Transport barriers in non-axisymmetric magnetic fields

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Abstract. A transport barrier which is characterized by the steeper temperature or/and density gradient has been observed both in the core region (internal transport barrier : ITB) and edge region (edge transport barrier : ETB) in the plasmas with non-axisymmetric magnetic fields compared with the gyro-Bohm type L-mode transport. A reduction of thermal conductivity and particle diffusivity in these plasmas compared with the L-mode levels is observed and the parameter dependence of thermal conductivity and particle diffusivity changes significantly from that in L-mode. A negative or at least a weak dependence of the thermal diffusivity on temperature is observed in the heat transport barrier, while a strong temperature dependence (temperature to the power of 1.5) is normally observed in the L-mode region. In this paper, transport barriers that have been observed in non-axisymmetric magnetic field configuration are discussed from the view point of the parameter dependence of transport coefficients, not the shape of temperature or density profiles. The transition between the L-mode and the improved mode with a transport barrier is also discussed based on the idea of a bifurcation of multi-state transport (discrete transport states with different relations of heat flux to temperature gradient that are self-consistent with the different turbulence states realized).

1. Introduction

The physics mechanism of transport barriers in confinement systems with non-axisymmetric magnetic fields has been studied since the high confinement mode (H-mode) and internal transport barrier (ITB) were experimentally observed in stellarators and heliotron devices. In the H-mode, steep pressure gradients and $E \times B$ flows localized near the plasma edge were observed in Wendelstein-7AS[1, 2, 3], CHS[4, 5, 6], LHD[7], Heliotron-J[8], TJ-II stellarator[9], H-1 heliac[10] and Tohoku-heliac[11, 12, 13]. The H-modes in these devices have been found to be sensitive to the edge rotation transform
(iota). Clear edge iota windows for the H-mode transition were observed in Wendelstein 7-AS and Heliotron-J, which has not been observed in the H-mode in a tokamak. Although there are some differences in the trigger mechanism of the H-mode, the role of the $E \times B$ shear flow on the suppression of turbulence and the reduction of transport is similar to that in a tokamak H-mode. The details of the characteristics of an H-mode in the confinement system with a non-axisymmetric magnetic field has been discussed[14].

In non-axisymmetric confinement devices, both the electron and ion ripple losses are sensitive to the radial electric field and electron and ion radial fluxes are functions of the radial electric field in the plasma. Since the electron radial flux is equal to the ion radial flux in the steady state condition, the radial electric field which gives identical ion and electron radial flux is called a "root" of ambipolar condition ($\Gamma_i = \Gamma_e$). The radial electric field with positive (negative) value is called electron (ion) root, because the electron (ion) flux exceeds the ion (electron) flux at zero radial electric field. Because the magnitude of the radial electric field is much larger in the electron root, the reduction of neoclassical transport due to the radial electric field is expected in the electron root rather than in the ion root. Therefore, the transition of the radial electric field is triggered by the neo-classical non-ambipolar ion and electron fluxes, and the radial electric field becomes positive (in the electron root) when the plasma collisionality is low enough, while it is negative (in the ion root) at higher collisionality[15, 16]. Associated with a transition of the radial electric field from negative (ion root) to large and positive (electron root), the reduction of the electron thermal diffusivity is observed in CHS[17], Wendelstein-7AS[18] and LHD plasmas[19, 20, 21, 22]. The coupling between the radial electric field shear and the turbulence levels has been investigated to understand the anomalous transport in the TJ-II stellarator [23].

Recently zonal flows have been identified experimentally using a heavy ion beam probe (HIBP) and probes and the coupling between turbulence and zonal flows has been investigated in CHS[24, 25] and in the H-1 Heliac[26, 27] as well as in tokamaks[28, 29, 30, 31, 32]. The physics mechanism for the formation of edge and internal transport barriers has been studied. In this paper, the physics mechanism to determine the mean $E \times B$ flow and its shear, zonal flows and spontaneous toroidal flow is described. The transport characteristics, especially the dependence of the heat flux on the temperature and temperature gradient, are discussed in comparison with that in tokamaks. The electron internal transport barrier associated with the transition of the mean $E \times B$ flow and other transport improvements (ion transport and particle transport) are also described.

2. Physical mechanism determining spontaneous, mean, zonal and GAM flows

Transport of the plasma in the optimized non-axisymmetric magnetic field system is governed by the turbulence as it is in the plasma in a tokamak, because the transport due to collisional processes (neoclassical transport) has been reduced to a level below
the transport due to turbulence (anomalous transport). Figure 1 shows the physical mechanism determining the spontaneous toroidal flow, the mean poloidal flow, the stationary zonal flow and geodesic acoustic modes (GAMs) in a non-axisymmetric magnetic field system. Turbulence in the plasma produces a mean flow in the direction of $E \times B$ (nearly the poloidal direction) and in the direction parallel to the magnetic field (nearly the toroidal direction) through the Reynolds stress. These mean flows tend to be localized at the region of the transport barrier, where the temperature an/or density gradient are significantly large. The mean flow in the direction parallel to the magnetic field is called the spontaneous toroidal rotation, while the mean flow in the direction of $E \times B$ is called the poloidal spin-up.

These mean flows are damped by the parallel and perpendicular viscosities [33]. The parallel viscosity is a damping mechanism through collisions between trapped particles (banana particles in a tokamak, ripple trapped particles in a non-axisymmetric magnetic system) and passing particles. This parallel viscosity is the dominant damping process.
in the poloidal direction. The perpendicular viscosity is the damping mechanism through momentum exchange due to turbulence in the region where a velocity shear exists. Therefore this viscosity is called shear viscosity or momentum diffusivity in momentum transport analysis. This perpendicular viscosity is a dominant damping process in the toroidal direction in a tokamak. In non-axisymmetric magnetic systems both parallel and perpendicular viscosities are important for determining the toroidal rotation.

