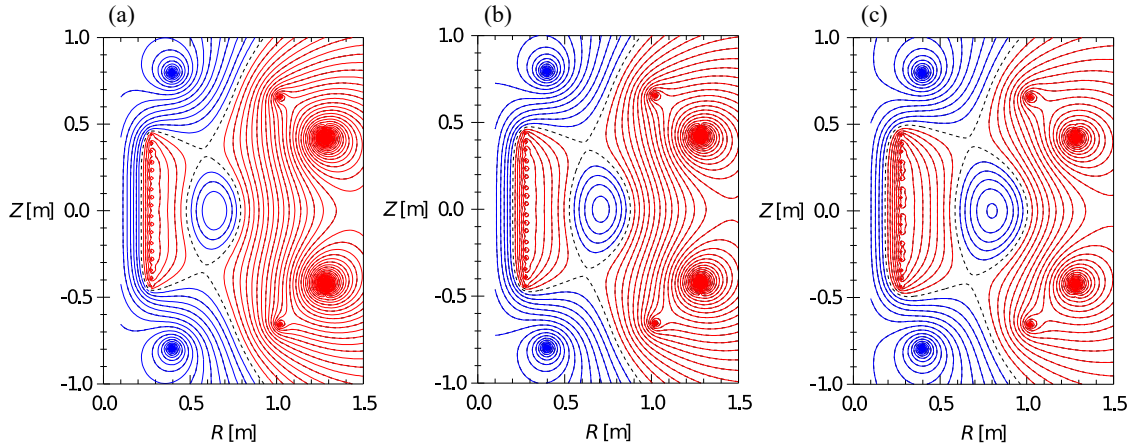


Figure A.5: Equilibrium magnetic configurations with D-coil current $I_{D-coil} =$ (a) -28.8, (b) -23.3, and (c) -18.8 kA turn.

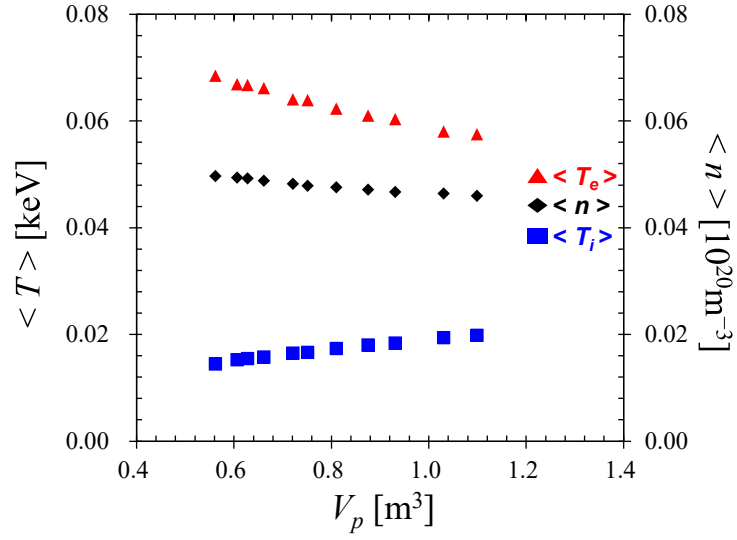


Figure A.6: Relationship between volume-averaged density, ion and electron temperature, and the plasma volume. This is the case when I_{D-coil} is changed.

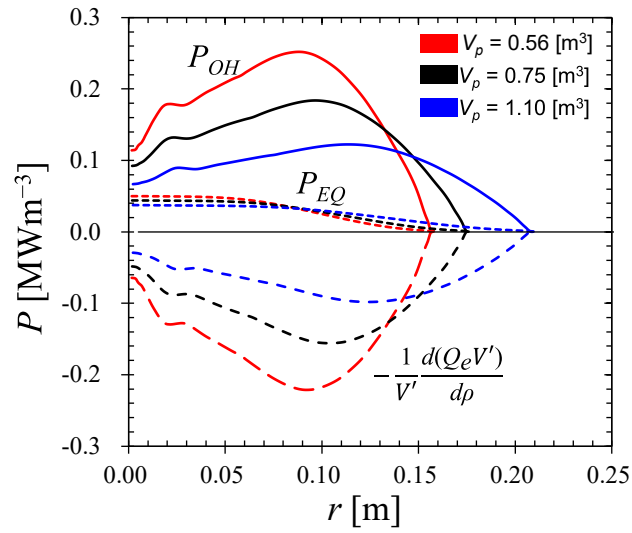


Figure A.7: Radial profile of ohmic power, equipartition power and electron heat transport term.

