Confinement Relevant Alfvén Instabilities in Wendelstein 7-AS

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Abstract. Bursting Alfvénic activity accompanied by strong thermal crashes and frequency chirping in a W7-AS shot is studied. A theory explaining the experimental observations is developed. A novel mechanism of anomalous electron thermal conductivity is found. In addition, a general consideration of the influence of the gap crossing on the Alfvén continuum in stellarators is carried out and a phenomenon of gap annihilation is predicted.

1. Introduction

Fast-ion-driven Alfvén instabilities are frequently observed in W7-AS [1] since neutral beam heating with $P_{\rm NBI} \lesssim 3.5\,\mathrm{MW}$ is used. In the majority of the cases, Alfvén modes affecting the energy and the fast particle confinement occur transiently during the build-up of optimized discharges. The relative fraction of fast particles is small at high densities and relatively low temperatures, in particular, during high-β operation, which explains why Alfvén instabilities are usually stabilized close to these operational boundaries. On the other hand, we should note that the fast ion population grows with temperature (it is proportional to $T_e^{3/2}$) and with increasing injection power; therefore, one can expect that fast-ion-driven instabilities will arise in future machines even in regimes with high β. The frequencies of the observed instabilities vary in a wide range from few tens of kHz to hundreds of kHz. A most interesting feature of these instabilities is that they can result in thermal crashes (the temperature can drop by up to 30%, see, e.g., Fig. 1). Note that this effect has never been observed in tokamaks (where Alfvén instabilities may influence the confinement of fast ions but not of the bulk plasma; only weak local thermal crashes during bursts of TAE (Toroidicity-induced Alfvén Eigenmodes) were observed in D-IIID [2]). In this work we analyse a low-frequency Alfvén instability observed in a particular shot of W7-AS and suggest an interpretation of this observation: we identify the instability, calculate its growth rate, suggest an explanation of the observed strong frequency chirping, and analyse the mechanisms that can result in thermal crashes. In addition, we present new features of the Alfvén continuum in stellarators.

2. An NBI-driven Alfvén instability accompanied by thermal crashes in the W7-AS shot #34723

2.1. Identification of the instability. Mechanism of the frequency chirping

The instability in the W7-AS shot #34723 (see Fig. 1) had a bursting character and was characterized by strong frequency chirping down, from about 70 kHz to 45 kHz, the instability being strongest at the final stage of the bursts, when thermal crashes occurred. The duration of the instability bursts was about 2.5 ms, the repetition period of the bursts was 8–10 ms. The plasma energy grew with time in average. Eventually, the bursts stopped, and only relatively weak steady-state Alfvénic activity with the frequency equal to that at the beginning of the bursts remained, without visible influence on plasma

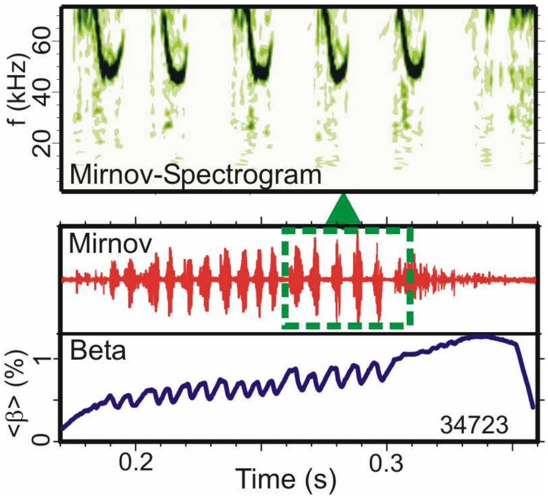

FIG. 1. Bursts of Alfvén instabilities in the W7-AS shot #34723.