Non-ambipolar radial flux and the neoclassical viscosity tensor are additional mechanisms driving these mean flows. Non-ambipolar radial flux produces a radial electric field and sustains the $E \times B$ mean flow. The non-ambipolar radial flux due to ion and electron ripple loss is a strong driving term of the $E \times B$ flow especially in non-axisymmetric magnetic field systems. In general (except for at the magnetic axis) the magnetic field is oblique to the symmetry direction (toroidal in tokamak and helical in non-axisymmetric system) because of the rotational transform of magnetic field. Because the plasma flow along the symmetry direction has minimum viscosity, the plasma tends to flow along the direction of minimum viscosity (symmetry direction). The neoclassical viscosity tensor (called poloidal viscosity in tokamak) acts as the forces to change the direction of plasma flow towards the symmetry direction. The neoclassical viscosity tensor is significantly large and sometimes larger than the shear viscosity and therefore the mean flow parallel to the magnetic field can be driven by the $E \times B$ flow. Because of the existence of these additional mechanisms driving these mean flows, the generation of the mean flows does not require the existence of turbulence.

On the other hand, zonal flows ($E \times B$ flows with finite frequency and radial wave number) are always generated by turbulence and these flows should vanish if the turbulence completely disappears. Since the zonal flows have no poloidal wave number these flow suppress the turbulence and reduce the radial transport by the $E \times B$ shear effect. The zonal flows have finite frequency and radial wave number and there is no net poloidal momentum when the momentum of the zonal flow is averaged in time and space. This is in contrast to the mean flows which have finite momentum by exhausting the momentum out of the plasma through the plasma periphery. The magnitude of the zonal flows are determined by the driving term due to turbulence and damping due to collisional flow damping. The other damping mechanism is non-linear flow damping which suppresses the growth of the zonal flow by consuming the free energy used to produce a zonal flow. In a non-axisymmetric system, bifurcation of the radial electric field occurs because the $E_r$ dependence of the ion radial flux is different than that of the electron radial flux. As the collisionality is decreased below the critical value, the radial electric field of the plasma changes from small negative (ion-root) to large positive (electron-root) and the radial flux also decreases significantly. The collisional flow damping is sensitive to the radial electric field and the mean $E \times B$ flow in a non-axisymmetric system. The collisional flow damping becomes small in the electron root plasma where the large $E_r$ reduces ripple loss and collisional damping. Therefore a larger zonal flow is expected in the electron root in a non-axisymmetric system.
Figure 2. (a) Spatio-temporal structure of stationary zonal flows evaluated from the correlation function in CHS. (b) Spatial correlation with the time lag of $\tau = 0, 1$ and 2 ms.[34] (c) Evolutions of wavelet spectrum of electric field fluctuations as a function of frequency, $W(E_A; f, t)$. The unit of the color bar is $V^2$/kHz. (d) Evolution of the zonal flow, $Z_A(t)$. The waveform of the zonal flow is the electric field numerically filtered in the frequency range around $\sim 0.5$ kHz. The blue dashed line indicates the evolution of the zonal flow amplitude using wavelet analysis, $A(Z_A; t)$ [35].

3. Coupling between fluctuations and stationary zonal flows and GAMs

Identification of zonal flows in the high temperature core was accomplished in CHS using twin HIIBPs located at two different toroidal positions separated by 90 degrees [24]. This zonal flow has a toroidal symmetry of $n = 0$ with finite radial wavenumbers and has a broad frequency spectrum from $\sim 0.5$ to $\sim 1$ kHz which is in contrast to the zonal flow with a sharp frequency peak at the GAM frequency (10 - 20 kHz) described later. This zonal flow is called stationary zonal flow, however, the flow pattern changes within ms and is not in a steady state because of the coupling between the zonal flow and turbulence. The change in the flow pattern of the stationary zonal flow is clearly demonstrated in the spatio-temporal pattern of the radial correlation function between the electric field at $r = 12$ cm and at another location varying from 10 to 14 cm. This observation shows that zonal flows with a finite radial wavelength of $\sim 1$ cm do exist, and the structure of a zonal flow changes completely in $\sim 2$ ms.

The linkage between a zonal flow and turbulence is examined with wavelet analysis to extract the relation between the intermittent characteristics of the turbulence and the zonal flow phase (or direction). Figure 2(c) shows examples of the evolution of the wavelet power spectrum of electric field fluctuations, $W(E_A; f, t)$. The wavelet analysis
Figure 3. (a) Diagrams of wavelet squared bicoherence near the positive peak (maxima) of the zonal flow and (b) an expanded view of the diagram[25].

successfully reveals the intermittent nature of turbulence manifested in the electric field fluctuations. Figure 2(d) shows the zonal flow component filtered out from the electric field fluctuation. Here the positive sign means that the resultant $E \times B$ -drift is in the ion diamagnetic direction to which the bulk rotation orients itself.

It has been confirmed qualitatively that the turbulence power is modulated by the phase of the zonal flow and that the fluctuation spectra clearly differ from each other at the maxima and minima of the zonal flow $Z_A(t)$. The electric field fluctuations measured with an HIBP have three characteristic frequency ranges for the fluctuation spectrum: zonal flows in the range of less than $\sim 1.5$ kHz, the coherent modes at 19 and 38 kHz which are conjectured as GAMs and broad band turbulence fluctuations in the range around 50 kHz. Wavelet bicoherence analysis is used to evaluate the dependence of the coupling constant on the phase of the zonal flow with the help of a conditional average. As seen in figure 3, a clear couplings between fluctuations and a zonal flow ($f < 1.5$ kHz) has been found near the positive peak (maxima) of the zonal flow. The strong couplings appear along the lines of $f_1 + f_2 \pm 0.5$ kHz, which is clearer in the expanded view in the frequency range (90 - 100 kHz) of the diagram. Couplings between fluctuations and a GAM zonal flow ($\sim 19$ kHz) are also observed in the medium frequency range
Figure 4. Radial profiles of (a) radial electric field and (b) shear in the radial electric field, (c) power spectra of the plasma potential in different radial regions, (d) radial profile of the spectral power density of stationary zonal flows (0.1-0.6kHz) of H-mode plasmas in the H-1 heliac. Hatched boxes indicate radial positions of the $E_r$ maxima in regions I and II[36].

