Observation of energetic-ion losses induced by various MHD instabilities in the Large Helical Device (LHD)

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Abstract
Energetic-ion losses induced by toroidicity-induced Alfvén eigenmodes (TAEs) and resistive interchange modes (RICs) were observed in neutral-beam heated plasmas of the Large Helical Device (LHD) at a relatively low toroidal magnetic field level (\(B_t \leq 0.75\) T). The energy and pitch angle of the lost ions are detected using a scintillator-based lost-fast ion probe. Each instability increases the lost ions having a certain energy/pitch angle. TAE bursts preferentially induce energetic beam ions in co-passing orbits having energy from the injection energy \(E = 190\) keV down to \(130\) keV, while RICs expel energetic ions of \(E = 190\) keV down to \(\sim 130\) keV in passing–toroidally trapped boundary orbits. Loss fluxes induced by these instabilities increase with different dependences on the magnetic fluctuation amplitude: nonlinear and linear dependences for TAEs and RICs, respectively.

(Some figures in this article are in colour only in the electronic version)

1. Introduction
A self-sustained D–T burning plasma can be realized under the condition that fusion-born energetic alpha particles (\(\alpha\)s) are confined long enough to heat the bulk plasma [1]. In addition to effective alpha heating of the bulk plasma, loss of \(\alpha\)s should be suppressed down to an acceptable level to prevent plasma-facing components from the localized damage due to the impact of escaping \(\alpha\)s. A better understanding of transport and loss process of these energetic ions is therefore essentially required to realize a nuclear fusion reactor. The principal concern is that D–T produced \(\alpha\)s and super-Alfvénic ions produced by neutral beam (NB) injection or ion cyclotron resonance heating may destabilize energetic-ion-driven magnetohydrodynamic (MHD) instabilities such as Alfvén eigenmodes (AEs) [2] and energetic-particle-continuum modes (EPMs) [3] because those instabilities can potentially cause radial transport and/or losses of energetic ions through wave–particle resonance processes. The effect of MHD instabilities driven by the bulk plasma such as sawtooth and tearing mode (TM) on energetic-ion transport is also of great concern since transport of energetic ion may be affected through not only wave–particle resonance but also stochastization of energetic-ion orbits [4].

Energetic-ion losses induced by AEs have been so far studied experimentally in several tokamaks [5–7]. Previously, detailed loss processes of energetic ions due to AEs were studied in TFTR and JT60-U [7–9]. In TFTR, a scintillator-based lost fast ion probe (SLIP) was employed to simultaneously measure the energy \(E\) and pitch angle \(\chi\) of escaping fusion products and beam ions [10, 11]. In those experiments, decreases in neutron emission rate and, at the same time, increase in energetic-ion losses due to global Alfvén eigenmodes (GAEs) were observed [7]. Nowadays, efforts on numerical simulations are being made to reveal the energetic-ion transport due to energetic-ion-driven MHD instabilities [12, 13]. In addition, MHD instabilities driven by the bulk plasma can potentially induce redistribution and/or losses of energetic ions. It has been experimentally observed that significant redistribution of energetic ions is caused by sawteeth in tokamaks [14–16]. Also, it has been recognized that a TM leads to redistribution and/or loss of energetic...
ions. A reduction in the NB current drive efficiency caused by beam-ion losses due to TMs was observed in DIII-D [17] and also energetic-ion loss by neoclassical TM has recently been discussed in ASDEX-U [18]. In both cases, beam-ion losses are interpreted to be due to stochasticization of beam-ion orbits. At present, energetic-ion losses due not only to energetic-ion-driven MHD modes but also pressure/current-driven MHD modes are intensively studied using the SLIP in NSTX, ASDEX-U and JET [19–21].

It is also of importance for helical/stellarator plasmas to understand anomalous radial transport and/or loss of energetic ions caused by various MHD instabilities. In experiments prior to this work, anomalous beam-ion losses induced by toroidicity-induced Alfvén eigenmodes (TAEs) and EPMs were measured with the SLIP on medium-scale stellarators/helical devices, i.e. CHS [22, 23] and W7-AS [24]. In this paper, we present characteristics of beam-ion losses due to both MHD instabilities, i.e. energetic-ion-driven TAEs and pressure-driven resistive-interchange (RIC) modes [25] in high $\beta$ discharges with $t_b > v_A$ beam ($t_b$ and $v_A$ indicate the velocity of the beam and Alfvén speed) of a large-scale helical device where energetic ions are well confined, compared with these medium-sized devices. Beam-ion losses are studied using the SLIP, which has recently been installed on the outboard side of the Large Helical Device (LHD) to investigate behaviours of co-going beam ions [26–28].

This paper is organized as follows. In section 2, a brief descriptive introduction of the LHD and the SLIP is presented. Characteristics of TAE-induced beam-ion loss, i.e. energy and pitch angles of escaping beam ions due to TAE burst, and high-frequency fluctuations seen on energetic-ion loss signals are presented in section 3. Anomalous losses of energetic ions induced by the RIC mode are also reported in section 3. Section 4 gives a summary. Finally, the equations to calculate Hz fluctuations induced by TAE mode are given in appendix A. The dispersion relation of RIC is also given in appendix B.