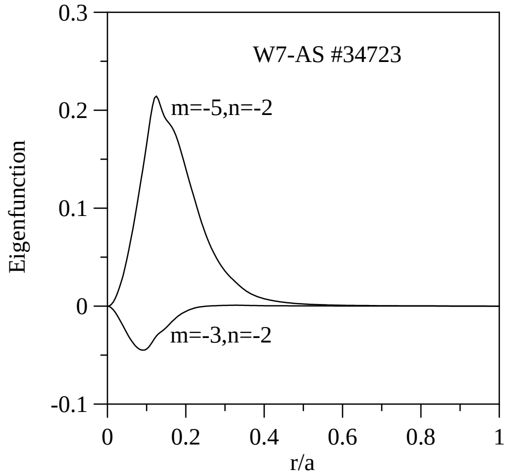

FIG. 2. Scalar potential of an Alfvén eigenmode calculated with the code BOA for the W7-AS shot #34723.

at this stage. The dominant poloidal mode numbers during the instability were $m_1 = 3$ and $m_2 = 5$.

Although the mentioned harmonics are coupled through the ellipticity coupling number, $\mu = m_2 - m_1 = 2$, the Ellipticity-induced Alfvén Eigenmode (EAE) instability was dismissed as explanation since the cylindrical branches of the Alfvén continuum (AC) with $m_1 = 3$ and $m_2 = 5$ intersect at $\iota = 0.5$, whereas $\iota(r)$ was predited less than 0.5 from equilibrium calculations. An eigenmode analysis carried out with the code BOA revealed eigenmodes with a toroidal mode number $n = 2$ (an eigenmode with the frequency near the AC branch $m/n = 5/2$ is shown in Fig. 2), the eigenmode frequencies being close to maxima/minima of the corresponding continuum branches. These branches, as well as the whole AC in the region below 120 kHz, were calculated with the AC code COBRA [3], see Fig. 3. Both Fig. 2 and Fig. 3 are obtained for a plasma with a slightly non-monotonic rotational transform [a profile with $\iota(0) = 0.447$, $\iota_{\max} = 0.45$ at $r = 0.25a$ and $\iota(a) = 0.405$, where a is the minor plasma radius, was used in Fig. 2; a similar profile but with $\iota_{\max} = 0.46$, in Fig. 3]. It was taken into account that the plasma consisted of the mixture of deuterium and hydrogen: in the considered shot, protons with the maximum energy of 48 keV were injected tangentially (the injection was balanced) into a deuterium plasma. The estimated fraction of hydrogen at $t = 0.29\,$s, immediately after a burst of the instability, is 20% in average and 60% in the plasma centre if the radial distribution of thermalized protons were as peaked as the beam ion distribution. This implies that the error bar in the Alfvén frequency in Figs. 2 and 3 is 15%. Note that the toroidal mode numbers are not known from the experiment. Therefore, we considered various magnitudes of n. We concluded that the mode with $n = 1$, which, on the first sight, is a good candidate because it has the frequency in the range of interest, was not observed. The basis for this conclusion was that the $m/n = 1/1$ AC branch lies much closer to the $m/n = 3/1$ branch than the $m/n = 5/1$ branch; thus, a satellite with $m = 1$ rather than with $m = 5$ should be present in the Fourier spectrum if $n = 1$; however, the $m = 1$ harmonic was not detected in the experiment. It follows from our analysis that the observed mode is not a gap mode; therefore, it can be identified as an unusual Global Alfvén Eigenmode (GAE) with the frequency lying above the Alfvén continuum.

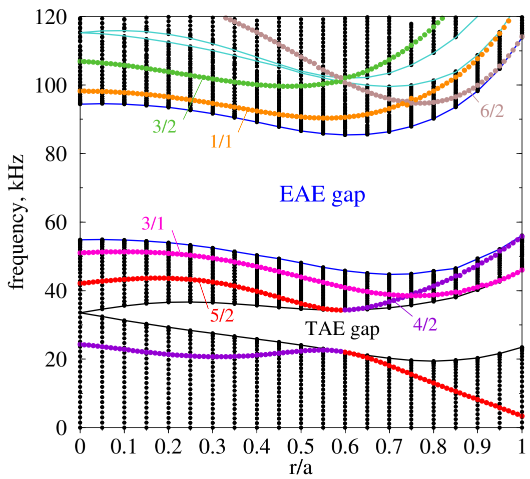

FIG. 3. The Alfvén continuum calculated by the code COBRA for the W7-AS shot #34723. Black vertical lines, AC at several radii; thick curves, selected AC branches; thin curves, boundaries of main gaps.