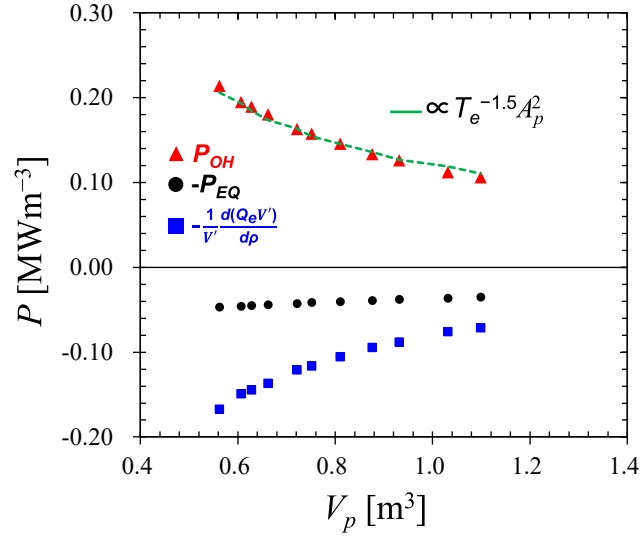


Figure A.8: Plasma volume dependence of terms at $\rho = 0.3$ in the power balance.

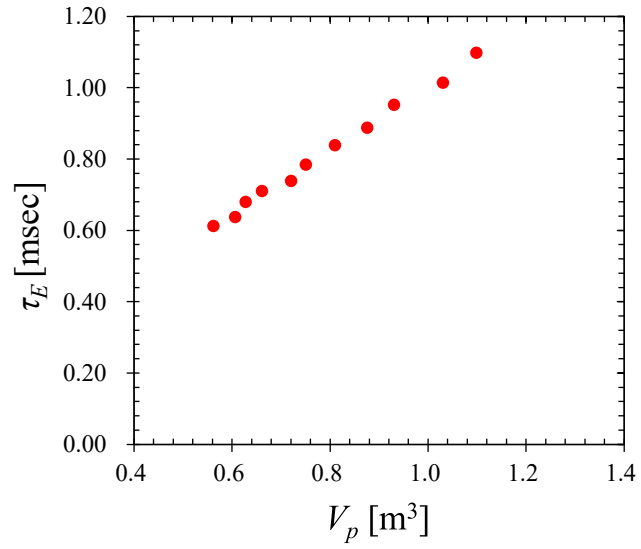


Figure A.9: Plasma volume dependence of the energy confinement time.

Appendix B

B.1 Dependence of particle species

In this section, the transport model that includes toroidal rotation effects for medium- Z and low- Z impurity species is evaluated to investigate the dependence of neoclassical transport on the particle species. The simulations are performed by referring to the negative magnetic shear mode discharge (JET#51976 at $t = 46.8$ s) shown in Ref. 163 to focus on the impurity transport process. The plasma density and temperature profiles shown in Fig. B.1 are used. The ITB of ion temperature is formed at $0.45 < \rho < 0.85$, and the bulk ion density gradient is also increased inside $\rho \sim 0.75$, where it is the position of the ITB-foot in the ion temperature profile. The comparison with experimental impurity profiles in this simulation is qualitative due to the high uncertainties in the plasma gradient, as described in Ref. 163. The bulk ion density is determined from $n_i = n_e - \sum Z n_Z$ to satisfy the quasi-neutrality condition since the impurity density is about 15% of the electron density. Fixed plasma profiles are used except for the bulk ion density. The safety factor profile, which has a negative magnetic shear, is determined by the current density profile with the function shown in Ref. 170, which is modeled to reproduce the hollow current density profile of discharge at JET#51976. The toroidal velocity profile is given by the function form with $V_{tor}^{\rho=0} = 250$ km/s at the plasma center, referring to Ref. 171. As the impurity species, carbon (C, $A = 12$, $Z_{max} = 6$), neon (Ne, $A = 20.18$, $Z_{max} = 10$), and nickel (Ni, $A = 58.69$, $Z_{max} = 28$) are assumed, and their toroidal velocity and Mach number are determined by Eqs. 2.71 and 2.72, respectively. The radial profiles of the Mach number for each impurity are shown in Fig. B.2. The radial profiles of the Mach number for each impurity are shown in Fig. B.2, and at the plasma center, $M_{D0} = 0.16$, $M_{C0} = 0.40$, $M_{Ne0} = 0.52$, and $M_{Ni0} = 0.87$. Since the Mach number decreases for lighter atoms, C decreases by about 1/2 compared to Ni. The turbulent diffusion coefficients of the impurities are referred to in Ref. 164 and have the same value for all impurity species. The following turbulent diffusion coefficient, a combination of constant and parabolic ones, is used.

