A POSSIBLE ROUTE TO SPONTANEOUS REDUCTION OF THE HEAT CONDUCTIVITY BY A TEMPERATURE GRADIENT-DRIVEN INSTABILITY IN ELECTRON-ION PLASMAS

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ABSTRACT

We have shown that there exist low-frequency growing modes driven by a global temperature gradient in electron and ion plasmas by linear perturbation analysis within the framework of plasma kinetic theory. The driving force of the instability is the local deviation of the distribution function from the Maxwell-Boltzmann distribution due to a global temperature gradient. Application of the results to the intracluster medium is being reduced to 5–7 orders of magnitude less than the mean free paths due to Coulomb collisions. This may provide a hint in explaining why hot and cool gas can coexist in the intracluster medium in spite of the very short evaporation timescale due to thermal conduction if the conductivity is the classical Spitzer value. Our results suggest that the realization of the global thermal equilibrium is postponed by the local instability, which is induced for a quicker realization of the local thermal equilibrium state in plasmas. The instability provides a new possibility to create and grow cosmic magnetic fields without any seed magnetic field.

Subject headings: conduction — cooling flows — galaxies: clusters: general — instabilities — intergalactic medium — plasmas

1. INTRODUCTION

One of the most important results obtained by the Japanese/US X-ray satellite ASCA (Tanaka, Inoue, & Holt 1994) is the first spectroscopical confirmation of the existence of at least two different temperature X-ray–emitting hot gasses in the cluster central region (Fukazawa et al. 1994; Ikebe 1995; Ikebe et al. 1999); the cooler is about 1 keV and the hotter is about 2–10 keV. Although the interpretation of the result is still controversial (Fabian 1994; Ikebe et al. 1999), there is a common feature among the different interpretations: a cool region with temperature $T_1$ is surrounded by a hot region with temperature $T_2$ maintaining a pressure equilibrium (see § 6.2). Since the total thermal energy in the hot region is much higher than that in the cool region, the hot region can be treated as a heat reservoir. Suppose that the interface separating two plasma regions with different temperatures has a thickness of $L$ and temperature variation in the interface region is almost linear in spatial scale. Assume that the electron mean free path $\lambda_e$ is smaller than $L$. The timescale $\tau_{ev}$ of the evaporation due to heat conduction can be roughly estimated by a timescale for electrons in the hotter region to diffuse across the interface, and is given as

$$\tau_{ev} \sim (\lambda_e/v_{th,e}) \times (L/\lambda_e)^2 \sim \epsilon^{-1}(4L/v_{th,e}) \sim 10^4 (L/4 \text{kpc})(k_B T_e/4 \text{keV})^{-0.5}(\epsilon/1.0)^{-1} \text{yr},$$

where $v_{th,e} \equiv (2k_B T_e/m_e)^{1/2}$ is the electron thermal velocity, $T_e$ is the electron temperature, and $m_e$ is the electron mass; $\epsilon$ is defined by $\epsilon \equiv \lambda_e/L$. In the cluster central region, the electron mean free path due to Coulomb collision is obtained as $\lambda_e \sim 0.5(k_B T_e/4 \text{keV})^2(n_e/10^{-2} \text{cm}^{-3})^{-1} \text{kpc}$ (Sarazin 1988). Therefore, the timescale of the evaporation is about $10^6$ yr and is much shorter than the age of cluster of galaxies, which is about the Hubble time. For maintaining the interface at least as long as the age of the cluster, the requirement is to reduce the electron mean free path by 4 orders of magnitude from the Coulomb mean free path to realize $\epsilon \sim 10^{-3}$ is required. Soker & Sarazin (1990) and Makishima (1997) proposed that the cool and hot phases are thermally insulated from each other by magnetic fields. However, their models involve another problem: how such special magnetic field structures are obtained in the cluster central region. In this paper, we show that there is an alternative possibility to reduce the electron mean free path by many orders of magnitude without any special assumption on magnetic field structure.