$(30 - 60kHz)$ of the turbulence frequency. These observations clearly demonstrate the coupling between fluctuations and stationary zonal flows and between fluctuations and GAMs.

4. Relation between mean $E \times B$ flow and zonal flow

It is well known that the $E \times B$ flow shear suppresses turbulence in the plasma. This $E_r$ suppression effect causes strong coupling between the turbulence and the mean $E \times B$ flow shear. Because the turbulence in the plasma couples with both the stationary zonal flow and mean $E \times B$ flow shear, the stationary zonal flow is expected to be related to the $E \times B$ flow through turbulence in the plasma. In the H-1 heliac, the contribution of a zonal flow to the spatial modulation of the radial electric field, $E_r$, profile was investigated. Figure 4 shows the radial profile of the radial electric field, its shear and spectral power density of stationary zonal flows measured with Langmuir probes. As
seen in figure 4(c), plasma potential fluctuations with low frequency (0.1-0.6kHz) and with zero toroidal (n) and poloidal (m) mode numbers are dominant in the H-mode and localized near the plasma edge. There are three \( E_r \) regions observed: a slightly positive \( \Delta E_r \) inside the top of the transport barrier, a negative \( E_r \) region (I) and a more negative \( E_r \) region (II). The zonal flow maximum spatially coincides with the maximum in negative \( \Delta E_r \), where the \( \Delta E_r \) shear becomes small. In general, the \( \Delta E_r \) shear tends to reduce the intensity of the zonal flow by the reduction of turbulence (driving terms) due to the shear suppression effect, while the \( \Delta E_r \) itself may enhance the intensity of the zonal flow by reducing the collisional flow damping (damping term). Therefore both the small \( E_r \) shear at the peak of \( E_r \) and the large \( E_r \) contribute to increase the zonal flow intensity as seen in figure 4. However, from these experiments it is not clear whether the \( E_r \) or \( \Delta E_r \) shear contribute to the enhancement of the stationary zonal flow observed.

Figure 5(a) shows the intensity of the zonal flow, \(|\Delta E_{ZF}|^2\), and the normalized fluctuation amplitude, \(|\varepsilon E_r/\nabla T_0|^2\) during the period of the electron internal transport barrier (e-ITB) and in the case without the ITB. A large intensity of the zonal flow and a small fluctuation amplitude are observed in the plasma with an electron internal transport barrier, where large \( E_r \) and large \( \Delta E_r \) shear exists inside \((\rho < 0.25)\) and at the boundary \((\rho \sim 0.3)\) of the ITB, respectively. This is in contrast to the small intensity of the zonal flow and the large fluctuation amplitudes which are observed in the plasma without an ITB. This experimental result clearly demonstrates the role of \( E_r \) in collisional damping rather than the role of \( \Delta E_r \) shear in turbulence suppression. The data in figure 5(a) are measured at \( \rho = 0.2 \) where the \( E_r \) shear is too weak (see the radial profile of the plasma potential) to suppress turbulence even in the ITB. This experimental result clearly demonstrates that the energy of the fluctuations is preferentially transferred into zonal flows when the collisional flow damping becomes weak in the electron ITB.

In order to study the effects of \( E_r \) and \( \Delta E_r \) shear on transport, the radial profiles of thermal diffusivity are calculated with a one-dimensional transport code based on the nonlinear current-diffusive interchange mode using typical \( T_e, T_i, n_e \) profiles. In this code, the \( E_r \) effect on the collisional damping rate and \( E_r \) shear effect on the reduction of thermal diffusivity are included. The radial profiles of \( E_r \) in figure 5(b) are mainly determined by neoclassical non-ambipolar fluxes. The reduction of turbulent transport in the entire region of strong positive \( E_r \) (not just the region of strong \( E_r \) shear) is quantitatively demonstrated by use of transport code analysis. The transport coefficient by bare fluctuations (without Zonal flow effects) does not show a noticeable reduction for \( r < r_{\text{inv}} \), where \( r_{\text{inv}} \) is the minor radius at the boundary between electron root (large positive \( E_r \)) and ion root (small negative \( E_r \)) owing to the increased temperature. By including the screening by Zonal flow thermal diffusivity is predicted to be quenched in the core of plasmas with electron ITB.
5. Temperature and temperature gradient dependence of transport in axisymmetric and non-axisymmetric systems

Figure 6 shows the radial profiles of electron temperature in the plasmas with an electron ITB for the ECH power scan and the radial profiles of thermal diffusivity at the highest ECH powers in JT-60U tokamak and the LHD heliotron. The electron temperature profile becomes peaked after the formation of the electron ITB and the central electron temperature strongly depends on the ECH power. As seen in figure 6(a)(b), the electron temperature profile observed in the electron ITB in LHD is peaked at the plasma axis, while the electron temperature profile observed in JT-60U has a flattened region near the
plasma center. Although the mechanism causing the flattening is not well understood, the flattening shows that the transport at the plasma core is worse than the L-mode level and the transport barrier is localized in the narrow region.