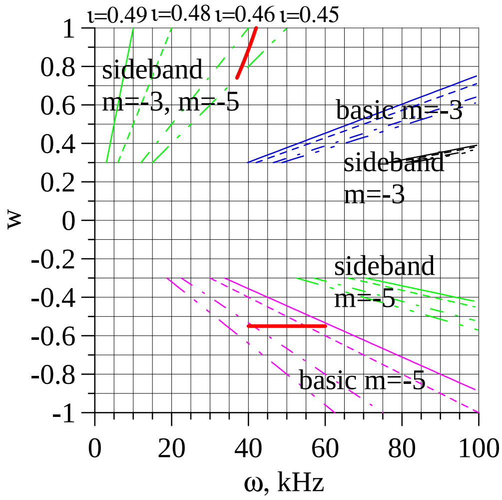

FIG. 4. The resonance velocity, $w \equiv v_\parallel/v_0$, vs ω for various ι_{max} in the W7-AS shot # 34723. The red curves show resonance velocities for the mode localized at ι_{max} and $40\,kHz < \omega = 0.82|k_\parallel(\iota_{max})|v_A < 60\,kHz$.

The Alfvén instabilities can be driven by fast ions when the following resonance condition is satisfied [4]:

$$\omega = [m\iota - n \pm (\mu_r\iota - \nu_r N)]v_\parallel/R, \tag{1}$$

where N is the number of the equilibrium magnetic field periods along the major azimuth of the torus, R is the major radius of the torus, μ_r and ν_r are the resonance coupling numbers determined by Fourier harmonics of the field line curvature [5]. In the considered shot the resonance with $\mu_r = 1$, $\nu_r = 0$ can lead to the strongest instability because, first, it involves beam ions with the energy close to the injection energy, $\mathcal{E}_0 = 48\,\mathrm{keV}$ (see Fig. 4), and second, the toroidicity-induced Fourier harmonic of the magnetic field (and, thus, the curvature) belongs to the dominant harmonics. Another resonance involving the particles with energy $\mathcal{E} \lesssim \mathcal{E}_0$ is the $\mu_r = 0$, $\nu_r = 0$ resonance (the "basic resonance"). However, the corresponding curvature harmonic is considerably smaller than the toroidal harmonic. Therefore, it can only lead to a relatively weak instability. Growth rates were calculated with the code GAMMA. We found that with the sideband resonance with $\mu_r = 1$, $\nu_r = 0$ being present, the instability is really very strong, $\gamma \lesssim \omega$, so that the perturbative approach used here is marginally applicable. This is the case when ω is less than a certain magnitude about 50 kHz. For higher ω the instability is much weaker.

On the other hand, the mode frequency strongly depends on $\iota_{\max} \equiv \max \iota(r)$. Therefore, a small change of ι considerably changes the mode frequency. In order to see this, we assume that $v_A(r) \approx \mathrm{const}$ in the region of the instability (the plasma density profile is very flat) and that the maximum of the rotational transform evolves from ι_1 to ι_2. Then, approximating the mode frequency as $\omega \approx Ck_\parallel(\iota_{\max})v_A$ with $C = \mathrm{const}$, we obtain the corresponding change of the mode frequency:

$$\frac{\omega_2 - \omega_1}{\omega_1} = \frac{\iota_2 - \iota_1}{\iota_1 - n/m} = -\frac{1}{3} \tag{2}$$

for $m/n = 5/2$, $\iota_1 \equiv \iota_{\max}(t_1) = 0.49$, and $\iota_2 \equiv \iota_{\max}(t_2) = 0.46$. Equation (2) is in reasonable agreement (slightly overestimates) with the frequency change obtained from AC calculations with COBRA.