Ramani & Laval (1978) found a new plasma instability that may relate to the reduction of the electron mean free path, henceforth the reduction of the heat conduction. They showed that the temperature gradient leads to an anisotropic electron-velocity distribution function and that the anisotropy of the velocity distribution function drives the instability like the Weibel instability (Weibel 1959; Fried 1959; Melrose 1986). They proposed that the chaotic magnetic field and electric field, produced by the plasma wave induced by the instability, scatter electrons so that the electron mean free path is reduced by many orders of magnitude. However, the instability found by them cannot be applied to reduce the electron mean free path by many orders of magnitude for the following reasons. Since they assumed that only electrons respond to the mode and ions were treated as fixed background particles, the phase velocity of the plasma wave must be faster than the ion sound velocity $v_{th,i} \equiv (2k_B T_i/m_i)^{1/2}$, where $T_i$, $n_i$, and $m_i$ are the ion temperature, ion number density, and ion mass, respectively. In astrophysical plasmas, the ion mass can be safely replaced by the proton mass. The unstable mode found by Ramani & Laval (1978) has nonzero real part of the wave frequency with a phase velocity of $\sim \epsilon v_{th,e}$. Therefore, the application limit of their analysis sets a relatively high lower limit on $\epsilon$ as $\epsilon > (m_e/m_i)^{1/2} \sim 0.025$ when $T_i = T_e$. (Hereafter, we refer to this lower limit as the wall of the square root of the mass ratio.) To explain the two-phase nature of the hot gas in the cluster central region, the instability found by them is practically not useful and the mechanism that can break the wall of the square root of the mass ratio is required.
The first application of the Ramani & Laval (1978) type instability to astrophysical plasmas was made by Levinson & Eichler (1992). They extend Ramani's & Laval's (1978) analysis to include the nonzero background magnetic field and found the low-frequency unstable mode, similar to the Ramani & Laval (1978) type modes. Pistnner, Levinson, & Eichler (1996) first applied Levinson's & Eichler's (1992) results to the cluster cooling flow. However, they applied their results to reduce the electron mean free path down to $10^{-3}$ and far less than the lower limit, bounded by the application limit of the wall of the square root of the mass ratio. Further, they neglected the cluster gravitational force in their Boltzmann equation in spite of taking into account a pressure gradient with a scale height comparable to the cluster gravitational scale height. As a result, they implicitly assumed the existence of the steady electric field for the electric force acting on an electron to balance with the pressure gradient force acting on the electron. Since the pressure gradient force acting on an ion balances with the cluster gravitational force acting on the ion, the electric force acting on an ion is as strong as the cluster gravitational force acting on the ion. The existence of such a strong electric field in a cluster of galaxies is unlikely. Further, the model adopted as the temperature distribution in the cluster central region obtained by White & Sarazin (1987) only describes the temperature distribution of the hottest phase gas. Therefore, the model adopted by them is inappropriate to answer why the multiphase gas clumps can coexist in the cluster of galaxies.

In this paper, we extend Ramani's & Laval's (1978) analysis including the response of the ion to examine whether we can break the wall of the square root of the mass ratio. We limit our attention to pressure equilibrium plasmas; in other words, the scale of the interface is much smaller than the pressure scale height. In this case, the gravitational force can be neglected. The plan of the paper is as follows. Readers who are not interested in the details of the derivation of the dispersion relation can skip §§ 3 and 4. In § 2, we set the model and summarize the basic assumptions and the application limit coming from the assumptions. In § 3, we deduce the solution of the Boltzmann equation under the situations set in § 2 using the Chapman-Enskog expansion of the electron and ion velocity distribution functions. We deduce the dispersion relations under low-frequency conditions in § 4. In § 5, we summarize the characteristics of the instability and modes found in § 4. The application of the results to the hot plasma in the central region of cluster of galaxies is discussed in § 6. Finally, § 7 summarizes other important implications from our results and remaining problems.

2. THE MODEL AND APPLICATION LIMIT

Consider the situation described in Figure 1. Two plasma regions with different temperatures contact, maintaining pressure equilibrium. The electron temperature in the transition region with a depth of $L$ varies almost linearly from $T_1$ to $T_2$ ($T_1 < T_2$) along the $x$-direction: $\delta_T = (T_2 - T_1)/L$. In clusters of galaxies, the pressure scale height is about the core radius of the cluster mass distribution, say $\sim 50$ kpc (e.g., Hattori, Kneib, & Makino 1999). Therefore, if the size of the interface is smaller than $\sim 10$ kpc, the pressure equilibrium condition for the interface region can be safely applied. The ion temperature is assumed to be the same as electron temperature, $T_i = T_e$. The mean free paths $\lambda_e$ and $\lambda_i$ of electron and ion, respectively, are assumed to be defined by the scattering due to plasma waves. If the scattering of the particle due to electric

