Nonideal magnetohydrodynamic Kelvin–Helmholtz instability driven by the shear in the ion diamagnetic drift velocity in a high-β plasma

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A two-dimensional nonideal magnetohydrodynamic eigenmode equation for the most dangerous perturbation ($k_{\parallel}=0$) is derived for a high-β plasma by making use of the generalized Ohm's law and becomes identical to the equation describing the two-dimensional hydrodynamic stability. When the pressure is nonuniform, the density is uniform, and there is no unperturbed electric field or gravity, a Kelvin–Helmholtz instability, which is driven by the shear in the ion diamagnetic drift velocity, is found. When the unperturbed ion pressure is proportional to the unperturbed total pressure, a necessary condition for the instability is that d^3B_0/dx^3 must change sign at least once between $x=x_1$ and $x=x_2$, where x is the direction of the nonuniformity perpendicular to the unperturbed magnetic field $B_0(x)$ and $x=x_1$ and x_2 are points where the x component of the velocity perturbation vanishes. An unstable $B_0(x)$ profile and the dispersion relation are obtained for a polygonal ion diamagnetic drift velocity profile.
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I. INTRODUCTION

The investigation of hydromagnetic stability of high-β plasmas is of interest in such varied fields as the study of magnetic confinement, space plasmas such as the auroras and magnetopause stability, and astrophysical plasmas. An ideal magnetohydrodynamic (MHD) Kelvin–Helmholtz (KH) instability driven by the shear in the $\mathbf{E}\times\mathbf{B}$ drift velocity, which occurs even in the cold magnetized plasma, has been intensively studied in fusion, space, and astrophysical plasmas. In this paper a KH instability driven by the shear in the ion diamagnetic drift velocity, which is a nonideal MHD drift in a high-β plasma, is studied for a configuration, in which the unperturbed ion pressure gradient is transverse to the magnetic field. Although the ideal MHD KH instability,[1–3] driven by the shear in the $\mathbf{E}\times\mathbf{B}$ drift velocity of ion and electron guiding centers, is a hydromagnetic counterpart of the hydrodynamic KH instability,[1] the present new instability is peculiar to the high-β plasma. This is because the ion diamagnetic drift, which is not entirely the motion of guiding centers, is an artifact of gyration, which is due to the magnetization and the gradient-B drift averaged over a thermal distribution of velocities for a straight field line geometry.[4] Although the velocity shear in the fusion plasma is known to stabilize the turbulence, the velocity shear is also important in destabilizing the KH instability and the present KH instability driven by the shear in the ion diamagnetic drift velocity may be relevant to turbulence generation in the fusion plasma and to the vortical deformation of auroral arcs following the substorm onset, which occurs in the near-Earth plasma sheet when the plasma pressure is highly disturbed.

In the present study a new KH instability is obtained by using the generalized Ohm's law and occurs due to the shear in the ion diamagnetic drift velocity and not due to the ideal MHD $\mathbf{E}\times\mathbf{B}$ drift. Therefore, the instability is essentially a nonideal MHD instability in a high-β plasma. There have been many studies investigating nonideal MHD effects on the KH instability, e.g., Hall MHD studies,[5–7] hybrid MHD studies[8–11] and kinetic studies.[12–14] However, those studies concentrated on the $\mathbf{E}\times\mathbf{B}$ driven KH instability and could not find a nonideal MHD instability. Although there is a study of a KH instability,[15] which is not driven by the $\mathbf{E}\times\mathbf{B}$ drift shear, their Hall MHD analysis is applicable only to thin current sheets, where ions are unmagnetized, and therefore their model is applicable to a parameter regime different from the present study. The stability of finite Larmor radius hydrodynamics was studied for the present configuration,[16–18] but they studied only a low-β plasma and assumed an electrostatic perturbation and that the magnetic field is constant in time.

