Development of Theory of Reversed-Shear Alfven Eigenmodes in Tokamaks

S. V. Konovalov,[1),2)] A. B. Mikhailovskii,[1)] E. A. Kovalishen,[1),3)] T. Ozeki,[2)] T. Takizuka,[2)] M. S. Shirokov[1)]

[1)]Institute of Nuclear Fusion, Russian Research Centre "Kurchatov Institute", Moscow, Russia
[2)]Naka Fusion Research Establishment, Japan Atomic Energy Research Institute, Naka-machi, Naka-gun, Ibaraki-ken, Japan
[3)]Moscow Institute of Physics and Technology, Dolgoprudnyi, Moscow Region, Russia

e-mail contact of main author: konoval@nfi.kiae.ru

Abstract. Alfven eigenmodes in tokamak discharges with reversed magnetic shear are studied theoretically by incorporating into the mode description the effects known from theories of Global, Kinetic and Toroidal Alfven Eigenmodes. In particular mode forming effect due to thermal plasma density gradient is examined. We show that this effect can essentially influence criteria on the existence of a localized solution of the eigenmode equation softening the requirements on the fast ion density and its gradient. Analysis of the finite Larmor radius effects reveals a new kinetic branch of Alfven eigenmodes, Kinetic Reversed-Shear Alfven Eigenmodes, KRSAEs. Conditions and damping rates for continuum damping of RSAE are examined. In its turn the influence of the hot ion density gradient effect, principal for the RSAE theory, is considered for TAE and EPM modes in the limit of large orbits of the hot ions.

1. Introduction

Recent experiments on large tokamaks demonstrated significant advantages of equilibrium configurations with reversed shear (RS) in achievement of reactor relevant regimes. This stimulated extensive theoretical studies of the Alfven modes in RS plasmas. The special interest among Reversed Shear Alfven Eigenmodes (RSAEs) is attracted by new magnetohydrodynamic (MHD) phenomenon, called the Alfven Cascades (ACs), which were observed on tokamaks JT-60U[1] and JET[2] in RS discharges with significant population of the hot ions. Theoretical description of ACs was suggested in a series of papers including [3-5].

As known, in order to show theoretically the existence of Alfven Eigenmodes (AEs), it is necessary to resolve the two following problems. The first one is to eliminate the singular point of the second-order differential equation for the Alfven modes. The second is to find a physical effect providing radial localization of the mode, i.e., a radial potential well. In the case of nonmonotonic (hollow) q-profile the singular point of the mode is eliminated due to that the parallel mode number $k_{||}(r)$ has an extremum related to the minimum $q(r)$. At the same time radial localization of the Alfven eigenmode structure can be supported by various physical effects. Similarly to theory of energetic particle modes[5] (EPMs), it was shown in Ref. 4 that fast ions serve not only as the drive of the AC modes but also provide their radial localization which was found to be depended on the hot ion density gradient. According to explanation of Refs. [6, 7], the physics of this effect is related to the cross-field drift of electrons compensating the electric charge of the hot ions with orbits larger than the mode radial scale. One more radial localization effect has been suggested by Breizman et al.[8] It is related to toroidal magnetohydrodynamic (MHD) corrections to the eigenmode equation which are of the second order in inverse aspect ratio expansion.

In the present report we extend the theory of AE in RS tokamak plasmas (RSAE) by incorporating the effect of thermal plasma density gradient taken from theory of cylindrical Global Alfvén Eigenmodes (GAE) (Section 2). In Section 3 we allow for kinetic (finite ion Larmor radius) effects and show a new branch of AEs, which is called the Kinetic Reversed-Shear Alfvén Eigenmodes (KRSAE). Continuum damping of RSAE is discussed in Section 4.

As incorporating effects from theories of the GAE, KAE, and TAE is shown to significantly influence the theory of RSAE it seems to be reasonable to consider what theory of the RSAE can give to its predecessors in return. Since the principal element of the theory of RSAE is the hot ion contribution into the eigenmode equation, we will discuss how this effect can modify results of the Alfven mode theory in the case of standard monotonic *q*-profile, namely, TAE (Section 5) and EPM (Section 6) modes. Summary is given in Section 7.