In order to compare the radial structure of heat transport between tokamak and heliotron plasmas with an electron ITB, the electron thermal diffusivity is normalized by the gyro-Bohm scaling of $T_e^3/2e/B^2$ as shown in figure 6(c). The normalized electron heat conductivity in the outer plasma is nearly identical for JT-60U and LHD. The normalized thermal diffusivity decreased towards the plasma center and reaches low levels close to $0.1 \text{ m}^2\text{s}^{-1}/(\text{keV}^3/2\text{T}^-2)$ in the LHD electron ITB. This is in contrast to the radial profile of the normalized electron thermal diffusivity observed in JT60U, where the minimum normalized thermal diffusivity is located at one-third of plasma minor radius ($r/a = 0.35$). It is interesting that the reduction of thermal diffusivity in a tokamak is localized in the narrow region with the $E \times B$ shear, whereas the reduction of thermal diffusivity in a heliotron extends to the plasma core where $E_r$ shear is small but $E_r$ is large enough. These observation are consistent to the predictions of transport modeling including the zonal flow screening effect as seen in figure 5. The difference in the radial profile of the normalized electron thermal diffusivity between LHD and JT-60U could be explained by the role of $E_r$ on the collisional flow damping of the zonal flow which has a strong contribution towards determining the thermal diffusivity.

The temperature dependence of the thermal diffusivity $\alpha = (T_e/\chi_e)(d\chi_e/dT_e)$ is an important parameter for studying the plasma with an ITB. In the L-mode, the parameter $\alpha$ is positive and typically it is 1.5, which is predicted from the gyro-reduced
Bohm scaling and is also consistent with the power degradation of the global energy confinement in LHD of $\tau_E \propto (P/n)^{-0.6}$ [40]. If the parameter $\alpha$ stays positive, the formation of an ITB would never occur, because a spontaneous increase of electron temperature during the formation of an ITB requires a negative $\alpha$. Therefore $\alpha$ would be the most reasonable parameter to confirm whether the plasma is in the L-mode regime or the ITB regime, especially near the boundary between L-mode and ITB mode in space and in time, and therefore negative $\alpha$ can be a definition of an ITB plasma. The temperature dependence is investigated by power balance analysis and cold pulse experiments.

As shown in figure 7, the temperature dependence parameter of the electron thermal diffusivity in the core region becomes negative (-1.5) associated with the formation of an ITB, while the temperature dependence parameter outside the ITB region remains positive, $\alpha = 1.6$, which is consistent with the temperature dependence ($\alpha = 1.5$) of the gyro-reduced Bohm transport. Pulse propagation analysis of the cold pulse experiment also shows the negative $\alpha$ inside the ITB and a positive $\alpha$ outside the ITB in LHD as shown in figure 7(b). The change of sign of the temperature dependence of the

Figure 7. (a) Electron thermal diffusivity as a function of electron temperature inside the ITB region in LHD and radial profile of the temperature dependence derived from the cold pulse propagation experiments in (b)LHD and (c)JT-60U[22, 39].
Figure 8. (a) $\nabla T_e$ and (b) $T_e$ dependence of $\chi_e$ at different normalized radii in the JT-60U NBI plasmas and (c)(d) those in LHD NBI plasmas. The $\chi_e$ is normalized by the gyro-Bohm $T_e$ dependence and $\nabla T_e$ is normalized by $(R/T_e)$, where $R$ is the major radius and $L_T = \nabla T_e/T_e$ in (a) and (c).[39].

electron thermal diffusivity is a key to the formation of the ITB in LHD. In contrast, as seen in figure 7(c), the radial profile of the temperature dependence factor, $\alpha$, in a tokamak is quite different from that observed in LHD. The temperature dependence is zero inside the ITB and negative outside the ITB in JT-60U. This data imply that the temperature dependence alone cannot explain the formation of an ITB in a tokamak. In order to cause the spontaneous transition to the ITB plasma, the temperature gradient dependence should play a role inside the ITB in the JT-60U tokamak, because the transport reduction of $\alpha = 0$ is not enough to cause the transition as described above.

Figure 8 shows the $\nabla T_e$ and $T_e$ dependence of $\chi_e$ (obtained from the static analysis i.e. the power balance analysis) in JT-60U and LHD. In figure 8(a) and (c), the thermal diffusivity is normalized by $T_e^{3/2}$, which is predicted by Gyro-Bohm transport, in order to eliminate the $T_e$ dependence of $\chi_e$ and make the $\nabla T_e$ dependence clear. The $\nabla T_e$ is normalized by $R/T_e$. Because the ratio of major radius to scale length of temperature is expected to give the threshold of turbulence in the critical temperature gradient transport model.
As seen in figure 8(a)(c), there is no clear critical gradient scale length observed in LHD, although an existence of the critical gradient scale length is suggested in JT-60U. In contrast, a clear $T_e$ dependence of $\chi_e$ is observed in LHD, while there is no clear $T_e$ dependence of $\chi_e$ in JT-60U. The dependence of heat transport on $T_e$ and $\nabla T_e$ changes significantly depending on the scale length of the temperature in JT-60U. The thermal diffusivity $\chi_e$ shows a Gyro-Bohm type $T_e$ dependence when the temperature gradient is small enough. When the temperature gradient exceeds a critical value, it depends mostly on the scale length in a tokamak. This implies that there is a strong nonlinear mechanism in transport. One of the candidates for the mechanism causing strong non-linearity is non-linear flow damping of the zonal flow in a tokamak plasma. In the LHD heliotron, the thermal diffusivity $\chi_e$ shows a clear $T_e$ dependence and it suggests that the nonlinear mechanism in transport is weaker. The weak non-linearity might be due to the stronger collisional flow damping of the zonal flow than that in tokamaks. More study on driving and damping mechanisms of zonal flows is required to clarify the mechanism causing the differences in the $T_e$ and $\nabla T_e$ dependence of heat transport between heliotrons and tokamaks.