$$D_W^{turb} = \begin{cases} C_1^{W,turb} & (0.0 \leq \rho \leq \rho_1^{W,turb}) \\ C_2^{W,turb} & (\rho_1^{W,turb} \leq \rho \leq \rho_2^{W,turb}) \\ \frac{C_3^{W,turb} - C_2^{W,turb}}{1.0^2 - \rho_2^{W,turb^2}} \rho^2 + \frac{C_2^{W,turb} 1.0^2 - C_3^{W,turb} \rho_2^{W,turb^2}}{1.0^2 - \rho_2^{W,turb^2}} \rho^2 & (\rho_2^{W,turb} < \rho \leq 1.0) \end{cases} \quad (B.1)$$

where, $\rho_1^{W,turb} = 0.2$, $\rho_2^{W,turb} = 0.5$, $C_1^{W,turb} = 1.0$, $C_2^{W,turb} = 0.1$, and $C_3^{W,turb} = 2.0$ are used. The turbulent diffusion coefficient of impurity is decreased in the ITB formation region, where the negative magnetic shear is formed, as confirmed by the analysis [172], and it is reproduced in the above formula. The impurity convection velocities of the turbulent are ignored. The density of impurity at the boundary is set as $n_{C,edge} = 3.0 \times 10^{17} \text{ m}^{-3}$, $n_{Ne,edge} = 0.7 \times 10^{17} \text{ m}^{-3}$, $n_{Ni,edge} = 0.1 \times 10^{17} \text{ m}^{-3}$.

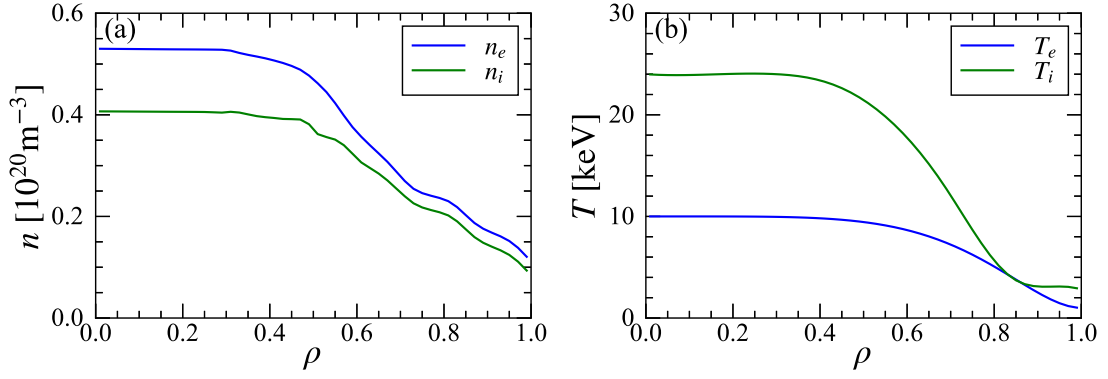


Figure B.1: Referenced radial profiles of (a) electron and ion density, (b) electron and ion temperature.

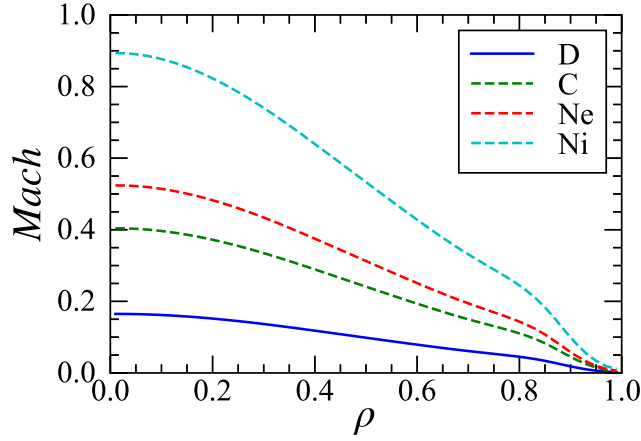


Figure B.2: Radial profiles of Mach number of the main ion (solid line), and the C, Ne, Ni (dashed line).