2. Effect of the thermal plasma density gradient on RSAE

In the recent paper [9] it was shown that taking into account the effect of thermal plasma density gradient does not change mathematical structure of RSAE equation[4,8]

$$\frac{d}{dx}\left[\left(x^2+S\right)\frac{d\phi}{dx}\right]+\left(Q_{eff}-S-x^2\right)\phi=0\,. \qquad (1)$$

Here ϕ is the electrostatic potential, x is the radial distance from the point r_0, where $q=q_{min}$, S is a parameter, characterizing the eigenvalue, $S=x_0^2\left(\left(\omega^2-\omega_A^2(0)\right)/\omega_A^2(0)\right)$, Q_{eff} is an effective potential well. The criteria for existence of a radially localized eigenfunction takes the form $Q_{eff}>1/4$, where $Q_{eff}=Q_{hot}+Q_{tor}+Q_c$ is an effective potential well given by additive contributions from the hot ions,[4,8] toroidal MHD corrections,[8] and thermal plasma density gradient,[9] $Q_c=-\left(d\ln V_A^2/dr\right)\left(q_{\min}/r_0 q''_{\min} m\right)\left(m-nq_{\min}\right)$, respectively. By comparison of relative amplitudes of these contributions we show that the localization effect of thermal plasma density gradient on AC mode can be stronger than the toroidal MHD effect as squared aspect ratio. Thus, the Alfvén Cascade modes can be theoretically demonstrated in cylindrical geometry approximation. The relative contribution of the hot ion density gradient effect, Q_{hot}, can dominate for the case of significant population of fast ions with large orbits. Such regime was realized in experiment of Ref. 2, where fast ion population was produced by powerful ICRF heating. For scenarios with parallel NBI, Q_c may compete with Q_{hot}. The mode localization effect due to the thermal plasma density gradient is localizing for the mode numbers satisfying the condition $q_{min}>m/n$. Then ACs correspond to the ``sub-Alfvénic'' modes. In the opposite condition, $q_{min} < m/n$, this effect is delocalizing. Then to provide AC eigenmodes existence other localizing mechanisms should be considered. The shift of the localization region of the eigenmodes (shift of the mode center) and the eigenfrequency shift caused by the thermal plasma density gradient were found to be sufficiently small.

3. Kinetic reversed-shear Alfven eigenmodes.

To show the KRSAE we follow (see for details [10]) the procedure of Refs. [11,12] originally developed for the problem of TAEs. Then, formally, taking into account the kinetic effect results in appearance in the l.h.s. of Eq. (1) of the term $-\Lambda\left(d^4\phi/dx^4\right)$, where

$$\Lambda = q_{\min}\rho_i^2\left(\frac{T_e}{T_i}+\frac{3}{4}\right)\frac{nq_{\min}-m}{mq''_{\min}}\left(\frac{m}{r_0}\right)^4, \qquad (2)$$

and ρ_i is the ion Larmor radius. We consider the parameter Λ to be positive, i.e. mode numbers satisfy the condition $q_{\min} > m/n$. In contrast to Eq. (1), the modified by kinetic effects AE equation has a localized solution for $Q_{eff} < 1/4$.

KRSAE modes are shown to posses the features of Alfven Cascades even for homogeneous thermal plasma density in cylindrical geometry approximation. The dispersion relation for these modes can be represented similarly to that of the KAE modes in shearless magnetic field, characterized by an effective radial wave vector $k_{r,eff}$, $\omega^2 = \omega_A^2\left(1+\rho_c^2 k_{r,eff}^2\right)$. The value $\Delta_{in} \approx 1/k_{r,eff}$ plays the role of the characteristic scale of the inner layer of KRSAE. Similarly to the case of KTAEs, this scale proves to be proportional to the squared ion Larmor radius.

