Short Wavelength Temperature Gradient Driven Modes in Tokamak Plasmas

A. I. Smolyakov,[1,2] M. Yagi,[3] and Y. Kishimoto[2]

[1]Department of Physics and Engineering Physics, University of Saskatchewan, Saskatoon, S7N5E2 Canada
[2]Department of Fusion Research, Japan Atomic Energy Research Institute, Naka, Japan
[3]Research Institute for Applied Mechanics, Kyushu University, Japan

(Received 24 October 2001; published 30 August 2002)

New unstable temperature gradient driven modes in an inhomogeneous tokamak plasma are identified. These modes represent temperature gradient (ion and electron) driven modes destabilized in the short wavelength regions with $k_\perp \rho_{i,e} \gg 1$, respectively. The instability occurs due to a specific plasma response that significantly deviates from Boltzmann distribution in the regions $k_\perp \rho_{i,e} \gg 1$.

DOI: 10.1103/PhysRevLett.89.125005 PACS numbers: 52.35.Kt, 52.35.Qz

Small scale instabilities driven by the ion and the electron temperature gradient instabilities [1] are believed to be responsible for particle and energy transport in a tokamak [2,3]. Both types of modes, the ion temperature gradient (ITG) and the electron temperature gradient (ETG), have been extensively studied over the past years [4–10]. In basic slab geometry, these modes are essentially sound waves destabilized by coupling to pressure fluctuations. In toroidal geometry, the modes are also driven by the unfavorable magnetic curvature, though the mode may remain essentially slablike in the region of the weak/negative shear [11,12]. In this Letter, we report on a new mechanism of the destabilization of the temperature gradient driven modes in the regime with large values of the Larmor radius parameter, $k_\perp^2 \rho_\alpha^2 \gg 1$, $\alpha = (e, i)$.

To illustrate the existence of new modes, first we consider a shearless slab case. Within the local theory, the parallel velocity and density perturbations are (see, e.g., [7,13])

$$\frac{\tilde{v}_{\alpha\parallel}}{v_{th\alpha}} = -\frac{e_\alpha \hat{\phi}}{T_\alpha} s_\alpha D_\alpha, \qquad \frac{n_\alpha}{n_0} = -\frac{e_\alpha \phi}{T_\alpha} l_\alpha - \frac{e_\alpha \hat{\phi}}{T_\alpha} D_\alpha, \tag{1}$$

$$\begin{aligned} D_\alpha = {} & \left(1 - \frac{\omega_{n\alpha}}{\omega}\right)[1 + s_\alpha Z(s_\alpha)]\Gamma_0(b_\alpha) \\ & + \frac{\omega_{T\alpha}}{\omega} s_\alpha \left[\frac{1}{2} Z(s_\alpha) - s_\alpha - s_\alpha^2 Z(s_\alpha)\right]\Gamma_0(b_\alpha) \\ & + \frac{\omega_{T\alpha}}{\omega}[1 + s_\alpha Z(s_\alpha)][\Gamma_0(b_\alpha) - \Gamma_1(b_\alpha)]b_\alpha, \end{aligned}$$

$$l_\alpha = 1 - \left(1 - \frac{\omega_{n\alpha}}{\omega}\right)\Gamma_0(b_\alpha) - \frac{\omega_{T\alpha}}{\omega}[\Gamma_0(b_\alpha) - \Gamma_1(b_\alpha)]b_\alpha.$$

Various plasma parameters are defined as follows: $\omega_{n\alpha} = -k_y c T_\alpha / e_\alpha B_0 L_n$, $\omega_{T\alpha} = -k_y c T_\alpha / e_\alpha B_0 L_{T\alpha}$, $L_n^{-1} = -n_0^{-1} \partial n_0 / \partial x$, $L_{T\alpha} = -T_\alpha^{-1} \partial T_\alpha / \partial x$, $s_\alpha = \omega / k_\parallel v_{th\alpha}$, $b_\alpha = k_\perp^2 \rho_\alpha^2 / 2$, $v_{tha}^2 = 2T_\alpha / m_\alpha$, $\rho_\alpha = v_{tha} m_\alpha c / (e_\alpha B_0)$; $\Gamma_{0,1}(b) = I_{0,1} \exp(-b)$, $\hat{\phi} = \phi - \omega/(k_\parallel c) A$ is an auxiliary potential, ϕ is the electrostatic potential, A is the magnetic vector potential, and $Z(s)$ is the standard plasma dispersion function.