The basic configuration for the present instability and the basic equations are presented in Sec. II. The unperturbed state for the instability is obtained in Sec. III. The two-dimensional (2D) eigenmode equation for the most dangerous mode ($k_{\parallel}=0$) is obtained in Sec. IV. A stability criterion for the present new instability is given and the dispersion relation for a polygonal ion diamagnetic drift velocity profile is presented in Sec. V. Discussion and summary are given in Sec. VI.

II. BASIC CONFIGURATION AND EQUATIONS

We consider a configuration in which the unperturbed magnetic field $\mathbf{B}_0(x)$ is in the z direction and is a function of only x. The unperturbed electric field $\mathbf{E}_0(x)$ is in the x direction and is a function of only x. The unperturbed pressure $p_0(x)$ and density $\rho_0(x)$ are functions of only x. The constant gravity acceleration $\mathbf{g}$, which may be important in astrophysical applications, is directed to $-x$ direction, i.e., $\mathbf{g}=-g\hat{\mathbf{x}}$, where $\hat{\mathbf{x}}$ is the unit vector in the x direction.

The stability of the present configuration is described by the following equations.

$$\frac{\partial \rho}{\partial t}+\nabla\cdot(\rho\mathbf{v})=0, \tag{1}$$

$$\rho\frac{d\mathbf{v}}{dt}=\mathbf{j}\times\mathbf{B}-\nabla p+\rho\mathbf{g}, \tag{2}$$

$$\nabla\times\mathbf{B}=\mu_0\mathbf{j}, \tag{3}$$

$$\frac{\partial\mathbf{B}}{\partial t}=-\nabla\times\mathbf{E}, \tag{4}$$

$$\nabla\cdot\mathbf{v}=0. \tag{5}$$

Here, ρ is the plasma density, $\mathbf{v}$ is the macroscopic velocity of the plasma, $\mathbf{B}$ is the magnetic field, p is the plasma pressure. Another equation, which relates $\mathbf{E}$ to $\mathbf{v}$ and $\mathbf{B}$, is the generalized Ohm's law, which is derived from the equation of motion for the electron fluid. From the quasineutrality one has $n_e\simeq n_i\simeq n$, where n_e and n_i are the electron density and the ion density, respectively, and also one has

$$\mathbf{j}=ne(\mathbf{v}_i-\mathbf{v}_e)\simeq ne(\mathbf{v}-\mathbf{v}_e), \tag{6}$$

where $\mathbf{v}_i$ and $\mathbf{v}_e$ are the average velocities of ion species and electron species, respectively. Substituting (6) into the equation of the electron fluid and assuming that a time scale of our concern is much longer than the electron gyration period yield

$$-ne(\mathbf{E}+\mathbf{v}\times\mathbf{B})+\mathbf{j}\times\mathbf{B}-\nabla p_e+nm_e\mathbf{g}=0, \tag{7}$$

where m_e is the electron mass and p_e is the electron pressure. Equation (7) is a generalized Ohm's law, which is used in the present analysis.

III. UNPERTURBED STATE

The unperturbed form of the generalized Ohm's law (7) can be written as

$$-n_0e(\mathbf{E_0}+\mathbf{v_0}\times\mathbf{B_0})+\mathbf{j_0}\times\mathbf{B_0}-\nabla p_{e0}+n_0m_e\mathbf{g}=0, \tag{8}$$

where the subscript 0 denotes the unperturbed state. From (2) one obtains

$$\rho_0(\mathbf{v_0}\cdot\nabla)\mathbf{v_0}=\mathbf{j_0}\times\mathbf{B_0}-\nabla p_0+\rho_0\mathbf{g}, \tag{9}$$

where $\mathbf{v_0}$ is expected *a priori* to be the sum of the $\mathbf{E}\times\mathbf{B}$ drift, the ion diamagnetic drift and the gravitational drift. Since all these drifts are in the y direction and are functions of only x, one can write in (9) $\mathbf{v_0}(x)=V_0(x)\hat{\mathbf{y}}$, where $\hat{\mathbf{y}}$ is the unit vector in the y direction. Then, Eq. (9) becomes simply