2. Experimental setups

LHD is a large helical device having number of toroidal periods $M$ of 10 and multipolarity $l$ of 2. The major radius $R_0$ and the plasma minor radius $a$ are 3.9 m and $\sim$0.6 m, respectively. The LHD is equipped with three negative-ion-source based NB injectors of which injection energy $E_b$ is up to 190 keV. One of the three tangentially injects NBs in the opposite direction of the equivalent plasma current whereas the others tangentially inject NBs in the direction of the equivalent plasma current in counter clockwise $B_t$ which is defined from the top view of the torus. The injected beam ions are super-Alfvénic in many cases and often destabilize TAEs [29, 30].

The SLIP important in this work has been lately developed and installed on the outboard side of the LHD torus. It is actually a magnetic spectrometer using a set of apertures and a scintillator plate (ZnS : Ag), providing energy $E$ and pitch angles $\chi = \arccos(v_{//}/v)$ of lost energetic ions simultaneously as a function of time, where $v_{//}$ is the velocity of ion parallel to the magnetic field and $v$ is the velocity of ion. Detailed description of the design concept and hardware of SLIP on the LHD are available in [26]. Figure 1(a) shows the installation...
position of SLIP and the typical orbit of a co-going energetic ion reaching the SLIP. Figure 1(b) shows the relation between the pitch angle at the plasma surface and at the SLIP position. As seen from figure 1(b), the differences of \( \chi \) between the SLIP position and that at the plasma surface from where the energetic ion is lost to the SLIP are in the range of \( \pm 10^\circ \).

A luminous image produced on the scintillator screen is transferred to the outside of the vacuum vessel and subsequently is divided by a half mirror into two paths. Each scintillation light image is monitored with a \( 4 \times 4 \)-photomultiplier tube (PMT) array and a CMOS camera, simultaneously. The relative sensitivity of PMTs is calibrated with an electro-luminescence sheet emitting a blue-green light uniformly within \( 10\% \) error. The detection ranges of the detectable light uniformly within \( 10\% \) error. The detection ranges of the detectable

In this paper, the discharge at \( t \sim 2.82 \) s is mainly discussed. The radial profiles of electron temperature \( T_e \) and \( n_e \) measured with Thomson scattering diagnostic (TSD) at this time are shown in figure 4(a). Hydrogen atom density \( n_0 \) is assumed to be scaled as \( 1/n_e^2 \), which was derived from the experimental result on the CHS having a similar magnetic configuration [39]. As seen from figure 4(a), \( n_0 \) does not change significantly towards the last closed flux surface \( (r/a = 1) \) and would rapidly decrease from the value at the edge of the ergodic layer where electron density and temperature are very low. The deposition profile of NB-generated energetic ions is shown in figure 4(b). It is calculated using the FIT code [40], which gives the deposition profile of beam ions averaged over the magnetic surface taking into account the prompt orbit loss. The MHD equilibrium is reconstructed using experimental data by the VMEC2000 code [41] under the free boundary condition. The reconstruction of MHD equilibrium was done using the pressure profile evaluated from the experimentally obtained data on the assumption of \( T_e = T_i \) and \( n_e = n_i \) having a multiplication factor of energetic-ion effect (in this case, the factor is \( \sim 1.2 \)) so that the calculated magnetic axis position agrees well with that predicted from the \( T_e \) profile.

3. Experimental results

3.1. Typical discharges with toroidal Alfvén eigenmodes and RICs excited in NB-heated high beta plasmas of the LHD

Figure 3 shows a typical discharge with TAEs obtained at \( B_t = 0.6 \) T (CCW direction) in the magnetic configuration of \( R_{ax} = 3.6 \) m. In this shot, three tangential NBs are tangentially injected. The electron temperature at the centre \( T_0 \) is \( \sim 0.8 \) keV, and the line-averaged electron density \( (n_e) \) is \( \sim 1.5 \times 10^{19} \) m\(^{-3} \). The volume-averaged beta value \( (\beta) \) is relatively high, \( \sim 2.0\% \), as seen from figure 3(a). The ratio of gyroradius of beam ions to averaged plasma minor radius \( \rho_b/(a) \) is fairly large, i.e. \( \sim 0.15 \) m/0.6 m. The ratio of the initial energetic-ion velocity to the Alfvén speed is \( v_b/(A_v) \sim 2 \). The volume-averaged energetic-ion beta value \( \langle \beta \rangle \) is comparable to \( \langle \beta \rangle \), estimated to be \( \sim 1\% \), on the assumption of classical slowing down of beam ions for \( \sim 5 \) MW beam deposition. Strongly excited TAE and RICs are identified in the quasi-stationary phase \( (2.8 \leq t \lesssim 3.2) \) s from the spectrogram analysis of MP signal. In high \( \beta \) plasmas in the inward-shifted configuration \( (R_{ax} = 3.6 \) m), TAE and low-frequency RICs often coexist. RICs are excited near the plasma edge due to the steep gradient of the bulk plasma pressure and can generate a magnetic island in the nonlinear evolution regime, while a linearly unstable interchange has no magnetic island [35]. The radial structure of the RICs can be derived from the data of SX/AXUV detector arrays [36–38].