By using Ampère's law and the Poisson equation, we obtain a general dispersion equation

$$\begin{aligned} & k_\perp^2 \delta^2 (l_i \tau + l_e + D_i \tau + D_e) - 2 s_e^2 (D_i \tau + D_e)(l_i \tau + l_e) \\ & = - k_\perp^2 \lambda_D^2 [k_\perp^2 \delta^2 - 2 s_e^2 (D_i \tau + D_e)], \end{aligned} \tag{2}$$

where λ_D^2 is the Debye length, $\lambda_D^2 = T_e/(4\pi n_0 e^2)$, $\delta^2 = c^2/\omega_{pe}^2$, and $\tau = T_e/T_i$. This general dispersion equation applies to both standard ITG and ETG modes as well as to new short wavelength modes. A similar dispersion equation was analyzed in Ref. [13].

The dispersion Eq. (2) is solved as a function of the $k_y \rho_i$ for fixed plasma parameters [13]: $\beta = 2 \times 10^{-4}$, $\rho_i/L_n = 2\sqrt{2} \times 10^{-2}$, $\rho_i/L_{Ti} = \sqrt{2} \times 10^{-1}$, $\rho_i/L_{Te} = \sqrt{2} \times 10^{-1}$, $k_x \rho_i = \sqrt{2} \times 10^{-1}$, $k_\parallel \rho_i = 2\sqrt{2} \times 10^{-3}$, $\tau = 1$, and $\lambda_D = 0$. The mode frequency and growth rate normalized to $k_\parallel v_{the}$ are shown in Figs. 1(a) and 1(b), respectively. Two new unstable branches exist in the regions $k_y \rho_i \geq 1$ and $k_y \rho_e \geq 1$. Recently, it has been emphasized that the ETG modes can be significantly modified by a finite value of the Debye length parameter [11,14]. Numerical solution of (2) shows that the electron short wavelength mode is strongly stabilized in a high temperature plasma for $\lambda_D/\rho_e \geq 1$; it is completely suppressed for $\lambda_D/\rho_e \geq 3$.

New modes occur due to a specific plasma response for $k_y \rho_\alpha > 1$. A standard notion is that for large values of the Larmor radius parameter, $k_y \rho_\alpha > 1$, the density response of the respective plasma component is Boltzmann due to decaying asymptotics of $\Gamma_{0,1}(b_\alpha) \sim 1/\sqrt{b_a}$ for large b_α. This, in fact, implicitly assumes that the ratio of $\omega_{n\alpha}/\omega$ is finite for large $k_y \rho_\alpha$. In turn, this requires that the mode eigenfrequency increases with $k_y \rho_\alpha$ (linearly or faster). The latter is true for drift wave type modes, where $\omega \sim \omega_{*\alpha}$; however, temperature gradient driven modes are basically sound waves whose frequency is of the order of $k_\parallel v_{th\alpha}$. If the wave frequency ω remains approximately constant, the ratio $\omega_{n\alpha}/\omega$ will increase with $k_y \rho_\alpha$ that compensates for the decaying $1/\sqrt{b_a}$ factor from $\Gamma_{0,1}(b_\alpha)$. Then the response functions are simplified for $k_y \rho_\alpha \gg 1$, giving

$$l_\alpha + D_\alpha = 1 + \frac{1}{2\sqrt{\pi}} \frac{e_\alpha}{e} \frac{v_{t\alpha}}{\omega L_{n\alpha}} s_\alpha Z(s_\alpha)\left(1 - \frac{\eta_\alpha}{2}\right) - \frac{1}{2\sqrt{\pi}} \frac{v_{t\alpha}}{\omega L_{T\alpha}} s_\alpha \left[\frac{1}{2} Z(s_\alpha) - s_\alpha - s_\alpha^2 Z(s_\alpha)\right], \tag{3}$$

where $\eta_\alpha = L_{n\alpha}/L_{T\alpha}$.