Figure B.3 shows the FSA in the radial direction and the 2D impurity profiles, with results almost consistent with the impurity densities from the experiments presented in Ref. 163 and Ref. 164. As Figs. B.3(d)-(f) show, the poloidal asymmetry is not observed, unlike the W case. It is also shown in Fig. B.4 that V_{sym}^{NCLASS} and V_{asym}^{NCLASS} almost agree with each other. Since the turbulent diffusion is increased in the plasma center region in this condition, the asymmetric structure does not appear even in the plasma center region where the Mach number is relatively large. On the other hand, as in the bulk ion density profile, the impurity density increases for all impurity species at the inside of the ITB-foot ($\rho \sim 0.75$), resulting in a step-like profile. The impurity density increases significantly for heavier impurities, and $R/L_{n_C} \sim 15$, $R/L_{n_{Ne}} \sim 25$, and $R/L_{n_{Ni}} \sim 57$ at $\rho = 0.55$ are obtained at $\rho = 0.55$, where the normalized impurity density gradient peaks. The increase in impurity density is due to the neoclassical inward particle pinch, and since the collision frequency is usually proportional to Z^2 , the neoclassical transport is increased for higher Z impurities. The neoclassical

convection velocities of Ne and Ni for the C are, respectively, $\frac{V_{N,asym}^{Ne}}{V_{N,asym}^C} \sim \frac{V_{T,asym}^{Ne}}{V_{T,asym}^C} \sim 1.7$, $\frac{V_{N,asym}^{Ni}}{V_{N,asym}^C} \sim \frac{V_{T,asym}^{Ni}}{V_{T,asym}^C} \sim 3.6$, where both $V_{N,asym}^{NCLASS}$ and $V_{T,asym}^{NCLASS}$ are increased for heavier particles.