4. Continuum damping of RSAE

It was taken in Refs. [4,7,13] that in the vicinity of the point r_0, the q - profile is determined by $q(r) = q_{\min} + (r-r_0)^2 q''_{\min}/2$. Representing parallel wave number $k_\parallel^2 \equiv (m/q-n)^2/R^2$ in terms of dimensionless coordinate $x = m(r-r_0)/r_0$ and expanding the squared frequency of Alfven continuum $\omega_A^2(x) \equiv v_A^2 k_\parallel^2(x)$ in a series in x^2 we get

$$\omega_A^2(x) \equiv \omega_A^2(0)\left[1-\frac{x^2}{x_0^2}\left(1-\frac{x^2}{x_1^2}\right)\right], \qquad (3)$$

where

$$\frac{1}{x_0^2} = \frac{r_0^2 q''_{\min}}{m^2 q_{\min}\left(1-nq_{\min}/n\right)}, \qquad \frac{1}{x_1^2} = \frac{1}{4x_0^2}\left[1+2\left(1-\frac{nq_{\min}}{m}\right)\right]. \qquad (4)$$

It was assumed in [4] that difference $1-nq_{\min}/m$ is small compared with unity, so that approximately $x_1 = 2x_0$. Analysis of [4,7,13] was addressed to the super-Alfvenic eigenmodes, $\omega^2 > \omega_A^2(x)$, localized in the region $x << x_1$. Meanwhile, according to Eq. (3), near the $x \approx x_1$ this condition is violated and eigenmodes should be damped due to continuum damping. The continuum damping of RSAE was studied in the recent paper [14]. Then with use of Eq. (3), eigenmode equation (1) is transformed to

$$\frac{d}{dx}\left[\left(S + x^2\left(1-\frac{x^2}{x_1^2}\right)\right)\frac{d\phi}{dx}\right] + \left(Q - S - x^2\left(1-\frac{x^2}{x_1^2}\right)\right)\phi = 0. \qquad (5)$$

In the region $x^2 >> (S,Q)$ the Eq. (5) reduces to

$$\frac{d}{dx}\left[x^2\left(1-\frac{x^2}{x_1^2}\right)\frac{d\phi}{dx}\right] + x^2\left(1-\frac{x^2}{x_1^2}\right)\phi = 0, \qquad (6)$$

which has a singular point $x = x_1$. Therefore, in order to match solutions of Eq. (6) for $x > x_1$ and $x < x_1$ we should bypass this point in correspondence with the Landau rule[15]. In the region far from the point $x = x_1$ solution of (6) was found in WKB approximation. Matching of the above solutions for $x \approx x_1$ yields

$$\phi = \hat{E} x^{-1}\left[\exp(-x) - i \operatorname{sgn}\omega \exp(x - 2x_1)\right], \tag{7}$$

where $\hat{E} \equiv E \exp(x_1)$ is a constant. This solution should be matched with solution of Eq. (1) which follows from Eq. (5) for $x << x_1$. The equation (1) was analyzed in [4,7,13] in the two limits: $Q - 1/4 << 1$ and $Q >> 1$ in assumption that its asymptotic behavior for $x >> \left(1, Q^{1/2}\right)$ is decaying and is given by Eq. (7) for $x_1 \to \infty$. In contrast to these works, we should look for solutions of Eq. (1) in the same limits but with allowance for the divergent parts of their asymptotic expressions. Then according to [14] for the case $Q - 1/4 << 1$ we arrive at the dispersion relation generalizing Eq. (F3) of Ref. [7] for the case of finite $1/x_1$:

$$(S/4)^{i\alpha} = 1 + 2i\alpha\left[2\psi(1) - \psi(1/4)\right] + \delta, \tag{8}$$

where $\psi(t) = \Gamma'(t)/\Gamma(t)$ is the psi-function, $\delta = 2\pi\alpha \exp(-2x_1)\operatorname{sgn}\omega$, $\alpha \equiv (Q - 1/4)^{1/2}$. It hence follows that, similar to Eq. (F5) of Ref. [7]

$$S = S_0\left(1 - \frac{i\delta}{\alpha}\right), \tag{9}$$

$S_0 = 256 \exp(-2\pi l/\alpha - 2C + \pi)$, $l = 1,2,3,....$, C is the Euler constant. Then with use of definition of S given after Eq. (1) we obtain from Eq. (9) the expression for the decay rate

$$\gamma \equiv \operatorname{Im}\omega = -4\pi\omega_A(0) S_0 \exp(-2x_1)/x_1^2 . \tag{10}$$