For the electron short wavelength branch, the ions can be taken adiabatic. Then in the electrostatic limit and $\lambda_D = 0$, $\tau = 1$, the dispersion equation reduces to $1 + l_e + D_e = 0$. Solution of this dispersion equation with (3) is shown in Fig. 1 by squares. In the fluid limit $\omega \gg k_\parallel v_{te}$ plasma dispersion functions can be simplified, giving

$$2 + \frac{1}{2\sqrt{\pi}} \frac{v_{te}}{\omega L_n}\left(1 - \frac{\eta_e}{2}\right) + \frac{1}{4\sqrt{\pi}} \frac{1}{s_e^2} \frac{v_{te}}{\omega L_n}\left(1 + \frac{\eta_e}{2}\right) = 0. \tag{4}$$

For the ion short wavelength branch $k_y \rho_i \gg 1$, electron finite Larmore radius is not important, $k_\perp \rho_e < 1$, and electron Landau damping can be taken into account in the first order in $\omega/k_\parallel v_{te} < 1$. Then the electron response function in this regime is

$$l_e + D_e = 1 + i s_e \sqrt{\pi}\left(1 - \frac{\omega_{ne}}{\omega} + \frac{\omega_{Te}}{2\omega}\right). \tag{5}$$

Electrostatic limit of (2) with the ion (3) and electron (5) response, and $\lambda_D = 0$, gives the ion mode frequency shown in Fig. 1 by circles.

In the fluid limit $\omega \gg k_\parallel v_{ti}$, one can get the following dispersion equation for the ion mode:

$$\frac{\omega_{ni}}{\omega} \frac{1}{2 s_i^2} \frac{1}{\sqrt{\pi} k_\perp \rho_i}\left(1 + \frac{\eta_i}{2}\right) + 2 + \frac{\omega_{ni}}{\omega} \frac{1}{\sqrt{\pi} k_\perp \rho_i}\left(1 - \frac{\eta_i}{2}\right) - i s_e \sqrt{\pi} \frac{\omega_{ne}}{\omega}\left(1 - \frac{\eta_e}{2}\right) = 0. \tag{6}$$

In the leading order one has from this equation $\omega^3 = v_i^3 k_\parallel^2 (1 + \eta_i/2)/(8\sqrt{\pi} L_n)$. The ion short wavelength instability exists even for adiabatic electrons; however, the mode growth rate is further increased due to the electron Landau damping reaching the maximum at $\eta_e \sim 5$. Effect of the ion temperature gradient on the ion short wavelength mode eigenfrequency is shown in Fig. 2 as a function of the $k_y \rho_i$ for a different value of η_i [from Eq. (2)]. For $\eta_i = 3$,

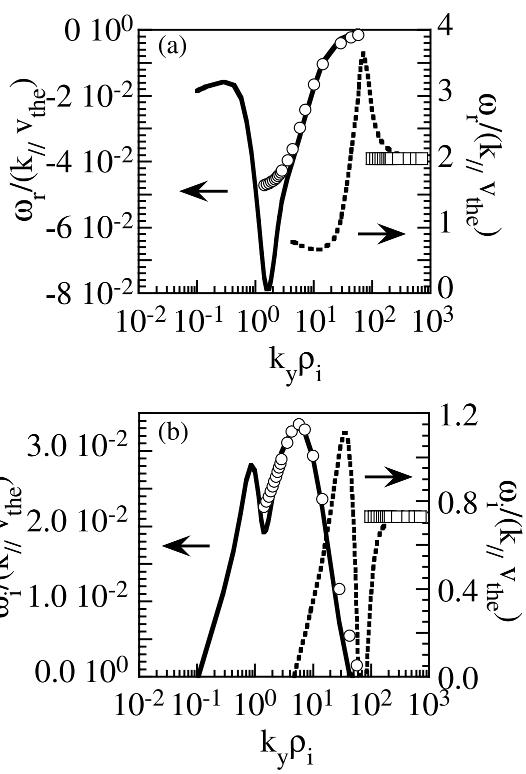

FIG. 1. Normalized wave frequency (a) and growth rate (b) for ITG (left panel) and ETG (right panel) modes. Solid line: standard ITG and short wavelength ion mode; dotted line: standard ETG and short wavelength electron mode; circle: simplified model $l_e + D_e + \tau(l_i + D_i) = 0$ with (3) and (5); square: $1 + l_e + D_e = 0$ with (3).