In LHD heliotrons, detailed experiment to investigate the $T_e$ and $\nabla T_e$ dependence of heat transport was performed using dynamic transport analysis of the discharges with repetitive pellet injection. After the pellet injection there are clear phases characterized by the sign of the time derivative of the $T_e$ gradients as indicated by phases I, II, III, IV. Both $T_e$ and the $T_e$ gradient increase in time as the electron density decreases in phases II and IV, which is a normal characteristic of heat transport in plasma. However, in phases I and III, the $T_e$ gradients decrease even with the increase of $T_e$. Figure 9(a) shows the normalized heat as a function of $T_e$ and the $T_e$ gradient in the weak $T_e$ dependence branch (phase II) and in the strong $T_e$ dependence branch (phase IV) after the pellets (12 pellets in series) are injected into the plasmas with different densities. The experimental data in phase II shows that the normalized heat flux gradually increases as the $T_e$ and $T_e$ gradient are increased, while that in phase IV shows a sharp increase of the normalized heat flux. It is noted that all the experimental data points of the 12 events are connected and located along curves (data points in phase II are on a weak $T_e$ dependence curve and data points in phase IV on a strong $T_e$ dependence curve). The data points in the transition phase (Phase III) are scattered between the weak and the strong $T_e$ dependence curves. When the plasma is on one of the branches, both $T_e$ and the $T_e$ gradient are uniquely determined for the given normalized heat flux, while they are not uniquely determined during the transition phase.

In order to investigate quantitatively the dependence of the normalized heat flux on the $T_e$ and $T_e$ gradient dependence, the normalized heat flux is given by $Q_e/n_e \propto T_e^\alpha(\nabla T_e)^\beta$, where $\alpha$ and $\beta$ are the $T_e$ and $T_e$ gradient dependence parameters, respectively. The difference between the $Q_e/n_e$ measured and $Q_e/n_e$ calculated with the parameters ($\alpha, \beta$) are investigated in a wide range of $Q_e/n_e$ for these two branches. Figure 9(b) and 9(c) show the contours of the square of differences between the normalized heat flux measured and that calculated, $\chi^2 = [Q_e/n_e - cT_e^\alpha(\nabla T_e)^\beta]^2$, where $c$ is a constant.
Figure 9. (a) Heat flux normalized by electron density as a function of electron temperature $T_e$ and temperature gradient in the weak $T_e$ dependence branch (phase II), in the transition period (phase III), and in the strong $T_e$ dependence branch (phase IV) at $\rho = 0.65$. The contours of the difference between the normalized heat flux measured and that calculated with the model of $cT_e^\alpha T_e^\beta$ with various values of $(\alpha, \beta)$ for (b) the weak $T_e$ dependence branch (phase II) and (c) the strong $T_e$ dependence branch (phase IV) at $\rho = 0.65$[41].

for various values of $(\alpha, \beta)$ in the weak and strong $T_e$ dependence branches. The $T_e$ dependence parameter $\alpha$ is 0.44 (-0.23, +0.25) for the weak $T_e$ dependence transport branch and 1.4 (-0.4, +0.3) for the strong $T_e$ dependence transport branch, while the $T_e$ gradient dependence parameter, $\beta$ is close to unity; 0.86 (-0.25, +0.25) and 1.1 (-0.2, +0.27) for the weak and the strong $T_e$ dependence branches, respectively. This strong $T_e$ dependence is consistent with the observation in tokamak when the $T_e$ gradient is
Figure 10. Temporal evolutions of the perturbed electron temperature with (a) the ∼ 420 μm diameter TESPEL injection, (b) the larger (∼ 670 μm diameter) TESPEL injection into the plasma with a line-averaged electron density of 0.5 × 10^{19} m^{-3} [46].

below the threshold [42]. The error bar is estimated from the range of α and β values for \( \chi^2 < 1.3 \chi^2_{\text{min}} \). In this experiment, the \( T_e \) increases as the \( T_e \) gradient is increased. The co-linearity between \( T_e \) and the \( T_e \) gradient causes large errors in the estimated α and β values. The sum of α and β is evaluated more accurately and it is 1.3 (-0.12, +0.12) for the weak \( T_e \) dependence branch and 2.5 (-0.24, +0.26) for the strong \( T_e \) dependence branch.

The strong \((T_e, \nabla T_e)\) dependence of \( \alpha + \beta = 2.3 \) observed in phase VI is consistent with the gyro-Bohm like dependence of \( \alpha + \beta = 2.5 \) in the collisionless regime. However the weak \((T_e, \nabla T_e)\) dependence of \( \alpha + \beta = 1.3 \) observed in phase II is much weaker than the gyro-Bohm dependence and it can be understood by taking account of the influence of a zonal flow. When zonal flows coexist with drift wave turbulence, the turbulent transport coefficient is strongly influenced by the damping rate of the zonal flow and deviates from the gyro-Bohm dependence[43]. The \((T_e, \nabla T_e)\) dependence \( \alpha + \beta \) depends on fluctuation levels and, when the collisional flow damping of a zonal flow is the dominant process, it can be 1.5, which is consistent with the \((T_e, \nabla T_e)\) dependence in phase II. These observation suggests that the collisional flow damping of a zonal flow in a non-axisymmetric system and is consistent with the fact that zonal flows are enhanced in the electron root plasma, where the collision flow damping becomes small.[24]

6. Non-local transport in non-axisymmetric system

The non-linearity of transport described above is due to the non-linearity of the driving and damping of turbulence and plasma flow. Since the scale lengths of the turbulence (\( \sim \))
poloidal gyro radius) and zonal flows (\textasciitilde tens of poloidal gyro radius) are different, the coupling between the turbulence and zonal flows is a candidate to cause the non-linearity of transport. Therefore the observation of nonlocal Te rise in a helical device will give new insight into the phenomenon of nonlocal Te rise, since helical systems have quite a different magnetic configuration (normally negative magnetic shear), which can affect the turbulence. The characteristics of the nonlocal Te rise invoked by the rapid edge cooling in helical plasmas have been investigated in LHD\cite{44, 45, 46}. In LHD, a variety of time responses of the nonlocal Te rise are obtained by changing the plasma parameters or the TESPEL size. Figure 10(a) shows a typical example of an abrupt increase in core Te just after the edge cooling. In this example, the \textasciitilde 420 \mu m diameter TESPEL is injected into the LHD plasma at a density of $0.5 \times 10^{19} \text{ m}^{-3}$ and as a consequence, the TESPEL reaches $\rho \sim 0.8$. As can be seen in figure 10(b), when the larger TESPEL (\textasciitilde 670 \mu m diameter) is injected into the LHD plasma, the time derivative of the core Te just after the edge cooling becomes much larger, compared with the case of figure 10(a). In this case, the line-averaged electron density and the TESPEL penetration depth are almost the same as those in the case of figure 10(a).