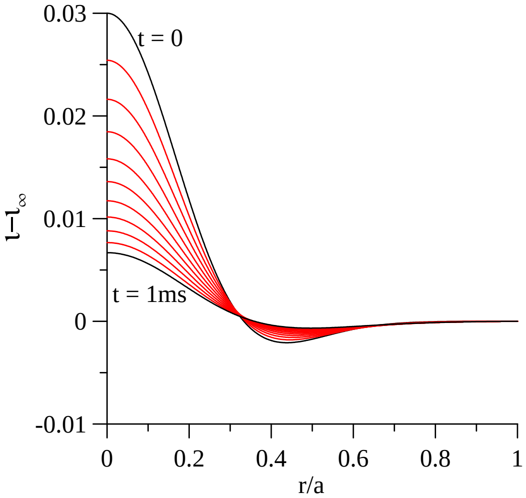

FIG. 5. Possible evolution of the rotational transform during a burst of Alfvén instability.

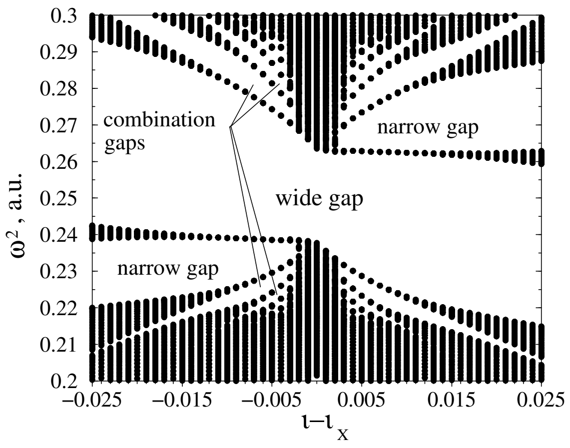

FIG. 6. Alfvén continuum near a crossing of two gaps, one of them being twice wider than the other.

The above considerations suggest the following explanation for the observed frequency chirping. The $\mu_r = 1$ resonance was responsible for the final stage of the instability, whereas other resonances [in particular, the basic resonance ($\mu_r = 0$, $\nu_r = 0$)] were responsible for the initial stage. This would be the case if the instability could influence the rotational transform, reducing its maximum magnitude, as inferred from Fig. 4. A possible mechanism leading to a change in the rotational transform is a redistribution of the beam ions with a certain sign of $v_\parallel$ by the instability. A simple estimate shows that such a redistribution can locally change the plasma current by about $\delta j = 10(1 - z_{\rm eff}^{-1})\,\mathrm{A/cm^2}$, where $z_{\rm eff} \gtrsim 2$ is the effective charge number. This is enough for the required local change of iota. For instance, $\delta\iota = 2\pi R\delta j/(Bc) = 0.04$ when $\delta j = 3.8\,\mathrm{A/cm^2}$. Moreover, the characteristic time of the evolution of the current density [and $\iota(r)$] is about $\tau_\iota \sim 1\,\mathrm{ms}$, see Fig. 5, which was obtained by solving the following model equation:

$$\frac{\partial\iota}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}\frac{D_\iota}{r}\frac{\partial}{\partial r}r^2(\iota - \iota_\infty), \tag{3}$$

where $D_\iota = c^2/(4\pi\sigma_\parallel)$, $\sigma_\parallel$ is the electrical conductivity along the magnetic field, $\iota_\infty = \iota_\infty(r)$ is the rotational transform in the infinite time after the burst. The obtained magnitude of τ_ι is a factor of two less than the duration of the instability burst. This is sufficient for ι to change during the burst.

Note that the sideband resonance involves particles with another sign of $v_\parallel$ than that for the basic resonance (see Fig. 4). This may explain why the frequency chirping down stopped after the mode frequency reached the minimum value of about 45 kHz and even a slight increase of ω was observed at the end of the instability burst.