Hence we have the estimation for $Q \approx 1$:

$$\gamma \approx -\omega_A(0)\exp(-2x_1)/x_0^2 . \tag{11}$$

For the case $Q >> 1$ analysis of Ref. [14] concluded at the following dispersion relation

$$\frac{1}{\Gamma(a + 1/2)} = -\frac{i}{\sqrt{\pi}} 2^{-\frac{3}{2}a - \frac{1}{4}} Q^{\frac{a}{2}}\left[1 - \sin(\pi a)\right]\exp\left(2Q^{1/2} - 2x_1\right)\operatorname{sgn}\omega, \tag{12}$$

where $a = \hat{S}/\left(2Q^{1/2}\right)$, $\hat{S} = S - Q$, according to [4,7,13] $\hat{S}/Q$ is a small parameter. The right-hand side of this equation is proportional to the small parameter $\exp(-2x_1)$. Then Eq. (12) can be solved by the method of successive approximations in this parameter, taking $a = a^{(0)} + a^{(1)} + ...$ In the zero order we have $a^{(0)} = -l - 1/2$, l=0,2,…, wich corresponds to eigenvalue $\hat{S} = -(2l + 1)Q^{1/2}$ found in Refs. [4,7,13]. The correction $a^{(1)}$ proves to be the following: $a^{(1)} = -i\left(2^{\frac{3}{2}l + \frac{1}{2}} Q^{\frac{a}{2}} \Big/ l!\sqrt{\pi}\right)\exp\left(2Q^{1/2} - 2x_1\right)$. Then the decay rate is given by

$$\gamma = -\frac{\left(\omega^2 - \omega_A^2(0)\right)}{\omega_A^2(0)} \frac{2^{\frac{3}{2}l + \frac{1}{2}} Q^{\frac{a-1}{2}}}{l!\sqrt{\pi}} \exp\left(2Q^{1/2} - 2x_1\right) \tag{13}$$

For $Q \approx 1$ Eq. (13) gives the same estimate for the decay rate as Eq. (11). At the same time, limiting expressions for the decay rate for $Q \to 1/4$ and $Q >> 1$ prove to be essentially different. In other words, the damping of RSAE depends essentially on the value of hot ion density gradient, which in its turn provides the threshold for the very existence of the RSAEs.

The presence of the continuum damping shows that the threshold $Q > 1/4$ is insufficient for excitation of the RSAE, i.e., comprehensive theory should include an "additional threshold" for the hot ion drive to prevail over the continuum damping.

Similar results on RSAE continuum damping rates were obtained in Ref. [5] by the test function method. Besides that presented analytical results of [14] give more accurate damping rates than calculated by variational method there is also difference in the formulations of the problem in [14] and [5]. The principal point is that in Ref. [14] parameter Q proportional to hot ion density gradient is provided mostly by the indirect effect of the hot ions, namely by the effect of the cross-field drift of electrons, compensating the charge of the hot ions[6,7], while in Ref. [5] this effect was missed. More details on this subject can be found in [14] and in the Section 6 of the present report.

5. TAE in the presence of the fast ions with large orbits

The effects of the density gradient of the energetic ions with large orbits on the TAE modes was studied in Refs. [16,17]. In Ref. [16] the case of low shear, $s < \varepsilon$, was examined. In Ref. [17] analysis was extended to the $\varepsilon < s < 1$ range. Formally TAE equations with accounting for the fast ion contribution take the form

$$\left(L_m^{(0)} + L_m^{(h)}\right)\phi_m - \frac{\widetilde{\varepsilon}}{4q^2R^2}\frac{\partial^2\phi_{m-1}}{\partial r^2} = 0,$$

$$\left(L_{m-1}^{(0)} + L_{m-1}^{(h)}\right)\phi_{m-1} - \frac{\widetilde{\varepsilon}}{4q^2R^2}\frac{\partial^2\phi_m}{\partial r^2} = 0, \qquad (14)$$

which defer from the standard equations for TAE (Eqs. (29.16, 29.17) of the book [12]) by the terms $L_m^{(h)}, L_{m-1}^{(h)}$ describing the contribution of the hot ions. For the case $m >> 1$