It is interesting that non-local transport appears both in LHD and tokamaks, although the $T_e$ and $\nabla T_e$ dependences of heat transport have significant differences between a heliotron and a tokamak. One of the candidate for the physics mechanism causing the non-local transport is energy transfer between high $k_r$ turbulence determining the transport and low $k_r$ zonal flows, because energy transfer between separated regions is required to explain propagation much faster than the diffusion of turbulence. It is important to investigate how the zonal flow changes during the non-local transport phenomena in order to understand turbulent transport.

7. Energy transfer between flows and turbulence and generation of parallel flow

Experimentally, direct measurements have been made of the radial-perpendicular component and radial-parallel component of the Reynolds stress in the boundary region in the TJ-II stellarator. As seen in figure 11(a), the perpendicular phase velocity of fluctuations (perpendicular to the magnetic field) is computed by means of the two point correlation technique using two floating potential signals measured by probes poloidally separated about 0.3 cm. The resulting radial profile of $v_\theta$ is radially flat for a mean plasma density below the threshold value ( $0.55 \times 10^{-19} \text{ m}^{-3}$ ), whereas above this critical density the perpendicular phase velocity reverses and the naturally occurring velocity shear layer appears in the proximity of the LCFS.

From the radial derivative of the mean parallel velocity and the radial-parallel component of the Reynolds stress, the radial-parallel contribution to the production of turbulent kinetic energy $P$ was experimentally evaluated in TJ-II as seen in figure 11(b). The production term combines the velocities cross correlation $<v_r M_\parallel>$ (momentum flux) with the mean velocity gradient ($\partial M_\parallel/\partial r$) and gives a measure of the amount of
energy per unit mass and unit time that is transferred locally between the mean flow and fluctuations. In low density plasmas, the production term $P$ as well as gradients in $\partial M_\parallel / \partial r$ are small and, in some cases, within experimental error bars. That means that, at low density, there is not any significant energy transfer between flows and turbulence. As the density approaches the threshold value to trigger $E \times B$ perpendicular flows the results show a different behavior. The production term clearly increases in the region where sheared flows are developed, reaching values up to $10^4 s^{-1}$. Furthermore, two different signs are found in $P$ thus implying that the turbulence can act as an energy sink ($P > 0$) for the mean flow or energy source ($P < 0$) near the shear layer [20]. This result suggests that parallel turbulent forces are relevant in momentum dynamics during the development of sheared flows in the proximity of the LCFS.

There are $E \times B$ flows and parallel flows driven by non-ambipolar flux in non-axi-symmetric systems. These flows become more significant in the plasmas where an internal barrier appears associated with the transition from ion root (small negative $E_\rho$) to electron root (large positive $E_\rho$). Figure 12(a) and (b) show the time evolution of the poloidal and toroidal flows in the NBI plasma with and without a 2nd ECH pulse. During the 2nd ECH pulse, the electron collisionality is low enough for the plasma to be in the electron root (positive electric field) as predicted by neoclassical theory. The positive radial electric field is mainly made up of the poloidal flow in the direction parallel to the $\langle E_\tau \times B_\phi \rangle$ drift. The large electric field of 10-20 kV/m and the large poloidal flow of 10-20 km/s are produced near the plasma center ($\rho = 0.3$) [figure 11(a)], where the sharp temperature gradient is observed. The neutral beam is injected to the plasma tangentially in the co-direction. Here the ”co” direction (positive flow velocity) is defined as parallel to the equivalent toroidal plasma current, which would produce the average transform actually produced by the external coil current. The ”counter” direction (negative flow velocity) is defined as anti-parallel to the equivalent
Figure 12. Time evolution of (a) central electron temperatures and densities, (b) poloidal flow velocity at $\rho = 0.38$[48], and (c) spontaneous toroidal flow velocity as a function of $E_r/B_\theta$ at $\rho = 0.4$ for tokamak and helical plasmas. The solid lines are best fits for the measured data and $\delta V_\phi = 0.15 \ (E_r/B_\theta) - 7.9$ for a Heliotron plasma and $\delta V_\phi = 1.34 \ (E_r/B_\theta) + 22.3$ for a tokamak plasma[49].

toroidal plasma current. As seen in figure 11(b), the plasma rotates parallel to the NBI when there is no 2nd ECH. This is simply because the toroidal momentum from the NBI causes the plasma toroidal flow. However, when the 2nd ECH is turned on, the toroidal flow velocity decreases and finally the plasma rotates anti-parallel to the NBI. These measurements clearly show the driven toroidal flow associated with a large poloidal flow (and large positive radial electric field) during the 2nd ECH pulse. The driven toroidal flow reaches $5 \times 10^4$ m/s at the plasma center and it is large enough to overcome the toroidal flow driven by a tangentially injected neutral beam. It should be emphasized that the direction of the toroidal flow is anti-parallel to the direction of the $<E_r \times B_\theta>$ drift. This is in contrast to the spontaneous flow in a tokamak, where the direction of the toroidal flow is parallel to the direction of the $<E_r \times B_\theta>$ drift, because the toroidal viscosity is nearly zero due to the toroidal symmetry.