2.2. Mechanisms of thermal crashes

2.2.1. Cooling of the plasma core due to ejection of fast ions

A simple mechanism that could be responsible for the thermal crashes during the bursts of Alfvén instabilities is the following. If a burst results in a loss of fast ions, the plasma heating decreases rapidly. This leads to a decrease of the plasma temperature, which may be significant if the energy confinement time, τ_E, is considerably less than the time between two subsequent bursts.

To study this scenario, we make some simplifying assumptions. We assume that all the injected ions have the same energy, $\mathcal{E}_0$, and pitch angle $\chi \equiv v_\parallel/v$. In addition, we assume that there is only one resonance with given μ_r, ν_r [see Eq. (1)] and that some fraction, ν_L, of the particles reaching the resonance magnitude of velocity, v_r, due to slowing down is promptly lost. Then, taking into account that $\mathcal{E}(t)\exp(t/\tau_\mathcal{E}) = \text{const}$, with $\tau_\mathcal{E}$ the particle energy loss time, during slowing down, we can write the following equation describing the temporal evolution of the beam ion energy density, W_b:

$$W_b(t) = W_{b0} \begin{cases} 1 - \frac{\nu_L \mathcal{E}_r}{\mathcal{E}_0}\left[1 - \exp\left(-\frac{t}{\tau_\mathcal{E}}\right)\right] & \text{for} \quad 0 \le t \le \tau_\text{burst} \\ 1 - \frac{\nu_L \mathcal{E}_r}{\mathcal{E}_0}\left[1 - \exp\left(-\frac{\tau_{burst}}{\tau_\mathcal{E}}\right)\right]\exp\left(-\frac{t-\tau_{burst}}{\tau_\mathcal{E}}\right) & \text{for} \quad t \ge \tau_\text{burst} \end{cases} \tag{4}$$

where W_{b0} is the beam energy density in the quiescent state, τ_burst is the burst duration. In order to describe the influence of the variation of W_b on the plasma temperature, Eq. (4) should be solved together with an equation for the plasma energy. The latter can be written as follows:

$$\frac{dT}{dt} = \frac{W_b}{3n_e\tau_\mathcal{E}} - \frac{T}{\tau_E}, \tag{5}$$

where τ_E is the plasma energy confinement time, n_e is the electron density.

Equations (4) and (5) were solved for $\tau_E = 5\,\text{ms}$, $\tau_\mathcal{E} = 3\,\text{ms}$, $\tau_\text{burst} = 2\,\text{ms}$, which corresponds to the W7-AS shot #34723, and $\nu_L\mathcal{E}_r/\mathcal{E}_0 = 0.9$. We obtained the temperature oscillations with the amplitude about 11%, which is less than what was observed experimentally. On the other hand, if we assumed that more resonances lead to particle loss, we would obtain a larger amplitude. However, the used magnitude of $\nu_L\mathcal{E}_r/\mathcal{E}_0$ seems to be unrealistically large because it implies that, at least, 90% of the resonant particles are immediately lost from the plasma. Therefore, the considered mechanism is hardly responsible for the observed thermal crashes, although it can considerably contribute to the changes of the temperature, especially, when τ_E is reduced because of the formation of magnetic islands, as described below.