$$\mathbf{j_0}\times\mathbf{B_0}=\nabla p_0-\rho_0\mathbf{g}. \tag{10}$$

Thus, one obtains from (8) and (10),

$$\mathbf{v}_{0\perp}=\frac{\mathbf{E_0}\times\mathbf{B_0}}{B_0^2}+\frac{1}{n_0eB_0^2}\mathbf{B_0}\times\nabla p_{i0}-\frac{M}{eB_0^2}\mathbf{B_0}\times\mathbf{g}, \tag{11}$$

where $M=m_i+m_e$, m_i is the ion mass, and m_e/M was neglected, compared to (1). If there is no j_{0z}, which is parallel to the unperturbed field line, $\mathbf{j_0}$ can be written from (10) as $\mathbf{j_0}(x)=j_0(x)\hat{\mathbf{y}}$. From (3) one also obtains

$$j_0(x)=-\frac{1}{\mu_0}\frac{d\mathbf{B}_0}{dx}. \tag{12}$$

Substitution of (12) into (10) yields

$$\frac{d}{dx}\left[p_0(x)+\frac{\mathbf{B}_0^2(x)}{2\mu_0}\right]=-\rho_0(x)g. \tag{13}$$

Since $\mathbf{E_0}=E_0(x)\hat{\mathbf{x}}$ is assumed, one obtains from (11)

$$V_0(x)=-\frac{E_0(x)}{B_0(x)}+\frac{1}{n_0(x)eB_0(x)}\frac{dp_{i0}}{dx}+\frac{Mg}{eB_0(x)}. \tag{14}$$

IV. EIGENMODE EQUATION

Let us assume that a perturbation has no variation in the z direction and has the form $\delta f(x)\exp[i(ky-\omega t)]$. Since a temporal growth of the instability is considered, k is real and ω is complex. From (1) one obtains

$$-i(\omega-kV_0)\delta\rho+\delta v_x\frac{d\rho_0}{dx}=0. \tag{15}$$

In the following the curl of the generalized Ohm's law (7) and the curl of the equation of motion (2) are taken in order to delete the ∇p_e and ∇p terms, respectively. First, by taking the curl of the generalized Ohm's law (7), one obtains

$$\frac{e}{M}\rho\frac{\partial\mathbf{B}}{\partial t}-\frac{e}{M}\rho[(\mathbf{B}\cdot\nabla)\mathbf{v}-(\mathbf{v}\cdot\nabla)\mathbf{B}]-\frac{e}{M}\nabla\rho\times(\mathbf{E}+\mathbf{v}\times\mathbf{B})$$
$$+(\mathbf{B}\cdot\nabla)\mathbf{j}-(\mathbf{j}\cdot\nabla)\mathbf{B}+\frac{m_e}{M}\nabla\rho\times\mathbf{g}=0. \tag{16}$$

By taking the perturbation of (16) and after some algebra one obtains x and y components of the perturbed equation of (16) as follows:

$$F(x)\delta B_x=0, \tag{17}$$

$$F(x)\delta B_y-n_0(x)e\,\delta B_x\frac{dV_0}{dx}-e\frac{dn_0}{dx}[-\delta E_z+V_0(x)\delta B_x]$$
$$+\delta B_x\frac{dj_0}{dx}=0, \tag{18}$$

where

$$F(x)=-en_0(x)i(\omega-kV_0)-ikj_0(x). \tag{19}$$

Since $F(x)\neq 0$ for $\omega_i\neq 0$, where ω_i is the imaginary part of ω, one obtains $\delta B_x=0$ from (17). From (4) one obtains

$$i\omega\delta B_x=ik\delta E_z. \tag{20}$$

Therefore, one obtains $\delta E_z=0$. Then, substitution of $\delta B_x=\delta E_z=0$ into (18) yields $\delta B_y=0$. Therefore, in the present configuration only the z component of the magnetic field is perturbed.