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Figure 3. (a) Time evolutions of NBI absorbed powers $P_{\text{abs}}$, the bulk plasma beta $\langle \beta \rangle$, the line-averaged electron density $\langle n_e \rangle$ and spectrogram of MP signal, where TAE ($m \sim 1, n = 1$) and RIC fluctuations (frequency $< 10 \text{kHz}$) are identified. (b) Power spectra of magnetic fluctuation. The dominant peak at 2 kHz frequency corresponds to RIC ($m/n = 1/1$). The smaller peak at $f \sim 4 \text{kHz}$ is the second harmonic of the RIC.

Figure 4. (a) Profiles of electron temperature, electron density and neutral density at $t = 2.82 \text{s}$. Neutral density is inferred from the scaling of $1/n_e^2$. (b) Radial profiles of deposited beam ions and the rotational transform $\iota/2\pi$ (inverse of safety factor) at $t = 2.82 \text{s}$, where the beam deposition profile is estimated using the FIT code, and $\iota/2\pi$ is calculated using the VMEC2000 code.

The scintillator image in the time frame of $t = 2.820$–$2.822 \text{s}$ captured by a CMOS camera is shown in figure $5(a)$. On this image, three primary-loss domains, named $D_1$, $D_2$ and $D_3$, are clearly recognized. Lost energetic ions of $[\rho_h(E), \chi] = [11–15 \text{ cm} (80–150 \text{ keV}), 25–35^{\circ}], [8–18 \text{ cm} (50–190 \text{ keV}), 35–45^{\circ}]$ and $[8–18 \text{ cm} (50–190 \text{ keV}), 45–60^{\circ}]$ would reach domains $D_1$, $D_2$ and $D_3$, respectively. Here, $\rho_h(E)$ is evaluated using $B_{\text{SLIP}} = 0.37 \text{T}$. The scintillation image expands into the region of $\rho > 17 \text{ cm}$ ($E > 190 \text{ keV}$), especially in the $D_3$ domain, although $E_b$ is less than 190 keV. This is due to the relatively low energy resolution of the SLIP, typically $\pm 20 \text{ keV}$ for $E/\chi = 150 \text{ keV}/60^{\circ}$, even larger for higher energy ions. The orbit calculation indicated that energetic ions have passing, passing and passing–toroidally trapped (or so-called ‘transition’) orbits would reach $D_1$, $D_2$ and $D_3$ domains, respectively. On the other hand, figure $5(b)$ shows these loss fluxes monitored with 16 PMTs whose frequency responses are 20 kHz for Nos 1–6, 8, 11, 12, 14, 15 and 16 PMTs and 200 kHz for Nos 7, 9, 10 and 13 PMTs. As seen from figures $5(a)$ and $(b)$, TAE/RIC-induced losses are dominantly detected on the above-mentioned three domains $D_1$, $D_2$ and $D_3$. The RIC-induced loss and TAE-induced one are detected mainly on $D_1$ and $D_2$, respectively. The loss flux detected on $D_3$ is caused mainly by Coulomb collisions in the non-uniform magnetic field structure, since the $D_3$ region is observed even in the phases without TAE and RICs. In subsequent subsections, the effects of respective MHD instability, TAE or RIC on beam-ion losses will be discussed in detail.

3.2. Beam-ion losses due to TAE

A strongly excited TAE having 60–80 kHz is observed in the time window from $t \sim 2.7$ to 3.3 s by means of MPs, as shown in figure 3. The toroidal and poloidal mode numbers
Figure 5. (a) Scintillation pattern captured with a CMOS camera. Three dominant loss domains (D1, D2, and D3) are observed. The spot in domain D3 appears even if TAE/RIC is not excited. (b) Time traces of magnetic fluctuations by TAE and RIC, and 16 SLIP signals. TAE-induced loss is clearly observed in $\Gamma_{\text{SLIP,9}}$, $\Gamma_{\text{SLIP,10}}$, and $\Gamma_{\text{SLIP,11}}$, details of which are discussed in section 3.2. RIC-induced loss is clearly observed as frequent spikes in $\Gamma_{\text{SLIP,5}}$, $\Gamma_{\text{SLIP,6}}$, $\Gamma_{\text{SLIP,13}}$, and $\Gamma_{\text{SLIP,14}}$, details of which are discussed in section 3.3.

are $n = 1$ and $m \sim 1$, respectively. The values of $n$ and $m$ were determined by MPs. The amplitude of TAE evaluated at the MP position is $\sim 5 \times 10^{-5}$ T at $t = 2.82$ s. Shear Alfvén continua calculated by the STELLGAP code [42] are shown in figure 6(a). A TAE eigenfunction is calculated using the AE3D code [43] for $m = 0$ to 10 assuming that ion density $n_i = n_e$ and MHD equilibrium is reconstructed using VMEC2000 (figure 4). As seen from figure 6(b), the eigenfunction of the observed TAE is composed of two dominant Fourier components and has a similar character to a core-localized TAE which can exist in the low magnetic shear region of a tokamak plasma. This TAE is the odd parity mode, and has the peak of the eigenfunction at the normalized radial position $r/a \sim 0.6$.