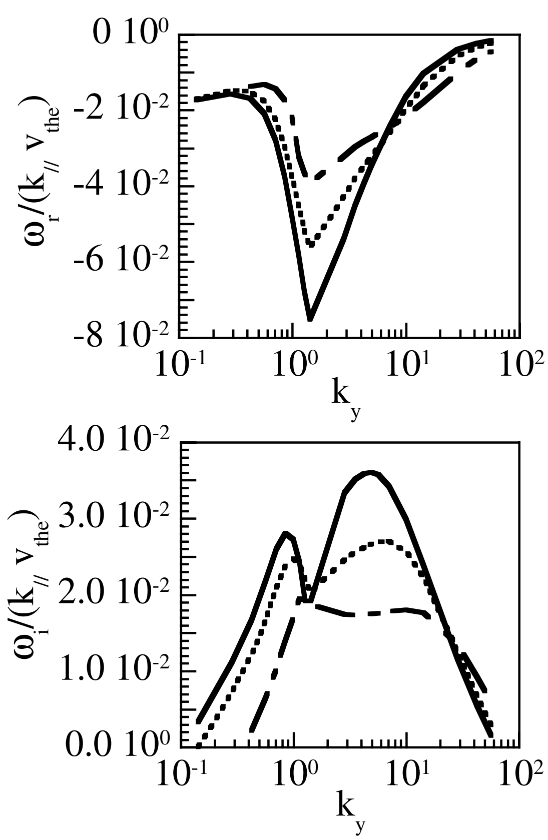

FIG. 2. The $k_y \rho_i$ dependence of the normalized real ω_r and imaginary ω_i parts of the eigenfrequency for ion mode. Solid line: $\eta_i = 5$; dotted line: $\eta_i = 4$; dotted-dashed line: $\eta_i = 3$.

the second peak disappears; however, there is still an instability in the short wavelength region.

The above shearless slab analysis is applicable to tokamak plasmas in the regions of a weak/negative magnetic shear located around the minimum q surface where the parallel wave vector $k_\| \neq 0$. Such a situation may occur near the double mode-rational surface [11,12,15] in the reversed shear regimes. Toroidal drift and mode coupling effects are weak in this region and can be neglected.

To investigate how the short wavelength modes may be affected by a finite magnetic shear and magnetic drift effects, we consider a simple nonlocal model based on the differential eigenmode equation. To illustrate the basic destabilization mechanism due to ion dynamics, we assume adiabatic electrons. A standard gyrokinetic equation in the ballooning space [4,8,16,17] gives

$$2\phi(\theta) = \int F_m J_0(k_\perp v_\perp/\omega_c) d^3v \frac{\omega - \hat{\omega}_*}{\omega - \hat{\omega}_D + i v_\| / qR\partial/\partial\theta} \times [J_0(k_\perp v_\perp/\omega_c)\phi(\theta)], \tag{7}$$

where $\hat{\omega}_* = \omega_{ni} + \omega_{Ti}(v^2/v_{ti}^2 - 3/2)$, $\hat{\omega}_D = \omega_D(\cos\theta + s\theta\sin\theta)(v_\perp^2/2v_i^2 + v_\|^2/v_i^2)$, and θ is the ballooning space variable. In the fluid limit $\omega > k_\| v_{ti}$ and $\omega > \omega_D$ one can obtain from (7) the following eigenmode equations [18,19]

$$2\phi(\theta) = \left(1 - \frac{\omega_{ni}}{\omega}\right)\Gamma_0(b)\phi(\theta) - \frac{\omega_{Ti}}{\omega} G_1 \phi(\theta) + \left[\left(1 - \frac{\omega_{ni}}{\omega}\right) G_2 - \frac{\omega_{Ti}}{\omega} G_3\right]\frac{\omega_D}{\omega}(\cos\theta + s\theta\sin\theta)\phi(\theta)$$
$$- \frac{v_{ti}^2}{q^2R^2\omega^2}\left(1 - \frac{\omega_{ni}}{\omega}\right)\left[C_1\frac{\partial^2\phi(\theta)}{\partial\theta^2} + C_2\frac{\partial\phi(\theta)}{\partial\theta} + C_3\phi(\theta)\right] + \frac{v_{ti}^2}{q^2R^2\omega^2}\frac{\omega_{Ti}}{\omega}\left[D_1\frac{\partial^2\phi(\theta)}{\partial\theta^2} + D_2\frac{\partial\phi(\theta)}{\partial\theta} + D_3\phi(\theta)\right]. \tag{8}$$