Next, by taking the curl of the equation of motion (2) one obtains

$$\rho\left(\frac{\partial}{\partial t}(\nabla\times\mathbf{v})-[(\nabla\times\mathbf{v})\cdot\nabla]\mathbf{v}+(\mathbf{v}\cdot\nabla)(\nabla\times\mathbf{v})\right)+\nabla\rho$$
$$\times\left[\frac{\partial\mathbf{v}}{\partial t}+(\mathbf{v}\cdot\nabla)\mathbf{v}\right]=(\mathbf{B}\cdot\nabla)\mathbf{j}-(\mathbf{j}\cdot\nabla)\mathbf{B}+\nabla\rho\times\mathbf{g}. \tag{21}$$

By taking the perturbation of (21) and after some algebra, one obtains x and z components of the perturbed equation of (21) as follows:

$$\rho_0 k[\omega-kV_0(x)]\delta v_z=-ikj_0(x)\delta B_x, \tag{22}$$

$$\rho_0\left[-i(\omega-kV_0(x))\left(\frac{d\delta v_y}{dx}-ik\delta v_x\right)+\delta v_x\frac{d^2V_0}{dx^2}\right]$$
$$-i(\omega-kV_0(x))\delta v_y\frac{d\rho_0}{dx}+\delta v_x\frac{dV_0}{dx}\frac{d\rho_0}{dx}$$
$$=-ikj_0(x)\delta B_z-\frac{ik}{\mu_0}\delta B_z\frac{dB_0}{dx}+ik\delta\rho g, \tag{23}$$

where

$$\delta j_x=\frac{ik}{\mu_0}\delta B_z \tag{24}$$

was used in the final equation. Since $\delta B_x=0$, one obtains $\delta v_z=0$ from (22). This means that in the present configuration, the perturbed motion is confined in the $x-y$ plane, which is transverse to the magnetic field. Substitution of (15) and (12) into (23) with the aid of (5) yields

$$\frac{d^2\delta v_x}{dx^2}+\frac{1}{\rho_0}\frac{d\rho_0}{dx}\frac{d\delta v_x}{dx}-\left[k^2-\frac{k}{\omega-kV_0}\frac{d^2V_0}{dx^2}\right.$$
$$\left.-\frac{k}{\omega-kV_0}\frac{1}{\rho_0}\frac{d\rho_0}{dx}\left(\frac{dV_0}{dx}-\frac{kg}{\omega-kV_0}\right)\right]\delta v_x=0. \tag{25}$$

Notice that on the right-hand side of (23) the first two terms are cancelled because of (12) and the magnetic field $B_0(x)$ enters only in the expression of $V_0(x)$ in (14). This equation is completely the same as the eigenmode equation describing the hydrodynamic KH instability with $k_z=0$ in nonuniform fluid in a gravitational field.[1]

V. STABILITY CRITERION AND DISPERSION RELATION

Since only the KH instability is of our interest, let us assume in the following $g=0$ and $\rho_0(x)=\rho_0=\text{const}$. The neglect of the gravity force is justified in terrestrial fusion plasmas and in most space plasma applications. Then, one obtains from (25) and (14),

$$\frac{d^2\delta v_x}{dx^2}-\left[k^2-\frac{k}{\omega-kV_0}\frac{d^2V_0}{dx^2}\right]\delta v_x=0, \tag{26}$$

$$V_0(x)=-\frac{E_0(x)}{B_0(x)}+\frac{1}{n_0eB_0(x)}\frac{dp_{i0}}{dx}. \tag{27}$$