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\caption{Sketch of one-dimensional distribution of temperature $T(x)$ in the interface region. $T_1$ and $T_2$ are the temperatures in the cool and hot regions, respectively. Also shown are the number density distribution $n(x)$ (dashed line) and the distribution of pressure (dotted line), which is constant along the $x$-direction.}
\end{figure}
field $E$ of the wave is the dominant source, the mean free path of each particle could be described as follows. Assume that the electric field associated with the plasma wave is coherent on the scale of order of the wavelength $\lambda$ and is chaotic over the scale. Then the momentum change of the particles can be described by the random walk manner in the momentum space. The mean free path is defined as the length of scale that the particles gain or lose their momentum compared to their original momentum during their move of scale. The momentum change of electrons during one step is $m_e \Delta v_e \sim eE\lambda/v_{th,e}$. The number of steps required for electrons to gain or lose their momentum compared to their original momentum is $[m_e v_{th,e}/(eE\lambda/v_{th,e})]^2 \sim (2k_B T_e/eE)^2$. Then we obtain

$$\lambda_e \sim \left(\frac{2k_B T_e}{eE\lambda}\right)^2 \lambda \sim \left(\frac{2k_B T_e}{eE}\right)^2 \frac{1}{\lambda}.$$  

Similarly, the mean free path for ions can be obtained as

$$\lambda_i \sim \left(\frac{2k_B T_i}{eE\lambda}\right)^2 \lambda \sim \left(\frac{2k_B T_i}{eE}\right)^2 \frac{1}{\lambda}.$$  

Since we are now assuming that electrons and ions have the same temperature, the electron and ion mean free paths due to the scattering by the wave electric field are identical. When the chaotic magnetic field associated with the plasma wave is the dominant source of the particle scattering, the particle mean free path would be defined in the following way. The magnetic moment of the particles, $\mu = (1/2m, v^2)/B$, is a conserved variable in a magnetized plasma, where $v_i$ is the particle velocity perpendicular to the magnetic field, written as $m_i v^2 \sim 2k_B T_i$, just after the perturbed magnetic field is induced. As the magnetic field strength gets stronger, $v_i$ increases to conserve $\mu$. The maximum allowable value of $v_i$ is $(3k_B T_i/m_i)^{1/2}$, since the kinetic energy of the particle must be conserved. Then it follows that an electron at the place with a magnetic field strength of $B_1$ is scattered back at the place with a magnetic field strength of $B_2 \sim 3/2B_1$ owing to the mirror effect. The mean free path is then defined as the length scale along the magnetic field from the place where the magnetic field strength is $B_1$ to the place where the magnetic field strength reaches $B_2$; since the critical value $B_{cr}$ of the magnetic field strength, in other words the place of the bouncing, does not depend on the particle species and only depends on their initial position and the magnetic field configuration. Therefore, the mean free paths of electrons and ions must also be the same in this case. These arguments show that the mean free paths of electrons and ions due to scattering by plasma waves may have the same orders of magnitude. Therefore, we assume that electrons and ions have a similar value of $\epsilon$.

In this paper, perturbations with a short-wavelength wave, such that

$$\lambda \ll \lambda_e,$$  

are treated. Therefore, the collision term in the Boltzmann equation for the perturbed variables can be neglected.

A self-consistency check of the obtained results with all these conditions constrains the allowable range of $\epsilon$. In addition, two further conditions provide important limits on $\epsilon$. First of all, the electron mean free path due to Coulomb collision must be shorter than $L$, at least before the onset of the instability, otherwise there is no chance for the electron distribution function to have anisotropy as described in the next section and, hence, no chance to have instability. This condition yields a lower limit on $L$ as

$$L > 0.5(k_B T_e/4 \text{ keV})^2 (n_e/10^{-2} \text{ cm}^{-3})^{-1} \text{ kpc}.$$  

Second, the growth timescale of the unstable mode, $\omega_i^{-1}$, must be shorter than the timescale of thermal diffusion when waves do not exist, in order for the waves to grow before the temperature difference will be erased owing to quick thermal conduction. This condition provides

$$\omega_i^{-1} < \epsilon^{-1} L/v_{th,e},$$  

where only the case that the electron mean free path is shorter than $L$ is considered. A further condition has to be required when applied to the clusters; that is, the growth timescale must be shorter than the age of the object. Conservatively the age of the universe is

$$\omega_i^{-1} < 1/H_0 = 10^{10} \text{ h}^{-1} \text{ yr},$$  

where $H_0 = 100 h \text{ km s}^{-1} \text{ Mpc}^{-1}$ is the Hubble constant.