Here $\omega_D = 2\varepsilon_n\omega_{ni}$ is the toroidal drift frequency, $\varepsilon_n = L_n/R$ is the toroidicity parameter, s is the shear parameter, and various coefficients are defined as follows: $G_1 = b[\Gamma_1(b) - \Gamma_0(b)]$, $G_2 = \Gamma_0(b) + b[\Gamma_1(b) - \Gamma_0(b)]/2$, $C_1 = \Gamma_0(b)/2$, $D_1 = \Gamma_0/2 + G_1/2$, $C_2 = k_\perp^{-1}\partial k_\perp/\partial\theta G_1$, $D_2 = -k_\perp^{-1}\partial k_\perp/\partial\theta G_4$, $G_4 = 2[bG_1 + b\Gamma_0 - b\Gamma_1/2]$, $C_3 = -(k_\perp^{-1}\partial k_\perp/\partial\theta)^2 G_5 + k_\perp^{-1}\partial^2 k_\perp/\partial\theta^2 G_1/2$, $D_3 = -(k_\perp^{-1}\partial k_\perp/\partial\theta)^2 G_7 - k_\perp^{-1}\partial^2 k_\perp/\partial\theta^2 G$, $G_2 = \Gamma_0(b) + b[\Gamma_1(b) - \Gamma_0(b)]/2$, $G_3 = \Gamma_0(b) + b[3\Gamma_1(b)/2 - 2\Gamma_0(b)] + b^2[\Gamma_0(b) - \Gamma_1(b)]$, where $b = k_y^2\rho_i^2(1 + s^2\theta^2)/2$.

The long wavelength limit of Eq. (8) has been considered in [4,19]. In this Letter, we consider a short wavelength limit $k_y^2\rho_i^2 > 1$, where we obtain for the function $F = \phi/\sqrt{k_\perp}$ the following Weber equation:

$$A\frac{\partial^2 F(\theta)}{\partial\theta^2} + \left(a_1 + \frac{3}{4}\frac{\omega_D}{\omega}a_0 - 2\sqrt{\pi}k_y\rho_i\right)F + \theta^2\left[\frac{3}{4}\frac{\omega_D}{\omega}a_0\left(s - \frac{1}{2}\right) - \sqrt{\pi}k_y\rho_i s^2 + \frac{v_i^2}{q^2R^2\omega^2}a_2\frac{s^4k_y^2\rho_i^2}{4} - \frac{3}{4}As^4\right]F = 0. \tag{9}$$

Here, $A = -v_i^2/(2q^2R^2\omega^2)^{-1}a_0$, $a_0 = 1 - (1 + \eta_i/2)\omega_{ni}/\omega$, $a_1 = 1 - (1 - \eta_i/2)\omega_{ni}/\omega$, and $a_2 = 1 - (1 + 3\eta_i/2)\omega_{ni}/\omega$. We have also used the expansion $\theta^2 < 1$. The last two terms in Eq. (9) are due to the C_2, C_3, D_2, and D_3 terms in (8). From (9) we have the dispersion equation

$$a_1 + \frac{3}{4}\frac{\omega_D}{\omega}a_0 - 2\sqrt{\pi}k_y\rho_i = -i\left\{A\left[\frac{3}{4}\frac{\omega_D}{\omega}a_0\left(s - \frac{1}{2}\right) - \sqrt{\pi}k_y\rho_i s^2 + \frac{v_i^2}{q^2R^2\omega^2}a_2\frac{s^4k_y^2\rho_i^2}{4} - \frac{3}{4}As^4\right]\right\}^{1/2}. \tag{10}$$

Strong coupling approximation is valid for modes well localized in the ballooning space θ. This assumption has been confirmed for the general case of Eq. (8) with a shooting code solution. An integral equation approach for (7) also gives similar results [20]. Solution of the dispersion equation (10) in the slab limit $\omega_D = 0$ is shown in Fig. 3. As follows from Fig. 3, the mode becomes unstable when the magnetic shear parameter exceeds some threshold value that depends on η_i. The mode growth rate decreases for lower values of the magnetic shear according to (10).

The finite value of the toroidal drift frequency provides an additional destabilization mechanism, creating a toroidal branch with a larger growth rate as shown in Fig. 4. Note that there are two unstable branches. For larger values of ω_D both branches merge into an interchange type mode with a growth rate $\gamma \simeq [3(1 + \eta_i/2)\omega_D v_{ti}/(16\sqrt{\pi}L_n)]^{1/2}$. The toroidal branch is weakly affected by the magnetic shear so that the growth rate remains significant even in the region of the negative shear as shown in Fig. 4.