The significant difference between tokamaks and a heliotron devices is the relationship between the poloidal field direction and the direction of dominant symmetry as seen in figure 12(c). The pitch angle of dominant symmetry is even larger than
that of the averaged poloidal field in a Heliotron device, while it is zero due to the toroidal symmetry in a tokamak. The formula for toroidal flow qualitatively explains the reversal of the sign of $v_{\phi}$ between CHS and tokamak plasmas, because $\alpha < \beta$ in Heliotron and $\alpha > \beta = 0$ in tokamak, where $\alpha$ is the pitch angle of the magnetic field and $\beta$ is the pitch angle of the minimum $\nabla B$ direction. The spontaneous toroidal flow observed in the plasma with an electron ITB in CHS is anti-parallel to the $< E_r \times B_\theta >$ drift direction. This is in contrast to the spontaneous toroidal flow in the direction parallel to the direction of the $< E_r \times B_\theta >$ drifts in tokamak plasmas, which is clearly demonstrate in figure 11(c). Here the spontaneous toroidal flow is given by the difference of magnitude of toroidal flow velocity between co-injection and counter-injection with a similar magnitude of momentum input and the radial electric field is given by the temperature gradient. In general, the toroidal flow in Heliotron plasmas is much smaller than that in a tokamak due to toroidal viscosity [33] which does not exist in tokamak plasmas. However, the spontaneous flow observed in CHS is comparable to that observed in the JFT-2M tokamak. This is because the radial electric field in CHS plasma with an internal transport barrier is much larger than that in the L-mode plasma in the JFT-2M tokamak. It should be noted that the ratio of spontaneous flow velocity to the radial electric field normalized by the poloidal field in the CHS Heliotron is 0.15, which is much smaller than that observed in a tokamak (1.34) by nearly an order of magnitude because of the parallel viscosity in the toroidal direction.

8. Ion internal transport barrier

Although the transition from L-mode to the high ion temperature mode plasma, which is characterized by the peaked ion temperature profile, has been achieved, the ion temperature gradient normalized $R/L_{Ti}$ was relatively small ($< 10$) compared with that observed in tokamaks. The high $T_i$ mode is characterized by a high central ion temperature and a low central ion thermal diffusivity, associated with a peaked electron density produced by NB fuelling with low wall recycling[50, 51].

More recently a clear ion internal transport barrier was observed associated with the formation of an electron internal transport barrier in the electron root[52]. As seen in figure13, a sharp temperature gradient of 12 keV/m appears at $\rho = 0.66$, which is much further outside of the electron transport barrier ($\rho < 0.3$). The ion temperature at the ion transport barrier is 0.3 - 0.4 keV and the normalized ion temperature gradient $R/L_{Ti}$ exceeds 30, which is comparable to the normalized electron gradient $R/L_{Te}$ observed inside the electron ITB[19]. The large separation of ITB locations between the ion transport barrier and the electron transport barrier is an interesting characteristic in non-axi-symmetric magnetic field systems, which is not observed in tokamaks. (The locations of electron ITBs and ion ITBs are close to each other[54]).

The high ion temperature plasma obtained by NBI heating (no additional ECH) in LHD has a moderate ion temperature gradient of $R/L_{Ti} = 10 - 15$ as seen in figure 13(b). In this discharge, there is no electron ITB observed ($T_e(0) < T_i(0)$). The radial
Figure 13. Radial profiles of (a) ion temperature gradient of an ion transport barrier plasma in CHS[53] and of (b) ion temperature of high $T_i$ discharges in LHD. In figure (a) the peak ion temperature gradient of 12keV/m gives the large normalized ion temperature gradient $R/L_{T_i}$ of 30 - 40.

profile of ion temperature observed in LHD is similar to that observed in the high $T_i$ mode without ECH heating[50, 51]. It should be noted that the change in $\nabla T_i$ appears at $R = 4.2m$, which suggest the formation of a weak internal transport barrier. Further study of ion transport for this high $T_i$ discharge should be done. One of the interesting features of this discharge is impurity exhaust. The large error bar in the ion temperature near the axis is due to the low intensity of the charge exchange line measured. The low intensity is mainly due to the hollow carbon density rather than the beam attenuation. It should be noted that the simultaneous achievement of good energy confinement and poor impurity confinement is a preferable transport characteristic in the transport of non-axisymmetric systems.

9. High density operation

Accessibility to high density operation is one of the important features of non-axisymmetric systems. In Wendelstein 7-AS, a high density regime with good energy confinement and low impurity concentration was achieved by building up the density rapidly with strong gas puffing at the very start of the discharge during NBI initiation[55, 56]. As discussed in the previous section, poor impurity confinement is achieved simultaneously with good energy confinement. In non-axisymmetric systems, simultaneous achievement of high density and poor impurity confinement has been also demonstrated experimentally. In LHD, a highly peaked density profile was created by injection of multiple pellets with a low recycling configuration[57, 58, 59]. The central densities achieved in Wendelstein 7-AS and LHD are $4-5 \times 10^{20} m^{-3}$ which is more than
Figure 14. Radial profiles of (a) electron density and (b) electron temperature in the normal confinement (NC), and attached and detached high density H-mode (HDH) in Wendelstein 7-AS (B = 2.5T)[56] and radial profiles of (c) electron density and (d) electron temperature in the super dense core plasma in LHD (B=2.64T)[58].

twice that of a tokamak with similar size and magnetic field strength. By replacing the plasma current in tokamak with the equivalent current in a helical system, the densities in helical systems exceed the Greenwald limit by a factor of 2.