Unlike tokamaks [44], a TAE with odd parity is often strongly excited having a large amplitude in the LHD. Although the reason is not yet clear at this moment, it may be attributed to the strong continuum damping at the very edge of the plasma having a high magnetic shear for an even parity TAE, as seen from the character of the TAE gap near the edge shown in figure 6(a). In high-$n$ and analytic approximation, the continuum damping rate increases strongly with the increase in the magnetic shear strength [45]. Moreover, the continuum damping rate for low-$n$ TAE increases more rapidly with the increase in the magnetic shear strength, compared with that for high-$n$ TAE [46–48]. This may be a dominant cause that even TAE is usually suppressed or only very weakly destabilized. A detailed numerical calculation is required to draw a decisive conclusion on the difference in excitation of even and odd parity TAEs, and is left for a future work.

The fluctuations of $n_e$ due to TAE are observed as $H\alpha$ fluctuations by FHA diagnostics in LHD plasmas at fairly low $B_t$ (typically $< 1$ T). Information on the radial structure of TAE can be derived through the cross correlation between the $H\alpha$ fluctuation $I_{\text{H\alpha}}$ signal and the MP signal $d\theta/dt$ [33]. The fluctuation components modulated by TAE in $I_{\text{H\alpha}}$ were
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Figure 6. (a) Shear Alfvén spectra of $n = 1$ calculated by the STELLGAP code, where $m = 0$–10 Fourier modes are included. In the calculation, profiles of $n_e$ and $\lambda/2\pi$ shown in figure 4 are used. The calculated and observed TAE frequencies are shown with thin and broken horizontal lines, respectively. (b) Eigenfunction of the TAE mode calculated with AE3D. This mode has odd radial parity.

Figure 7. (a) Profiles of the local $H\alpha$ ($H\alpha$ emissivity) fluctuation induced by $m = 1$ and $m = 2$ Fourier components of TAE eigenfunction calculated by AE3D where $n_0 \propto n^{-2}_e$ is assumed. (b) Profile of local $H\alpha$ fluctuation induced by TAE in the cross section of the LHD plasma viewed by the FHA diagnostic, where the mode pattern is uniform in the toroidal direction. The eigenfunction of TAE is converted from that in the flux coordinate used in (a) to the real one based on VMEC2000 calculation. (c) The $H\alpha$ fluctuation profile obtained experimentally with FHA (open circles) and that calculated from the result shown in (b) that was calculated with the AE3D code (open squares). The mode pattern shown in (b) is rotated 10 times having 1000 steps for each rotation.

calculated by the product of the coherence $\gamma^2_{H\alpha-MP}$ and the power spectral density of $\tilde{I}_{H\alpha}$, where the coherence between TAE-induced fluctuations and background fluctuations due to edge turbulence is assumed to be negligibly small [49]. Thus, the derived TAE-induced $H\alpha$ fluctuations integrated along the line of sight are shown in figure 7(c). On the other hand, the predicted $H\alpha$ fluctuations were calculated by TAE eigenfunctions obtained by the AE3D code (figure 6) and plasma parameters including assumed neutral density profile shown in figure 4. The predicted electron density fluctuations were obtained based on the ideal MHD theory. The plasma displacement due to shear Alfvén waves such as TAE was evaluated under the condition that the magnetic perturbations are expected to be purely transverse. Each local $H\alpha$ emission pattern ($H\alpha$ emissibility) induced by $m = 1$ and $m = 2$