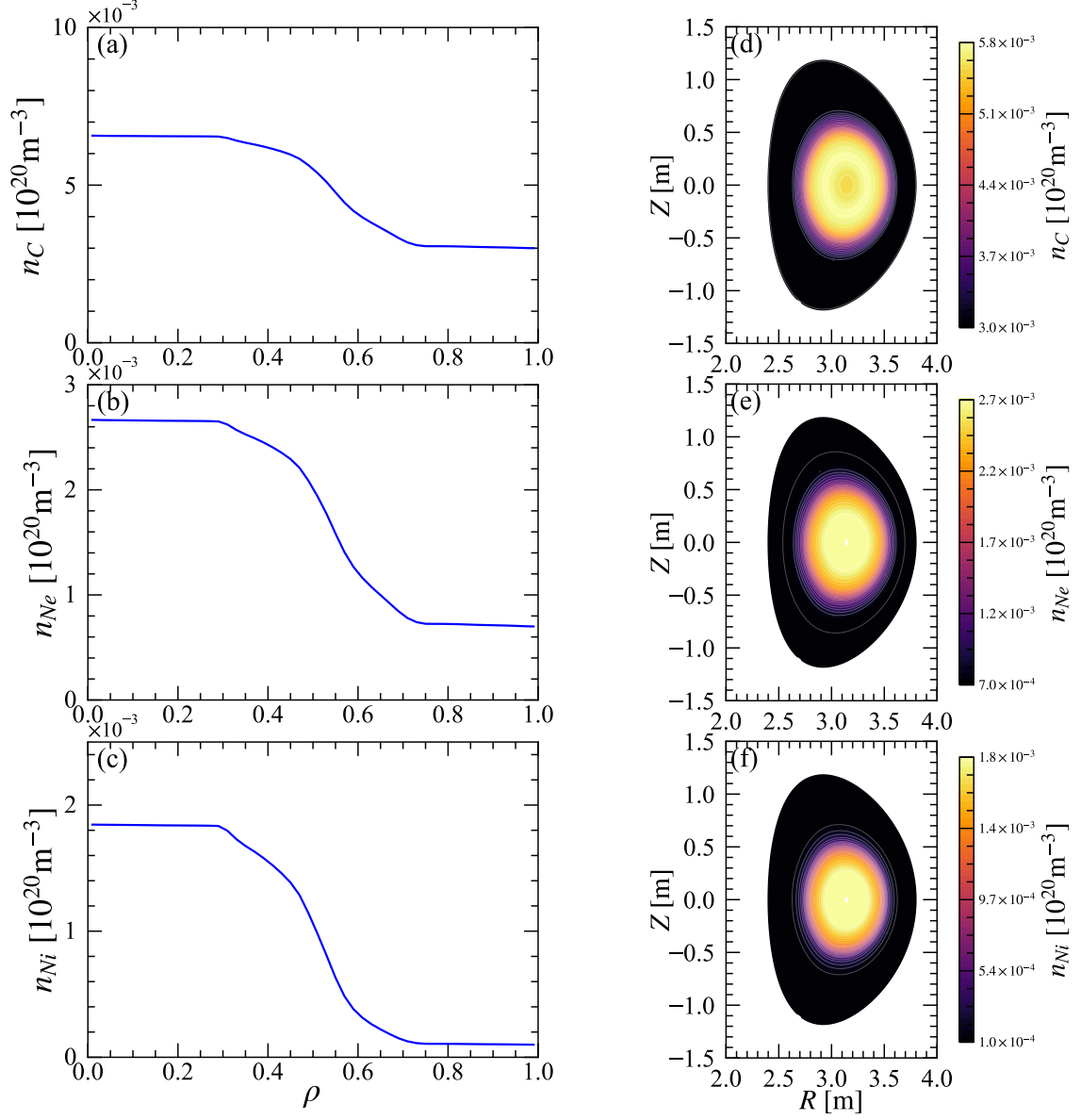


Figure B.3: Radial profiles of the density of (a) C, (b) Ne, (c) Ni of LFS value on the midplane, and calculated 2D plots of density of (d) C, (e) Ne, (f) Ni on the poloidal cross-section.

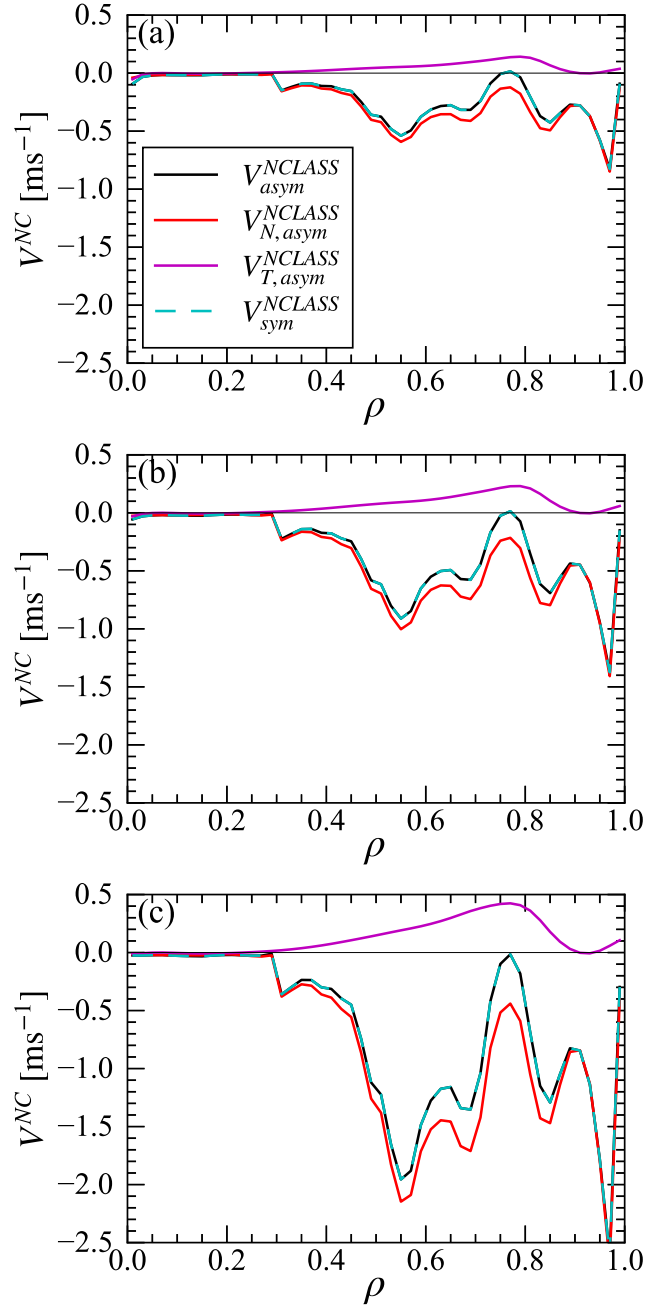


Figure B.4: Radial profiles of neoclassical convective velocities of (a) C, (b) Ne, (c) Ni.

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