The equation (26) is the Rayleigh's stability equation.[19,20] We assume that $\delta v_x=0$ at $x=x_1$ and $x=x_2$. Following Rayleigh[19,20] one multiplies (26) by δv_x^*, which is the complex conjugate of δv_x, and integrates by parts the resulting equation from $x=x_1$ to $x=x_2$. Then, one obtains

$$\int_{x_1}^{x_2}\left(\left|\frac{d\delta v_x}{dx}\right|^2+k^2|\delta v_x|^2\right)dx$$
$$-\int_{x_1}^{x_2}\frac{k}{\omega-kV_0}\frac{d^2V_0}{dx^2}|\delta v_x|^2dx=0. \tag{28}$$

From the imaginary part of (28) one obtains

$$\omega_i\int_{x_1}^{x_2}\frac{|\delta v_x|^2}{|\omega-kV_0|^2}\frac{d^2V_0}{dx^2}dx=0. \tag{29}$$

From (29) a necessary condition for ω_i to be nonzero is that $d^2V_0(x)/dx^2$ must change sign at least once between $x=x_1$ and $x=x_2$. This is the Rayleigh's criterion for the instability.[19,20]

We now concentrate on a new instability, which is driven by the shear in the ion diamagnetic drift velocity. In order to show existence of such an instability, we first obtain a necessary condition for the instability and then we have to show that there is indeed an unstable configuration whose growth rate ω_i and ω_r are given explicitly. For this purpose, we consider a case where the unperturbed electric field $E_0(x)=0$ and $V_0(x)$ is given by the ion diamagnetic drift velocity as follows:

$$V_0(x)=\frac{1}{n_0eB_0(x)}\frac{dp_{i0}}{dx}. \tag{30}$$

We consider that n_0 is uniform, but $p_{i0}(x)$ is nonuniform and consider a particular case, where the ion plasma pressure $p_{i0}(x)$ is proportional to the total pressure $p_0(x)$, i.e., $p_{i0}(x)=Kp_0(x)$, where K is the constant. When both plasma species have an equal unperturbed temperature, i.e., $T_{i0}(x)=T_{e0}(x)$, where T_{i0} and T_{e0} are ion and electron unperturbed temperatures, respectively, K is equal to 1/2. However, in Earth's plasma sheet, $T_{i0}(x)$ is much larger than $T_{e0}(x)$ and K is typically 0.9. In this case $V_0(x)$ in (30) is given by

$$V_0(x)=-\frac{K}{\mu_0n_0e}\frac{dB_0}{dx}, \tag{31}$$

where (13) was used. A necessary condition for the KH instability driven by the shear in the ion diamagnetic drift velocity is therefore that d^3B_0/dx^3 must change sign at least once between $x=x_1$ and $x=x_2$. This is only a necessary condition for the instability and not the sufficient condition. In the following we take advantage of previous studies of the stability of parallel flows to obtain an unstable profile of B_0. A well-known velocity shear profile, which has an inflexion point and is unstable, is the hyperbolic tangent velocity profile.[21] For analytical tractability we adopt in the following a velocity profile $V_0(x)$ shown by the dotted line in Fig. 1. This is a function of x, which is given by

$$V_0(x)=\left\{\begin{array}{cc}0 & x\leqslant -a\\ U_0(x/a+1) & |x|<a\\ 2U_0 & x\geqslant a\end{array}\right\}, \tag{32}$$

where U_0 is a constant velocity, which is given by $U_0=-(K/2\mu_0n_0e)[dB_0/dx]_{x=a}$. This velocity profile is known to be unstable.[19,20,22] From (31) one obtains

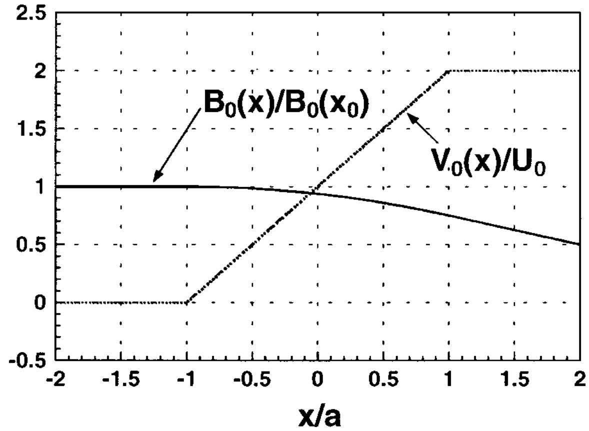

FIG. 1. Profiles of unstable $V_0(x)$ (dotted curve) and unstable $B_0(x)$ (solid curve).