Figure 14(a) and (b) show radial profiles of electron density and temperature in the normal confinement (NC), and the attached and detached high density H-mode (HDH) in Wendelstein 7-AS (B = 2.5T). Although the density profile of the normal confinement (NC) is peaked at the magnetic axis, the density profile of the high density H-mode (HDH) is broad. The significant difference between NC and HDH discharges is energy confinement and impurity confinement time. The energy confinement was up to twice that of standard scaling, while the impurity confinement time was reduced and approaching the energy confinement time by the reduction of the inward pinch velocity which normally exists in the NC discharge. A super dense core (SDC) plasma develops naturally in LHD as a peaked, high density profile is generated by multiple pellet injections from the outside mid-plane. As shown in figure 14(c)(d), A core region
with electron densities $\sim 4.6 \times 10^{20}\text{m}^{-3}$ and temperatures $\sim 0.85\text{ keV}$ is maintained by an internal diffusion barrier (IDB) located at normalized minor radius $0.3 < \rho < 0.5$. Unlike the tokamak internal transport barrier (ITB) discharges, there is no increase in the temperature gradient, possibly because the density becomes extremely large in LHD. The radial width of the IDB is $\sim 0.1\text{m} (\rho \sim 0.2)$. It should be pointed out that there is no impurity accumulation causing the termination of the discharge observed in LHD. The density is limited not by radiation collapse but by MHD instabilities triggered by the increase of pressure in the IDB region[60].

10. Summary

Transport barriers in confinement systems with non-axisymmetric magnetic fields are summarized focusing on the physics mechanisms determining the turbulence level in the plasma. There are differences in transport barrier characteristics between plasmas with axisymmetric magnetic fields (tokamak plasmas) and plasmas with non-axisymmetric magnetic fields (helical plasmas), although in both cases the turbulence in the plasma has a dominant contribution to the transport.

(1) The radial-perpendicular component of the Reynolds stress due to turbulence in the plasma produces a zonal flow with finite $k_r$ and zero $k_\theta k_\phi$ as well as a mean flow ($k_r = 0$). There are two types of zonal flows: one is a stationary zonal flow with the frequency below $1/\tan\text{damping}$ and the other is a zonal flow excited at the GAM frequency. The stationary zonal flow, which is considered to be more effective for the suppression of turbulence, is driven by turbulence and damped by collisional flow damping and non-linear flow damping.

(2) The internal transport barrier in a helical plasma extends to the plasma core where there is no strong $E_r$ shear, while that in a tokamak tends to be localized in the narrow region where the strong $E_r$ shear exists. The mechanism causing these differences is considered to be the $E_r$ effect on the reduction of collisional flow damping of the zonal flow. In a helical plasma, the radial electric field is expected to reduce the collisional flow damping, which is much larger than in a tokamak because of the ripple (if $E_r = 0$), enhance stationary zonal flows and hence reduce the turbulence level. On the other hand, $E_r$ does not contribute to the collisional flow damping and $E_r$ shear can be a dominant process to reduce the turbulence in a tokamak.

(3) In general, the heat flux normalized by density depends on the temperature, $T$, and the temperature gradient, $\nabla T$. Therefore the ratio of normalized heat flux to $\nabla T$ (thermal diffusivity) should have some dependence on $T$ and on the scale length of the temperature $T/\nabla T$. The thermal diffusivity in a helical plasma has a clear temperature dependence but has no clear scale length dependence. This is in contrast to tokamaks where there is a clear threshold of scale length of $T$ above which a significant enhancement of the thermal diffusivity exists and the $T$ dependence is often masked by the $\nabla T$ dependence. This difference can be also explained by the differences in the damping mechanism of zonal flows. The $T$ and $\nabla T$ dependence of heat flux measured
in a helical plasma is consistent with that expected by zonal flow theory assuming that collisional flow damping is the dominant term. The existence of a threshold of scale length in tokamak plasmas suggests that the non-linear flow damping of zonal flows is dominant.

(4) Non-local transport observed in helical plasmas is similar to that observed in tokamaks, although the $T$ and $T/\nabla T$ dependence of transport is quite different between helical plasma and tokamak plasmas. This fact implies that the physics mechanism to transfer the turbulence energy from one location to another location is not plasma current but the fluctuations with small $k_r$ (stationary zonal flow is one of the candidate). The coupling between turbulence(high-k)-turbulence(low-k) and turbulence-flow, that should cause the non-local transport is quite crucial for understanding transport both in tokamak and helical plasma.

(5) In helical plasmas the mean $E \times B$ and spontaneous parallel flow are driven by the radial-parallel component of the Reynolds stress due to turbulence and neoclassical non-ambipolar flux due to collisional processes. The viscosity tensor acts as damping and driving terms and the relation between $E \times B$ and spontaneous parallel flow is strongly affected by the magnitude of the viscosity in the poloidal and toroidal directions. There is a difference in the direction and magnitude of the parallel flow between helical and tokamak plasmas. For the given radial electric field, the spontaneous parallel flow is small and anti-parallel to the direction of the $<E_r \times B_\theta>$ drift in a helical plasma, while in a tokamak it is large and parallel to the direction of $<E_r \times B_\theta>$. The difference of the direction is due to the differences in symmetry, helical symmetry or toroidal symmetry.

(6) The simultaneous achievements of good energy confinement and poor impurity confinement and high density and poor impurity confinement are observed in a helical plasma. A peaked ion temperature with moderate gradient ($R/L_{T_i} = 10-15$) is achieved in NBI heated plasma by reducing the recycling. By applying ECH a strong ion temperature gradient ($R/L_{T_i} = 30-40$) is obtained associated with the formation of an electron ITB in the electron root. The ion ITB and electron ITB are located in different regions of the helical plasma, which is in contrast to that the proximity of the electron ITB and ion ITB in a tokamak plasma. The density achieved in a helical system is twice of that in a tokamak with similar size and magnetic field (equivalent Greenwald factor of 2). Even with quite high density operation, the impurity is not a problem affecting the plasma performance.

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