Fourier components in TAE eigenfunction in circular plasma is shown in figure 7(a). The local $H\alpha$ emission induced by the whole TAE eigenfunction is given in figure 7(b), on the assumption that the mode pattern of local $H\alpha$ emission is uniform in the toroidal direction. The $Z$-profile of $H\alpha$ emission integrated along the lines of sight shown in figure 1(c) in the $Z$-direction is shown as the profile of the normalized minor radius (figure 7(c)). Detailed equations used for this calculation are given in appendix A. Thus, the calculated profile agrees well with the structure obtained using FHA. This result indicates that this FHA diagnostics, newly introduced on the LHD, has the potentiality to give useful information on the internal structure of TAEs excited in high beta plasmas characterized by relatively high densities at considerably low toroidal fields.
As shown in figure 8(a), a sharp drop and a relatively slow rise in the charge-exchanged fast neutral fluxes \( \Gamma_{\text{neutral}} \) in the range of energy from 180 to 190 keV are observed, correlated with each TAE burst, while \( \Gamma_{\text{neutral}} \) in the range \( E = 130 – 160 \) keV exhibits a rapid rise and slow decay. The temporal evolution of charge-exchanged neutral fluxes will indicate the effect of clump–hole pair creation in the velocity space caused by TAE burst [50]. The time evolution of \( \Gamma_{\text{neutral}} (E = 130 – 160 \) keV) related to the clump formation indicates the rapid radial transport of energetic ions, suffering from rapid changes in ion energy and pitch angle, and would be lost outside the plasma during the process. The loss flux of energetic ions detected on the \( D_2 \) domain of the SLIP has a similar shape of the pulse. The characteristic decay time of the beam-ion loss flux to SLIP, \( \Gamma_{\text{SLIP,10}} \), is evaluated to be \( \sim 4.5 \) ms, which is comparable to the estimated slowing down time through Coulomb collision of beam ions with electrons (\( \sim 5 \) ms) at the peak of the TAE eigenfuction without TAE. During the slowing down of the formed clump of energetic ions between TAE bursts, certain number of ions will be lost and detected by the SLIP. This process would lead to the slow decay of the SLIP signal \( \Gamma_{\text{SLIP,10}} \) shown in figure 8(a). These observations tell us that a substantial transport/loss of energetic ions is induced by TAEs. That is, the peak just after the TAE burst is thought to be caused by the TAE burst. It should be noted that the signal \( \Gamma_{\text{SLIP,10}} \) is also modulated more frequently, compared with the change by the TAE burst. This frequent modulation is caused by the RIC mode excited near the plasma edge [25], and will be discussed in more detail in the next subsection. After the frequent modulations due to RIC modes in \( \Gamma_{\text{SLIP,10}} \) were removed using a numerical band-pass filter, the increment of the beam-ion loss rate due to TAE bursts, \( \Delta \Gamma_{\text{SLIP,10}} \), was evaluated. Moreover, time averaging, at every \( 10 \) \( \mu \)s time interval, was also applied to remove the remaining white noise in \( \Gamma_{\text{SLIP,10}} \). Thus, the derived smoothed \( \Gamma_{\text{SLIP,10}} \) is also overlaid in the trace of \( \Gamma_{\text{SLIP,10}} \). In figure 8(b), \( \Delta \Gamma_{\text{SLIP,10}} \) evaluated with the difference of \( \Gamma_{\text{SLIP,10}} \) between the pulse peak and that just before each TAE burst is plotted versus the magnetic fluctuation amplitude measured at the MP position. The error bar corresponds to the level of white noise. In this figure, the curve of \( \Delta \Gamma_{\text{SLIP,10}} \propto (b_{\text{TAE}}) \) is also shown. As seen from figure 8(b), \( \Delta \Gamma_{\text{SLIP,10}} \) induced by TAE increases non-linearly or approximately quadratically with the increase in the TAE magnetic fluctuation amplitude, although the data points are scattered appreciably. This may suggest a diffusive-type loss rather than a convective one [51]. These energetic ions may not be lost promptly from the plasma region by TAE. Presumably, first, energetic ions would be transported by TAE excited around \( r/a \sim 0.6 \) to the plasma edge region, as discussed in [50, 52], being scattered into the loss orbit.

The temporal changes in \( E \) and \( \chi \) associated with TAE burst may be obtained through comparison among the loss flux signals \( \Gamma_{\text{SLIP,1}}, \Gamma_{\text{SLIP,2}}, \ldots, \Gamma_{\text{SLIP,16}} \) having various \( E \) and \( \chi \). The temporal evolutions of \( \Gamma_{\text{SLIP,14}}, \Gamma_{\text{SLIP,13}}, \Gamma_{\text{SLIP,10}} \) and \( \Gamma_{\text{SLIP,6}} \) are shown together with NPA signals in figure 9(a). The signals \( \Gamma_{\text{SLIP,10}} \) and \( \Gamma_{\text{SLIP,6}} \) and the charge-exchanged neutral having \( E = 130 \) keV \( \Gamma_{\text{neutral}}(E=130\text{keV}) \) are rapidly increased and slowly decayed by the TAE burst exhibiting similar wave forms, although \( \Gamma_{\text{SLIP,6},1} \) is appreciably deformed by RICs. In contrast to them, the signals \( \Gamma_{\text{SLIP,13}} \) and \( \Gamma_{\text{SLIP,14}} \) and the charge-exchanged fast neutral having \( E > 180 \) keV \( \Gamma_{\text{neutral}}(E=180\text{keV}) \) are suddenly decreased and slowly recovered by the TAE burst having similar wave forms. That is, the hole formation in the velocity space by the TAE burst is seen as a sharp drop on the wave of \( \Gamma_{\text{neutral}}(E=180\text{keV}) \) which will be proportional to the source signal of energetic ions supplied by NBI. Energetic ions in this energy range will also be lost by the TAE burst suffering from a certain change in the pitch angle, which will be detected as \( \Gamma_{\text{SLIP,14}} \) and \( \Gamma_{\text{SLIP,13}} \). At the same time with the hole formation, the formed clump with a certain energy drop is recognized as a pulse increased in the signal \( \Gamma_{\text{neutral}}(E=130\text{keV}) \). Thus, the generated energetic ions with some energy drop will be lost and will be detected as \( \Gamma_{\text{SLIP,6}} \) and \( \Gamma_{\text{SLIP,10}} \). The above-mentioned characteristics of these signals indicate clearly that the \( E \) and \( \chi \) of lost ions are quickly changed during each TAE burst. This is also clearly seen on the scintillation pattern shown in figure 9(b). If the fast ions are kicked by TAE into another position, this may lead to another loss orbit. The increase in \( \Gamma_{\text{SLIP,6}} \) by the TAE burst shows a certain time lag for that in \( \Gamma_{\text{SLIP,10}} \), although the
signal seems to be still deformed by RIC-induced losses. This time lag suggests that further pitch-angle scattering by TAE bursts needs a certain time.