2.2.2. Enhancement of electron thermal conductivity by Alfvén waves

An Alfvén mode in a stellarator includes numerous Fourier harmonics (although only several of them are significant). For some harmonics, $k_\parallel(r)$ vanishes at certain radii, around which magnetic islands may be formed. When the electrons are fast enough, so that the magnetic islands almost coincide with the islands associated with the resonances of the electron motion, $\omega = k_\parallel v_\parallel$, the electron transport may be enhanced by the destruction of the magnetic flux surfaces. However, because of the low plasma temperature in the considered experiment ($T(0) = 290\,\text{eV}$), only superthermal electrons "feel" the magnetic islands. Furthermore, the plasma density was high ($n \sim 10^{14}\,\text{cm}^{-3}$), so that the frequency of the electron–ion collisions was large, $\nu_{ei} \approx z_\text{eff} \times 10^6\,\text{s}^{-1}$. We find that $\omega^t_\text{isl} \ll \nu_{ei} \ll \omega^u_\text{isl}$, where $\omega^t_\text{isl} \sim mv_\parallel \delta\iota/R$ and $\omega^u_\text{isl} \sim mv_\parallel \iota/R$ are the characteristic frequencies of the particles trapped in the island and the untrapped particles, respectively, $\delta\iota$ is the variation of ι within the island. For $\iota = 0.45$, $\delta\iota = 0.01$, and $m = 5$, we obtain $\omega^t_\text{isl} = 2.5 \times 10^5\,\text{s}^{-1}$ and $\omega^u_\text{isl} = 10^7\,\text{s}^{-1}$. This suggests that the anomalous thermal conductivity of electrons results from the joint action of the waves and the Coulomb collisions. By similarity with the neoclassical theory, we can introduce the "anomalous collisional", "anomalous plateau", and "anomalous banana" regimes. In the experiment, the electrons are in the anomalous plateau regime, close to the boundary with the collisional one. Below we evaluate the thermal conductivity for the collisional regime, somewhat underestimating the effect of the waves.

Integrating the equation of the particle motion, $\dot{r} = v_\parallel \widetilde{B}_r/B + c\widetilde{E}_\theta/B$, we obtain

$$\xi_r = \frac{imc}{Brk_\parallel(\omega - k_\parallel v_\parallel)}\widetilde{E}_\parallel - \frac{i\widetilde{B}_r}{k_\parallel B}, \tag{6}$$

where ξ_r is the radial particle displacement, $\widetilde{\mathbf{B}}$ and $\widetilde{\mathbf{E}}$ are the perturbed magnetic and electric fields, respectively. The first term in Eq. (6) depends on the velocity of the individual particle, whereas the second one describes the oscillations of the particles together with the plasma. Therefore, only the first term describes the deflection of the particles from the perturbed flux surfaces. Using this term, we can estimate the characteristic radial displacement of a particle during its random walk, Δ. We define Δ via $\Delta^2 = \int dv_\parallel f_M(v_\parallel)|\xi_r|^2$, where ξ_r is determined by Eq. (6) with the last two terms omitted, the integral is taken over the region of untrapped particles (i.e., $|v_\parallel - \omega/k_\parallel| > 2[\omega_{Be} c\widetilde{E}_\parallel/(k_\parallel B)]^{1/2}$), f_M is the Maxwellian distribution, $\omega_{Be} = eB/(M_e c)$. As a result, we find $\Delta(\widetilde{E}_\parallel)$. Then we express $\widetilde{E}_\parallel$ in terms of $\widetilde{B}_r$, using the equations $\widetilde{E}_\parallel = 4\pi i\,\widetilde{j}_\parallel/(\omega\epsilon_\parallel)$ and $\widetilde{\mathbf{j}} = (c/4\pi)\nabla\times\widetilde{\mathbf{B}}$, where $\epsilon_\parallel = \omega_{pe}^2/(k_\parallel^2 v_e^2)$ is the plasma dielectric permeability along the magnetic field, $\widetilde{\mathbf{j}}$ is the perturbed current, ω_{pe} is the plasma frequency, $v_e = (T_e/M_e)^{1/2}$, M_e is the electron mass. We obtain the coefficient of heat conductivity of the electrons as follows:

$$\chi_e = \Delta^2\nu_e = \frac{\nu_e}{(2\pi^{1/2})}\left(\frac{m}{rk_\parallel}\right)^{1/2} k_\perp^3\rho_i^3 \frac{v_A^3}{v_e\omega^{3/2}\omega_{Be}^{1/2}}\left(\frac{\widetilde{B}_r}{B}\right)^{3/2}, \tag{7}$$

where $\rho_i = v_i/\omega_{Bi}$, $v_i = (T_i/M_i)^{1/2}$, ω_{Bi} is the ion gyrofrequency, $\nu_e = \nu_{ei} + \nu_{ee}$.