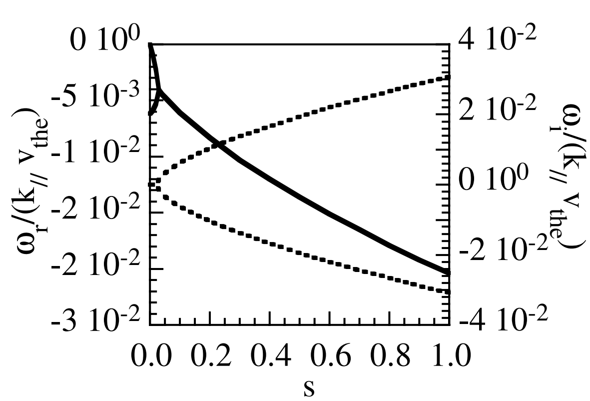

FIG. 3. Ion short wavelength mode in the fluid slab limit, $\omega_D = 0$. Real (solid line, left panel) and imaginary (dashed line, right panel) parts of the normalized wave frequency $\omega/k_\| v_{the}$ are shown as functions of the shear parameter s for $\eta_i = 3$, $\alpha_* \equiv k_y\rho_i k_\| v_{ti}/(2\omega_*) = 0.28$, and $k_y\rho_i = 4\sqrt{2}$. Here $k_\| \equiv 1/(qR)$.

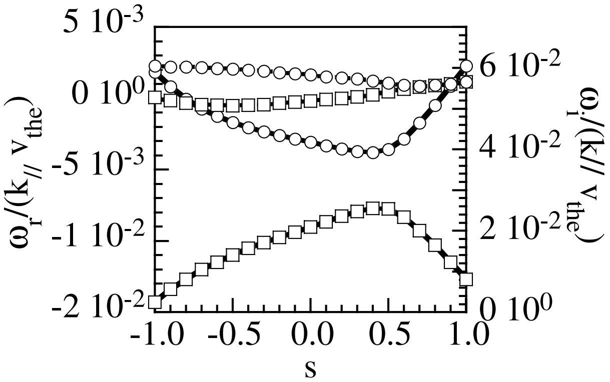

FIG. 4. Real (solid lines, left panel) and imaginary (dashed lines, right panel) parts of the normalized, $\omega/k_{\|}v_{the}$, wave frequency of the unstable ion short wavelength mode in the fluid toroidal limit as a function of the shear parameter s; $\eta_i = 3$, $k_y\rho_i = 4\sqrt{2}$, $\alpha_* = 0.28$, $k_* \equiv k_y\rho_i k_{\|}v_{ti}/\omega_D = 1$. Squares and circles indicate two different branches.

New short wavelength temperature gradient driven modes can produce a significant level of anomalous transport. A simple mixing length estimate gives for the toroidal interchange mode $D_i \simeq \rho_i^2 v_{thi}\sqrt{\varepsilon_d}/(L_n a_m^{3/2})$, where a numerical factor a_m corresponds to the wave vector with the maximal growth rate, $k_\perp\rho_i \simeq a_m > 1$. The toroidal short wavelength mode is closely related to the "ubiquitous" mode that includes the effect of trapped electrons [18]. Trapped electrons will further increase the mode growth rate and, respectively, the turbulent diffusion associated with such a mode. One of the most interesting features of the short wavelength ion mode is its ability to produce a finite level of the electron transport because the mode growth rate remains finite well into the region $k_\perp^2\rho_i^2 > 1$. The electron transport produced by the short wavelength modes will be substantially larger than that of the standard ETG modes. It is interesting to note that, for moderate plasma pressures $\beta = 1\%$–2%, the ion short wavelength mode extends into the region $k_\perp \simeq c/\omega_{pe}$ where the electron transport may be further increased by the electromagnetic effects [9,21]. A new mode existing in the intermediate region may provide coupling between standard ITG and ETG modes thus leading to a multiple-length-scale turbulence state [22].

In summary, we have identified new short wavelength branches of the temperature gradient driven modes. These modes are closely related to acoustic type modes that are destabilized by both the temperature gradient effects and plasma toroidicity and persist in the region of the weak/negative magnetic shear. A modified plasma response in the region of $k_y\rho_\alpha \gg 1$ is essential for a new instability.