Interestingly, high-frequency fluctuations of $\Gamma_{\text{SLIP}}$ correlated with the TAE mode frequency ($\sim 70$ kHz) were observed using a correlation analysis between $\Gamma_{\text{SLIP}}$ and the MP signal, as seen from the spectrogram of coherence shown in figure 10(a). High-frequency ion loss fluxes correlated with TAE are often observed in the ranges of energies and pitch angles of the lost ions: $E/\chi \sim 180$ keV/45–55° and 6–23 keV/25–35°. The $E$ and $\chi$ ranges match with the regions indicated by shaded circles in figure 10(b). The energy of the lost energetic ions in regions 5 and 15 shown in figure 10(b) satisfies the fundamental and the side-band resonance conditions of TAE, respectively, that is $v_i \geq v_A$ and $>v_A/3$ at the TAE mode peak [53]. The Alfvén speed at the peak position of the TAE eigenfunction ($r/a \sim 0.6$) is $2.2 \times 10^4$ ms$^{-1}$. $v_i$ of the lost beam ions in region 5 is $3.4 \times 10^2$ ms$^{-1}$ ($\rho_b \sim 15$ cm, $\chi \sim 50^\circ$ and $B_{\text{SLIP}}$ at SLIP position is 0.37 T) corresponds to $v_i/v_A \sim 1.5$. The velocity in region 15 is $9.2 \times 10^2$ m s$^{-1}$ ($\rho_b \sim 5$ cm, $\chi \sim 30^\circ$) corresponds to $v_i/v_A \sim 0.4$. The fluctuations in $\Gamma_{\text{SLIP}}$ suggest the existence of fundamental and side-band resonant interactions between TAE and energetic ions.

3.3. Energetic-ion losses due to RIC

As mentioned briefly in the previous subsection, these low frequency instabilities also induce energetic-ion losses, which are clearly detected by the SLIP. In this high beta shot, RIC with $m/n = 1/1$ is strongly excited. The mode frequency of about 2 kHz is related to mainly poloidal plasma rotation (figure 3(b)). The fluctuations in AXUV signals in the frequency range 1.5–2.5 kHz would give the information of the eigenfunction of the $m/n = 1/1$ RIC mode. The radial profiles of the AXUV fluctuation amplitude related to the $m/n = 1/1$ RIC mode in the time window $t = 2.75–2.83$ s are shown in figure 11(a), together with the fitted AXUV fluctuation data obtained using a model profile of the local emissivity of AXUV. In this analysis, the parameters $x_0$ and $w$ in the model function exp$\left\{[(r/a - x_0)/w]^{2}\right\}$ were adjusted so that the calculated AXUV fluctuation profile taking into account the lines of sight of the AXUV detector array (figure 1(d)) should best fit with the experimental data [37]. Here, $x_0$ and $w$ correspond to the location of the rational surface and the island width. The model function of the local emissivity corresponds to the radial structure of the $m/n = 1/1$ RIC mode. The best fit result was obtained with $x_0 = 0.9$ and $w = 0.2$. Figure 11(b) shows the thus obtained radial profile of the emissivity fluctuation caused by the $m/n = 1/1$ RIC mode. The peak position $x_0$ agrees well with the position of the $\iota/2\pi = 1$ rational surface shown in figure 4(b). This approximately corresponds to the eigenfunction of the RIC mode. On the other hand, the dispersion relation of RIC is given in appendix B and the eigenmode function can be derived from the eigenmode equation. In this paper, however, we did not make direct comparison of the RIC-induced fluctuation profile derived from experimental data (figure 11(b)) with the eigenfunction calculated from the RIC eigenmode equation, because the experimentally obtained eigenfunction is applicable to a qualitative discussion on the interaction of the RIC mode with energetic ions.

Energetic-ion losses induced by RIC are observed mainly in domains $D_1$ and $D_3$ on the $\rho_b-\chi$ plane shown in figure 5(a). The energetic-ion loss fluxes are increased and modulated

![Image](image.png)

Figure 9. (a) Time traces of TAE magnetic fluctuations, $\Gamma_{\text{neutral}}$ and $\Gamma_{\text{SLIP}}$, in which RIC-induced modulations are removed. (b) Scintillation pattern obtained by a CMOS camera in the phase $t = 2.820–2.822$ s. Directions of the pitch angle scattering and the energy drop of energetic ions by TAE bursts are indicated with arrows.
Figure 10. (a) Spectrogram of coherence $\gamma$ between $\Gamma_{\text{SLIP},5}$ and MP signal. Losses correlated with TAE oscillations are observed with coherence analysis. (b) Regions where high correlation between MP signal and $\Gamma_{\text{SLIP}}$ are observed (shaded circles). In particular, regions 15 and 5 have higher coherence.