It is clear that the plasma confinement can be essentially affected by the anomalous transport only when the wave exists in a considerable part of the plasma cross section. This is not the case for the mode considered above (Fig. 2), which is localized at the plasma centre. Nevertheless, below we show that it affects a considerable part of plasma by generating a kinetic Alfvén wave (KAW). Such a wave, in addition, generates significant $\widetilde{E}_\parallel$ due to a small radial wave length.

The KAW is generated because the frequency of the considered GAE lies above the corresponding AC branch, $\omega_A(r)$ (in contrast to the conventional GAE). Our analysis shows that in this case a kinetic Alfvén wave (KAW) is generated due to "tunnel" interaction with the continuum, as in the case of TAE [6]. The amplitude of the radiated KAW depends on the distance of the mode frequency from the continuum and in our case can be comparable to the GAE amplitude. The absorption length of KAW is

$$l_{\rm abs} = 2\sqrt{\frac{2}{\pi}}\frac{v_e}{v_A}\frac{\omega}{\omega_A}\left(\frac{1}{k_\perp^2\rho_i^2} - \frac{3}{4}\right)k_r^{-1}, \tag{8}$$

where

$$k_\perp^2\rho_i^2 = \frac{\omega^2 - \omega_A^2(r)}{\omega_A^2(r) + \frac{3}{4}\omega^2}. \tag{9}$$

One can show from Eq. (9) that $k_\perp \approx 0$ at the point where the KAW is generated, and $k_{\perp\max}^2\rho_i^2 \approx 0.3$. This leads to $l_{\rm abs} \sim 10\,\mathrm{cm}$. Due to this, the destabilization of the ideal GAE localized in the central region affects a considerable part of the plasma. Analysis of the dispersion relation of KAW shows that the coupling of Fourier harmonic of the wave results in transformations of KAW branches near rational-ι points where $k_\parallel$ of the branches coincide, in our case, from a $m/n = 5/2$ into a $m/n = 4/2$ branch at $r/a \sim 0.5$;

the latter propagates to the point $r/a \sim 0.7$, where $\omega = \omega_A$ (see Fig. 3). At this point the wave is reflected because the KAW is evanescent when its frequency is below the corresponding AC branch. Thus, we conclude that the instability affects about two thirds of the plasma radius, which roughly coincides with the experimentally observed inversion radius of the thermal crashes.

Let us evaluate the amplitude of the wave required to account for the observed thermal crashes. Taking $\tau_E = a^2/(\mu_0^2\chi_e) = 1\,\mathrm{ms}$ (where $\mu_0 = 2.4$), $k_\perp^2\rho_i^2 = 1/3$, $k_\parallel = 10^{-3}\,\mathrm{cm}^{-1}$, we obtain from Eq. (7) that $\widetilde{B}_r/B = 4\times10^{-4}$. This estimate is relevant to KAW, for which $\widetilde{B}_\theta$ is few times larger than $\widetilde{B}_r$. On the other hand, Eq. (7), which was used to evaluate $\widetilde{B}_r$, underestimates the anomalous thermal conductivity because the plateau regime rather than the collisional regime takes place in the experiment. Therefore, we can expect that $\widetilde{B}/B$ required to provide $\tau_E \sim 1\,\mathrm{ms}$ will be about the above estimate for $\widetilde{B}_r/B$. This magnitude of $\widetilde{B}/B$ is quite reasonable from the point of view of available experimental data: the Mirnov-measured amplitude of $\widetilde{B}_\theta/B$ outside the plasma was as large as 10^{-4}, whereas soft X-ray measurements indicated that the instability was localized in the core, which implies that its amplitude inside the plasma well exceeded the level of 10^{-4}.