$$L_m^{(h)} = L_{m-1}^{(h)} = \frac{\omega\kappa_h\Omega_h m}{rv_A^2}\frac{n_h}{n_c}\frac{M_h}{M_i}, \qquad (15)$$

where $\kappa_h = d\ln n_h/dr$, n_h is the hot ion density. Then the key parameter determining the features of the eigenmodes modified by the hot ion density gradient effect was found to be

$$H = -\frac{r_m\kappa_h\Omega_h}{4ms^2\omega}\frac{n_h}{n_c}\frac{M_h}{M_i}. \qquad (16)$$

It was shown[17] that TAEs are essentially modified if the population of hot ions is not too small, so that the parameter *H*, determined by Eq. (16), is of order or higher than unity, $|H| \geq 1$. An estimation can be done (cf. [7]) assuming $M_h = M_i$, $r_m\kappa_h \approx 1$, $s \approx 0.3$, $m = 3$, $n_h/n_c \approx 10^{-3}$, $\Omega_h/\omega \approx 10^3$. Then we get $H \approx 1$. Note, that, according to (16), the parameter H increases essentially as the shear decreases. Not only the amplitude but also the sign of this parameter is important. According to Eq. (16) the sign of *H* coincides with that of the mode frequency, ω, or by another words, it depends on the direction of the wave propagation with respect to equilibrium magnetic field. It was shown that depending on the sign and absolute value of the parameter H, along with standard TAEs new varieties of TAEs can be realized. These new modified TAEs, called in [17] TAE-H, TAE-H^+ and TAE-H^-, have frequencies within the Alfven gap. As for the case of standard TAE, the frequency of TAE-H^- lies near the bottom of the gap, while frequencies of TAE-H and TAE-H^+ are close to its top. New varieties of TAEs differ from standard modes by their spatial structure and energy. All new TAEs are the waves with positive energy and continuum dissipation leads to damping of all varieties of TAE. Also it was shown that similarly to standard TAEs, TAEs-H$^-$ have odd parity, while TAEs-H and TAEs-H$^+$ are even.

6. Compensating electron effect and the Energetic Particle Modes.

Recent works [18, 19] were addressed to analysis of the effect of the cross field drift of electrons compensating the charge of the fast ions (indirect effect of the fast ions, according to terminology of [7]) on the energetic-particle modes (EPMs) in the normal shear discharges. According to [7], in the limit of extremely large orbits of the fast ions, the perturbed current continuity equation with account for the compensating electron effect takes the form

$$\frac{d}{dr}\left[\left(\frac{\omega^2}{v_A^2}-k_\parallel^2\right)\frac{d\phi}{dr}\right]-k_y^2\left(\frac{\omega^2}{v_A^2}-k_\parallel^2\right)\phi+\frac{4\pi e_h \omega k_y}{cB_0}\frac{dn_h}{dr}\phi=0 \ . \tag{17}$$

Substituting here $d/dr \rightarrow ik_x$, we arrive at the local dispersion relation

$$\frac{k_\parallel^2 c^2}{\omega^2}-\frac{c^2}{v_A^2}-\varepsilon_{11}^{e,h}=0\,, \tag{18}$$

where $\varepsilon_{11}^{e,h}$ is the perpendicular dielectric permittivity of the compensating electrons determined by

$$\varepsilon_{11}^{e,h}=-\frac{4\pi e_h k_y c}{k_\perp^2 B_0\omega}\frac{dn_h}{dr}\,, \tag{19}$$

where $k_\perp^2=k_x^2+k_y^2$. Neglecting the inertial term in Eq. (18) we can qualitatively estimate the eigenfrequency of the mode provided by the compensating electron effect (Compensating Electron Alfven Eigenmode, CEAE)

$$\omega=\omega_{CE}\equiv-\frac{1}{4}\frac{k_y\omega_A^2}{\Omega_h\kappa_h}\frac{n_c M_c}{n_h M_h}\,. \tag{20}$$

With use of the expression (19) it is straightforward to show that the CEAEs are the positive energy modes.