Figure 11. (a) Radial profiles of the amplitude of RIC fluctuations measured with AXUV array (open circles) and from the calculation (open squares). Note that the signal $dI_{\text{AXUV}}$ obtained by the AXUV array corresponds to the signal integrated along the lines of sight shown in figure 1(d). The parameters $x_0$ and $w$ in the model profile of the local emissivity $\exp\left(\frac{(r/a - x_0)^2}{w^2}\right)$ were adjusted so that both profiles should have good agreement, where $x_0$ and $w$ indicate the location of the rational surface and the island width. Here, $x_0 = 0.9$ and $w = 0.2$ give good agreement as shown in this figure. (b) Radial profile of local fluctuation of $m/n = 1/1$ RIC derived from the AXUV fluctuation data in (a) through the above model fitting. $x_0$ agrees well with the predicted location of the $\iota/2\pi = 1$ rational surface.

synchronizing with the RIC mode oscillation as seen from figure 5(b). The zoomed time evolutions of the magnetic fluctuation of RIC and the beam-ion loss flux signals $\Gamma_{\text{SLIP},6}$ and $\Gamma_{\text{SLIP},13}$ are shown in figure 12(a). The pitch angle $\chi$ of these lost beam ions is close to that of the so-called transition particles whose orbit randomly changes between co-passing and toroidally trapped orbits [54]. It should be noted that the particle gyro radius $\rho_h$ corresponds to the ions with a slightly reduced energy from the injection energy of NB. The fluctuations of $\Gamma_{\text{SLIP}}$ correlated with RIC were observed using a coherence analysis (figure 12(b)). It should be noted that the $\Gamma_{\text{SLIP}}$ fluctuation is not only correlated with the fundamental RIC mode of $\sim 2\,\text{kHz}$ but also the second harmonic having $\sim 4\,\text{kHz}$ (figures 12(b) and 3(b)). Moreover, it is clearly found from the phase analysis that the phase difference between $\Gamma_{\text{SLIP},13}$ and RIC is approximately fixed at $\sim 185^\circ$ having an appreciable variation of about $\pm 50^\circ$ (figure 12(c)). The reason why the phase does not stay rigorously constant is thought to be that RIC induced losses would be induced through stochasticization of energetic-ion orbits due to large amplitude RIC oscillations. The increment of loss fluxes due to RICs increases linearly rather than non-linearly with the magnetic fluctuation amplitude, as shown in figure 12(d). This dependence suggests a convective-type loss [51].

In the edge region of the LHD, the plasma rotates mainly in the poloidal direction due to the high toroidal viscosity.
The $m/n = 1/1$ RIC mode frequency of $\sim 2$ kHz will reflect the poloidal plasma rotation. In this situation, we discuss whether or not trapped energetic ions could interact with this mode. The resonance condition for helically trapped ion is $\omega \sim N \Omega$, where $\omega$ and $\Omega$ indicate the frequency of the mode and the toroidal precession frequency, and $N$ is the harmonics number [4]. Note that in the LHD configuration, once energetic ions are deeply trapped in a helical ripple well, those ions can drift toroidally while they rotate poloidally along the valley of the two helical winding coils [55, 56]. The poloidal precession of helically trapped ion $\omega_{ht}$ is expressed as [57],

$$\omega_{ht} = -\frac{2l}{E_{\perp} v_{\parallel} B_0 r (l/2\pi)} \cdot \omega_{h0},$$

where

$$\kappa^2 \approx \frac{1}{2\varepsilon_h v_{\perp}^2}, \quad \omega_{h0} = \frac{E}{B_0 r (l/2\pi)} \quad \text{and} \quad E = \frac{1}{2} m_h v_{\perp}^2.$$

Here, $E(\kappa)$ and $K(\kappa)$ represent the first and the second kinds of the complete elliptic integrals, respectively. The parameter $\varepsilon_h$ stands for the ratio of amplitude of the fundamental helical perturbation component $B_{21}$ to magnetic field strength $B_0$, where the two subscripts of the field amplitude stand for the poloidal and toroidal numbers of the helical field. $B_0$ is the strength of the resultant magnetic field, i.e. $B_0 \approx B_{00}$. For simplicity, we discuss the deeply helically trapped limit, that is, $v_{\parallel} \sim 0$. The minor radius and magnetic field strength are chosen as the peak position of RIC, $r = 0.54$ m and $B_0 = 0.6$ T, respectively. Since the toroidal precession frequency of the helically trapped particle $\omega_{ht}$ is approximately expressed as $\omega_{ht} = (l/M) \cdot \omega_{h0}$, $\omega_{h0}$ is estimated to be $\sim 10$ kHz for the helically trapped ion of $E = 130$ keV using $\varepsilon_h = 0.2$, which is the value at $r/a = 0.9$ under the present plasma condition [53]. Since $\omega_{ht}$ is much higher than the mode frequency ($\sim 2$ kHz), the trapped ion would not be a candidate for resonant interaction. However, this kind of resonant interaction may not be ruled out completely, because the above discussion is still crude and more detailed particle orbit calculations are needed. Again, it should be noted that toroidally trapped energetic ions are an extremely small fraction of energetic ions in the LHD because of tangential beam injection and the relatively large aspect ratio of the plasma ($R_0/r \sim 6$).

As another more likely mechanism, energetic ions travelling near the plasma edge as shown in figure 1(a) may fall into an escaping orbit easily by the large amplitude of magnetic
fluctuations of the RIC mode near the edge, which are related to the eigenfunction shown in figure 11(b) \((r/a \sim 0.9)\) [37]. This loss is thought to be caused by stochasticization of the orbit of the energetic ion by RICs.