Enhancing the electron heat conductivity, the considered mechanism has almost no influence on the ion transport (on both diffusion and heat conductivity) because the ions weakly deflect from the perturbed flux surfaces, as follows from Eq. (6) (for the ions $\omega \gg k_\parallel v_i$). On the other hand, although the electron diffusion is enhanced, the electron confinement time was weakly affected in the considered W7-AS shot because of the very flat electron density profile.

3. Annihilation of gaps in the Alfvén continuum of stellarators

Calculations of the AC in W7-AS with COBRA show that the high-frequency part of AC consists of extremely thin walls (with the relative width $\Delta\omega/\omega \sim 10^{-3}$ and less) squeezed by wide gaps. Such compression of the continua (which is observed in other devices as well, e.g, in LHD [5]) completely changes the shape of AC branches, which is of importance for the properties of Alfvén eigenmodes. The wave functions of the compressed continua are characterized by strong ballooning. To elucidate the properties of such continua, we consider a model configuration with only two coupling Fourier harmonics, $\exp(i\mu_1\theta - i\nu_1 N\phi)$ and $\exp(i\mu_2\theta - i\nu_2 N\phi)$, and study the behaviour of AC near the crossing point of the two concomitant gaps.

We show that the vicinity of the crossing point is characterized by the presence of long chains of strongly coupled Fourier harmonics of the wave, coupling between different chains being much weaker. Restricting the calculations to one chain ("the chain model"), one can take into account very far interactions between the harmonics. Analytical study of the chain model reveals that the gaps "annihilate" at the crossing: The width of the joint gap equals the difference of the widths of the two separate gaps. Numerical calculations (Fig. 6) confirm this prediction and show the appearance of multiple "combination" gaps (i.e., gaps that result from the joint action of the two harmonics). The gaps are separated by narrow continuum threads, the width of each thread decreasing exponentially with $|\iota - \iota_X|^{-1}$, where X refers to the crossing point. The obtained pattern is confirmed by calculations of AC in the stellarator NCSX [7], in which the chain model was not used.

Employing the WKB approach, we suggest an analytical treatment of the problem. The problem is reduced to the Schrödinger equation with a periodic potential (like in solid state physics). The continua near the main gap correspond to the bound states trapped in the potential wells. This trapping explains the wave function ballooning observed in the W7-AS continuum. The spectral width of the states (i.e., the width of the continuum walls) is determined by tunnelling between the wells and exponentially decreases with $|\iota - \iota_X|^{-1}$, as observed in numerical calculations.

4. Summary and Conclusions

(I) An interpretation of experimental observations of Alfvénic activity in the W7-AS shot #34723 is suggested, and a relevant theory is developed.

The observed instability is identified as an unconventional GAE mode (with the frequency above AC) accompanied by the generation of kinetic waves (KAW). The calculated GAE eigenmodes are localized at the plasma centre with the dominant mode numbers $m_1 = 5$, $m_2 = 3$ (satellite harmonic) and $n = 2$. An interesting feature of KAW is that the mode with $m/n = 5/2$ is transformed into the mode with $m/n = 4/2$ at $r/a \sim 0.5$. The generation of the KAW extends the region affected by the instability up to $r \sim 0.7a$, where the radially propagating 4/2 KAW is reflected.

The performed analysis of the resonances between the waves and beam ions predicts enhancement of the instability during the frequency chirping down, which agrees with the appearance of the thermal crashes at the final stage of the instability bursts. It is shown that the frequency chirping may be caused by small local changes of the rotational transform due to the redistribution of the beam ions.

It is revealed that the presence of Alfvén waves, especially KAW, can strongly enhance the electron thermal conductivity. The considered mechanism is simple and robust: it is based on the fact that the collisional energy transfer can be strongly increased by a wave of a finite amplitude because the waves increase the excursions of the particles from the magnetic flux surfaces. The mechanism explains the oscillations of the plasma energy content without noticeable density variations in the W7-AS shot #34723.

(II) A general consideration of the behaviour of Alfvén continuum near a point where two gaps cross is carried out. The phenomenon of gap annihilation at the crossing point is predicted. An analytical description of compressed continua is suggested.

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