In allowance for the hot ion contribution, one should substitute in Eq. (18)

$$\varepsilon_{11}^{e,h}\rightarrow\varepsilon_{11}^{h}\equiv\varepsilon_{11}^{e,h}+\varepsilon_{11}^{i,h}\,, \tag{21}$$

where $\varepsilon_{11}^{i,h}$ is the hot ion permittivity. As a model for $\varepsilon_{11}^{i,h}$ we take, following to [20], the expression corresponding to the approximation of strongly circulating Maxwellian ions with inhomogeneous number density n_h and uniform temperature T_h. In this case one has

$$\varepsilon_{11}^{i,h}=\frac{1}{k_\perp^2 d_h^2}\left[1-\left(1-\frac{\omega_{*h}}{\omega}\right)\delta_h\right], \tag{22}$$

where d_h is the hot ion Debye radius,

$$\delta_h=\left\langle J_0^2(\xi_{\perp h})\right\rangle+\frac{1}{\omega}\sum_{p=-\infty}^{\infty}\left\langle\frac{J_0^2(\xi_{\perp h})J_p^2(\xi_{\parallel h})(pv_\parallel/qR)^2}{\omega-k_\parallel v_\parallel-pv_\parallel/qR}\right\rangle, \tag{23}$$

$\omega_{*h} = k_y c T_h \kappa_h / (e_h B_0)$ is the diamagnetic drift frequency of the hot ions, $\xi_{\perp h} = k_\perp v_\perp / \Omega_h$, $\xi_{\| h} = k_\perp \Lambda$, $\Lambda = q\left(v_\perp^2/2 + v_\|^2\right)/\left(\Omega_h v_\|\right)$ is the radial width of the magnetic drift motion of the energetic particles.

The most important for the instability problem is the case $\omega << \omega_{*h}$. In this case the expression for ε_{11}^h takes the form

$$\varepsilon_h = \varepsilon^{e,h}\left(1 - \delta_h\right). \qquad (24)$$

The theory of EPM (see e.g. [5] and references therein) corresponds to the approximation neglecting the unity compared with δ_h. However, as can be seen from Eq. (23), for the large orbit case $\left(k_\perp \rho_h, k_\perp \Lambda_h\right) >> 1$, one has

$$\delta_h << 1. \qquad (25)$$

In other words, the term with δ_h can be treated as a small correction, which contradicts to the approximation of the EPM theory.

Then for the case of large orbits of fast ions, EPMs are substituted by the CEAEs. The latter modes of positive (in contrast to EPMs) energy are found to undergo strong continuum damping. The propagation of the CEAE was found to be opposite to hot ion diamagnetic drift direction. Thus these modes are also damped due to their resonant interaction with the hot ions. Then the general picture of Alfven instabilities in the positive shear discharges looks more favorable than predicted by the EPM theory.

7. Summary

We have extended the theory of RSAE by incorporating the mode forming effect due to thermal plasma density gradient. The effect is localizing for the mode numbers satisfying the condition $q_{\min} > m/n$ and delocalizing otherwise. The shifts of the localization region of the eigenmode and of the eigenfrequency due to the termal plasma density gradient were found[9] to be sufficiently small.

Accounting for the kinetic finite Larmor radius effect in the RSAE equation reveals[10] kinetic branch of these modes, KRSAE. Experimental confirmation of their existence could be a subject of future works.

We studied continuum damping of RSAE and calculated rigorously the damping rate. The damping rate depends in a complex way on the hot ion density and its gradient. Therefore the damping condition should modify RSAE onset threshold as well.

In the approximation of the large orbits of hot ions the dominant hot ion contribution to the RSAE equation comes from the electrons compensating the charge of the fast ions (indirect fast ion effect on the mode). Incorporating this effect in the theoretical description of the TAE modes revealed[17] new varieties of TAEs with different energy, eigenfrequency and parity of the modes.

The cross field drift of the compensating electrons was shown to be of crucial importance in the theory of EPM to provide quasineutrality of perturbations. In the limiting case of the large orbits of hot ions, the unstable EPM modes were found to be substituted by the new kind of eigenmodes, called Compensating Electron Alfven Eigenmodes[18,19] which are heavily damped due to continuum dissipation and resonance interaction with fast ions.

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