4. Summary

In NB-heated plasmas in the LHD, anomalous losses of energetic ions due to TAE and RIC modes were observed using the SLIP placed on the outboard side of the LHD. The radial profile of Hα fluctuations due to TAE was derived through a correlation analysis between Hα fluctuations and magnetic fluctuations. The radial structure of Hα fluctuations calculated using the eigenfunction of TAE obtained with the AE3D code agrees well with the experimentally obtained Hα fluctuation profile. This comparison has demonstrated that the FHA newly installed on the LHD is a powerful diagnostic tool to get information on the TAE internal structure in high beta plasmas obtained at considerably low toroidal fields. The energetic ions lost by TAE bursts belong to a class of co-passing particles. From the temporal analysis of SLIP signals obtained by multi-PMTs, it was shown that \(E\) and \(\chi\) of the lost energetic ions are quickly changed by each TAE burst. The relation between the magnetic fluctuation amplitude and the increment of loss flux by TAE is approximately quadratic and suggests a diffusive-type loss. High-frequency fluctuation components in magnetic fluctuations were observed with low toroidal fields, which are related to the edge, and would also be important for tokamak plasmas. This is one of the future interesting research targets.

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Appendix A.

The amplitude of Hα fluctuations \(\tilde{I}_{\text{Hα}}\) induced by AE is expressed as

\[
\tilde{I}_{\text{Hα}} \propto \int n_0 \tilde{n}_e (n_{e\text{c}} v_e) \, dl,
\]

where, \(n_0\), \(n_e\), \(l\) and \((n_{e\text{c}} v_e)\) indicate the neutral density, electron density, the path length along the line of sight of an Hα detector and the electron-impact excitation rate coefficient averaged over the Maxwell distribution function. The subscript AE stands for the contribution from AE activities. In this formula, the fluctuations of neutral density and \((n_{e\text{c}} v_e)\) are assumed to be negligible because \(T_e\) is in the range from 10 to 700 eV under the present experimental condition, then

\[
\tilde{I}_{\text{Hα}} \sim \frac{1}{\omega} \left( \frac{m \cos \theta \phi}{B_{\phi}} - \frac{i \sin \theta \partial \phi}{B_{\phi} \partial r} \right) \xi_r. \tag{A8}
\]

where \(k, v, e, \phi, B\) and \(\omega\) indicate the wave number, velocity, electrical field induced by TAE, the electrostatic potential, magnetic field strength and the TAE frequency. With large aspect ratio approximation \((R \gg r)\), the displacement for TAE in cylindrical coordinates \((r, \phi, \theta)\) with low \(n\) can be simply expressed as

\[
\xi_r \sim \frac{m \phi}{r \omega B_{\phi}}, \tag{A9}
\]

The displacement induced by TAE is approximately derived from the scalar potential obtained from the AE3D code as

\[
\xi_{\text{AE}} = \int v_k \, dt \approx - \int \frac{(e_k \times B)}{B^2} \, dt = - \int \frac{(\nabla \phi \times B)}{B^2} \, dt = - \frac{i \phi_k (k \times B)}{\omega B^2}. \tag{A7}
\]
where \( i = \sqrt{-1}; \) then, (A6) is rewritten as

\[
\tilde{l}_{\text{He-AE}} \sim \frac{1}{\omega} \int n_0 \frac{m}{B_p} \left[ -2 \frac{r}{R} n_e \left( m \cos \theta \frac{\phi}{r} - i \sin \theta \frac{\partial \phi}{\partial r} \right) + \frac{dn_e}{dr} \frac{m \phi}{r} \right] \, dr.
\]

Hence, the amplitude of \( \text{He} \) fluctuation is expressed as

\[
\tilde{l}_{\text{He-AE,dms}} \sim \frac{1}{\omega} \int n_0 \frac{m \phi}{B_p} \sqrt{\frac{m^2 \phi^2}{r^2} \left( \frac{dn_e}{dr} - \frac{2}{\omega} n_e \cos \theta \frac{\phi}{r} + \left( \frac{2}{\omega} n_e \sin \theta \frac{\partial \phi}{\partial r} \right)^2 \right)^2} \, dl.
\]

(A10)

(A11)

Appendix B.

The dispersion relation of RIC mode in a cylindrical plasma is given as follows [59]:

\[
2 \rho \left( 1 + \frac{g}{\omega} + \rho \frac{m^2}{r^2} \right) \frac{e^{i(\phi \tau)/4}}{\varepsilon_{k_1'}} \left( \frac{\omega_\nu}{\omega} + \rho \frac{m^2}{r^2} \right),
\]

(B1)

\[
\rho = \rho_e / a,
\]

(B2)

\[
g = \frac{m \Omega}{r} \frac{m \Omega}{dr},
\]

(B3)

\[
\omega_\nu = -\frac{m}{r} \frac{1}{n_0} \frac{dn_0}{dr},
\]

(B4)

\[
k_1' = (m n - n)/(r - r_0),
\]

(B5)

where \( \nu, \varepsilon, \omega, \rho_e, a, \Omega, r, r_0, m \) and \( n \) indicate the ion–electron collision frequency normalized by electron cyclotron frequency, the inverse of aspect ratio, the frequency of the mode, the bulk ion Larmor radius evaluated with the electron temperature, the plasma minor radius, the curvature of the magnetic field, the radial position, the location of the rotational surface, and the poloidal and toroidal mode numbers, respectively. Index ‘0’ indicates the rational surface. The mode is of interchange type and therefore localizes around the rotational transform (\( i/2\pi = 1/q; q \) is the safety factor).

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