The authors acknowledge useful discussions with Y. Idomura, J. Q. Li, S. I. Itoh, and K. Itoh. We are also grateful to Dr. H. Ninomiya and Dr. A. Kitsunezaki for their continuous encouragement. This research was supported by Natural Sciences and Engineering Research Council of Canada, by the Grant-in-Aid for Scientific Research and the Grant-in-Aid for International Scientific Research of Ministry of Education, Culture, Sports, Science and Technology of Japan, and by collaboration program of Research Institute for Applied Mechanics of Kyushu University.

[1] Yu. Tserkovnikov, Sov. Phys. JETP **5**, 54 (1957); L. I. Rudakov and R. Z. Sagdeev, Sov. Phys. Dokl. **6**, 415 (1961); B. B. Kadomtsev and A. B. Timofeev, Sov. Phys. Dokl. **7**, 826 (1962); A. A. Galeev, V. N. Oraevskii, and R. Z. Sagdeev, Sov. Phys. JETP **17**, 615 (1963); A. B. Mikhailovskii and L. I. Rudakov, Sov. Phys. JETP **17**, 621 (1963).

[2] A. M. Dimits *et al.*, Phys. Plasmas **7**, 969 (2000).

[3] F. Jenko, W. Dorland, M. Kotchenreuther, and B. N. Rogers, Phys. Plasmas **7**, 1904 (1999).

[4] F. Romanelli, Phys. Fluids B **1**, 1018 (1989); F. Romanelli, L. Chen, and S. Briguglio, Phys. Fluids B **3**, 2469 (1991).

[5] J. Y. Kim and W. Horton, Phys. Fluids B **3**, 3194 (1991).

[6] H. Biglari, P. H. Diamond, and M. N. Rosenbluth, Plasma Phys. **1**, 109 (1989); G. S. Lee and P. H. Diamond, Phys. Fluids **29**, 3291 (1986).

[7] Y. C. Lee, J. Q. Dong, P. N. Guzdar, and C. S. Liu, Phys. Fluids **30**, 1331 (1987); J. Q. Dong, P. N. Guzdar, and Y. C. Lee, Phys. Fluids **30**, 2694 (1987).

[8] J. Q. Dong, W. Horton, and J. Y. Kim, Phys. Fluids B **4**, 1867 (1992).

[9] W. Horton, B. G. Hong, and W. M. Tang, Phys. Fluids **31**, 2971 (1988).

[10] T. S. Hahm and W. M. Tang, Phys. Fluids B **2**, 1815 (1990); **1**, 1185 (1989).

[11] Y. Idomura, M. Wakatani, and S. Tokuda, Phys. Plasmas **7**, 2456 (2000).

[12] Y. Kishimoto *et al.*, Nucl. Fusion **40**, 667 (2000).

[13] Y. Chen and S. Parker, Phys. Fluids **8**, 2095 (2001).

[14] J. Q. Dong, H. Sanuki, and K. Itoh, Phys. Plasmas **8**, 3635 (2001).

[15] Z. Chang *et al.*, Phys. Rev. Lett. **77**, 3553 (1996).

[16] J. Y. Kim and M. Wakatani, Phys. Plasmas **2**, 1012 (1995).

[17] T. M. Antonsen and B. Lane, Phys. Fluids **23**, 1205 (1980); P. J. Catto, W. M. Tang, and D. E. Baldwin, Plasma Phys. **23**, 639 (1981); E. A. Frieman and Liu Chen, Phys. Fluids **25**, 502 (1982).

[18] B. Coppi and F. Pegoraro, Nucl. Fusion **17**, 969 (1974); B. Coppi, S. Migliuolo, and Y-K. Pu, Phys. Fluids B **2**, 2322 (1990); B. Coppi and G. Rewoldt, Phys. Rev. Lett. **33**, 1329 (1974).

[19] S. C. Guo, L. Chen, S. T. Tsai, and P. N. Guzdar, Plasma Phys. Controlled Fusion **31**, 423 (1989).

[20] A. Hirose, M. Elia, A. Smolyakov, and M. Yagi, Phys. Plasmas **9**, 1659 (2002).

[21] A. Hirose, Phys. Fluids B **3**, 1599 (1991).

[22] S.-I. Itoh and K. Itoh, Plasma Phys. Controlled Fusion **43**, 1055 (2001).