$$B_0(x) = -\frac{\mu_0 n_0 e}{K}\int_{x_0}^{x} V_0(x)dx + B_0(x_0). \tag{33}$$

The solid curve in Fig. 1 shows the profile of unstable $B_0(x)$ normalized by $B_0(x_0)$, which is obtained from (33) by assuming $x_0 = -2a$ and $dB_0/dx = -B_0(x_0)/(4a)$ at $x=a$. Next, one has to obtain explicit dispersion relation for (32). A velocity profile defined by

$$V_0(x) = \left\{ \begin{array}{cc} -U_0 & x \leqslant -a \\ U_0 x/a & |x| < a \\ U_0 & x \geqslant a \end{array} \right\} \tag{34}$$

which is a shifted form of (32), is known to be unstable[19,20,22] and the dispersion relation for the velocity profile (34) is given by[19,20,22]

$$\left(\frac{\omega}{k}\frac{1}{U_0}\right)^2 = \left(1 - \frac{1}{2ka}\right)^2 - \frac{1}{4(ka)^2}e^{-4ka}. \tag{35}$$

The normalized growth rate $\omega_i a/U_0$ for (34) calculated from (35) is plotted in Fig. 2 as a function of ka for $ka < 0.64$ and the real part of the frequency ω_r for (34) is zero for $ka < 0.64$ as is expected from the anti-symmetric nature of the velocity profile (34). The velocity profile (32) is obtained from (34) by transforming the unperturbed velocity $V_0(x)$ to

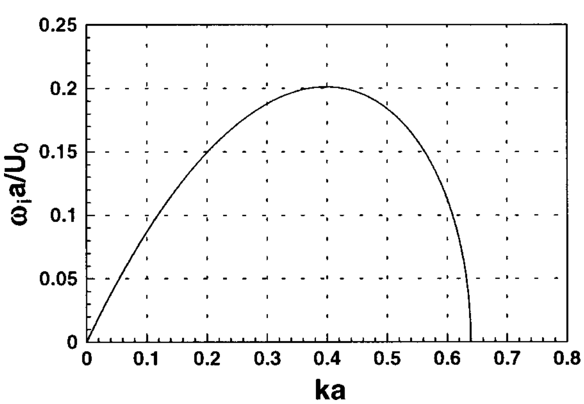

FIG. 2. The normalized growth rate $\omega_i a/U_0$ versus the normalized wave number ka for the velocity profile $V_0(x)$ shown in Fig. 1.

$V_0(x) + U_0$. Therefore, the real frequency ω_r for the velocity profile (32) is obtained by Doppler shifting ω_r obtained for the velocity profile (34) by kU_0. This gives $\omega_r = kU_0$ and the same ω_i as shown in Fig. 2 for $ka < 0.64$ for the unperturbed velocity profile (32). Therefore, we could show that the B_0 profile shown in Fig. 1 is indeed subject to the new KH instability and could obtain the explicit growth rate ω_i plotted in Fig. 2 and $\omega_r = kU_0$ for the profiles of V_0 and B_0 shown in Fig. 1. This is sufficient evidence for the existence of the new instability, which is driven by the shear in the ion diamagnetic drift velocity.

VI. DISCUSSION AND SUMMARY

In the present paper a nonideal MHD KH instability is derived in the frame of the nonideal MHD model. In the frame of the fluid model, any velocity shear would provide a KH type instability. Therefore, as Eq. (27) shows the ion diamagnetic drift and **E**×**B** drift for a fluid contribute the same way to the velocity shear and thus to the KH instability. However, the physical nature of the present new KH instability driven by the shear in the ion diamagnetic drift velocity is important and its difference from the nature of the ideal MHD KH instability driven by the shear in the **E**×**B** drift velocity should be emphasized. Since the ion diamagnetic drift is not entirely a motion of guiding centers, the present nonideal MHD instability driven by the shear in the ion diamagnetic drift velocity is peculiar to the nonuniform high-β plasma. In order to make the present calculation more relevant to collisionless plasmas one would need to carry out the analysis kinetically by solving the ion gyrokinetic equation. A kinetic analysis would capture the special nature of the diamagnetic flow as a finite Larmor radius (FLR) effect. A kinetic treatment of the divergence of the pressure tensor and calculation of the pressure tensor using the solution of the linearized gyrokinetic equation would introduce new diamagnetic terms which do not appear in the present fluid treatment.

The present results obtained for a slab geometry could be applicable to the plasma configuration in the equatorial plane in the near-Earth plasma sheet, where the magnetic field is transverse to the plane and the ion diamagnetic drift dominates the **E**×**B** drift. For such a near-Earth plasma, the present calculation indicates that as long as d^3B_0/dx^3 does not change sign in the region where there is a velocity shear and the density is considered to be uniform, the region is stable. Therefore, in the quiescent state of the magnetosphere, the near-Earth plasma would not be subject to the present instability. However, the field lines in the near-Earth plasma sheet are curved unfavorably and in such a configuration, pressure driven instabilities such as ballooning modes are excited when the plasma β exceeds the critical value.[23–37] Therefore, when the magnetospheric substorm is initiated possibly by the onset of such pressure driven instabilities,[23–37] the pressure profile in the near-Earth plasma sheet would be highly disturbed by ballooning modes in the near-Earth plasma sheet or in the outer edge of the ring current and the instability criterion for the present instability may be satisfied in the near-Earth plasma. In such a situation the present instability would be responsible for the vortex formation in breakup auroral arcs, which is observed and known as ''the westward traveling surge'' of discrete auroral arcs. It follows that a more accurate treatment of the present instability in the near-Earth plasma sheet should take into account an unfavorable field-line curvature, which also allows the development of pressure driven instabilities. Here, we should note that the ion diamagnetic drift effect in ballooning instability causes a net drift of ballooning modes in the direction of the ion diamagnetic drift and it also has a stabilizing influence on ballooning modes.[23] Since the ion pressure gradient in the near-Earth plasma sheet is predominantly earthward and the ion diamagnetic drift direction is westward, this may explain the observed predominant westward motion of ''the westward traveling surge.'' Although a ballooning stability analysis of the plasma sheet,[23] which used a local approximation in the radial direction and neglected a shear in the ion diamagnetic drift velocity, showed that the growth rate of ballooning modes decreases with increasing k (east–west wave number), it showed that there is a substantial range of k in which ballooning modes are unstable.[23] Therefore, it is highly conceivable that the development of ballooning modes in the near-Earth plasma sheet may set a favorable pressure profile for the development of the ion diamagnetic drift driven KH instability in the near-Earth plasma sheet.

It is known that even the most quiescent operating regimes (''*H*-modes'') of fusion machines display spontaneous electromagnetic disturbance.[38] Considering the nonuniform nature of pressure distribution in high-β fusion plasmas, the present nonideal MHD KH instability driven by the shear in the ion diamagnetic drift velocity seems to be ubiquitous in fusion plasmas. However, in order to show that the present instability is also relevant to turbulent transport in tokamaks, it is at least necessary to prove that the instability could survive in the presence of magnetic shear.

In summary, a nonideal MHD KH instability, which is driven by the shear in the ion diamagnetic drift velocity, is found for a high-β plasma. This instability may be relevant to turbulence generation in a high-β fusion plasma, which may affect magnetic confinement, and may also be important in the stability of a high-β plasma and the creation of a certain type of vortex structure in space and astrophysical plasmas.

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