3. DISTRIBUTION FUNCTION

When $\epsilon < 1$, the discussion given in the previous section ensures that the solution of the Boltzmann equation of electrons and ions can be expanded in powers of $\epsilon$ (Chapman & Cowling 1960),

$$f_e = f_{m,e} + f_e^{(1)} + f_e^{(2)} + \cdots,$$

$$f_i = f_{m,i} + f_i^{(1)} + f_i^{(2)} + \cdots,$$

where $f_{m,e}$ and $f_{m,i}$ are Maxwellian distribution functions for electrons and ions, respectively; $f_{e}^{(k)}$ and $f_{i}^{(k)}$ ($k = 1, 2, \cdots$) describe the deviation of distribution functions from the Maxwellian in order of $\epsilon^k$ for electrons and ions, respectively. Since we assume that $T_i(x) = T_e(x) \equiv T(x)$ and the charge neutrality must be maintained, the Maxwellian parts for both electrons
and ions are written as \( f_{m,e} = n_0(x)[\pi v_{th,e}(x)]^{-3/2} \exp \{ -[v/v_{th,e}(x)]^2 \} \) and \( f_{m,i} = n_0(x)[\pi v_{th,i}(x)]^{-3/2} \exp \{ -[v/v_{th,i}(x)]^2 \} \), respectively. Here \( n_0(x) \) is the electron number density, \( v_{th,e}(x) \equiv [2k_B T(x)/m_e]^{1/2} \), and \( v_{th,i}(x) \equiv [2k_B T(x)/m_i]^{1/2} \).

The pressure equilibrium assumption gives

\[
\frac{\nabla n_0}{n_0} = -\frac{\nabla T}{T} = -\frac{1}{L} \delta_T.
\]

(10)

Once this condition is satisfied, the pressure is time independent everywhere, even if a secular variation of the temperature due to thermal conduction is taken into account. Then the time dependence of the density can be related to that of the temperature as

\[
\frac{1}{n_0} \frac{\partial n_0}{\partial t} = -\frac{1}{T} \frac{\partial T}{\partial t}.
\]

(11)

The Boltzmann equation under the pressure equilibrium condition without background electric and magnetic fields leads to

\[
\frac{\partial f_e}{\partial t} + v_x \frac{\partial f_e}{\partial x} = \left( \frac{\partial f_e}{\partial t} \right)_c,
\]

(12)

where \( \left( \frac{\partial f}{\partial t} \right)_c \) is the collision term and the equilibrium electric field does not appear because of the pressure equilibrium condition. The secular time variation of the distribution function is ascribed to the secular variation of the temperature due to thermal conduction. Since the cool region with a temperature of \( T_1 \) is considered to be immersed in a heat bath with a temperature of \( T_2 \), the temperature of the interface increases monotonically owing to the heat conduction. Hence, the time evolution of temperature can be described as

\[
\frac{\partial T}{\partial t} = \epsilon \delta_T \frac{v_{th,e}}{L} T.
\]

(13)

Then,

\[
\left| \frac{\partial f_e}{\partial t} \right| \sim \frac{1}{T} \left| \frac{\partial T}{\partial t} \right| f_e \sim \epsilon \delta_T \frac{v_{th,e}}{L} f_e.
\]

(14)

The simplest choice of the collision term is the Krook operator, such that

\[
\left( \frac{\partial f_e}{\partial t} \right)_c = -v_c (f_e - f_{m,e}),
\]

(15)

where \( v_c = v_{th,e}/\lambda_e \) is the collision frequency assumed to be constant. Substituting the collision term by the Krook operator, ordering the Boltzmann equation in \( \epsilon^1 \) and \( \epsilon^2 \) provides

\[
f^{(1)}_e = -\frac{1}{v_c} v_x \frac{\partial f_{m,e}}{\partial x},
\]

(16)

\[
f^{(2)}_e = -\frac{1}{v_c} \left( \frac{\partial f_{m,e}}{\partial t} + v_x \frac{\partial f^{(1)}_e}{\partial x} \right).
\]

(17)

It shows that the system can be treated as steady state in the first order of \( \epsilon \), but the secular variation cannot be negligible in the second order of \( \epsilon \). The time derivative of the Maxwell-Boltzmann distribution function can be described as

\[
\frac{\partial f_{m,e}}{\partial t} = -\frac{1}{2} \left( 5 - \frac{2v^2}{v^2_{th,e}} \right) \frac{1}{T} \frac{\partial T}{\partial t} f_{m,e} = -\epsilon \delta_T \frac{v_{th,e}}{L} \left( 5 - \frac{2v^2}{v^2_{th,e}} \right) f_{m,e}.
\]

(18)

The first- and the second-order solutions for the electron distribution function are thus obtained as follows:

\[
f^{(1)}_e = -\epsilon \delta_T \frac{v_x}{v_{th,e}} \left( 5 - \frac{m_e v^2}{k_B T} \right) f_{m,e},
\]

(19)

\[
f^{(2)}_e = \frac{\epsilon^2 \delta_T}{2} \left[ \left( 5 - \frac{2v^2}{v^2_{th,e}} \right) + \frac{v^2_x}{v^2_{th,e}} \left( 5 - \frac{4v^2}{v^2_{th,e}} \right) + \frac{v^2_x}{2v^2_{th,e}} \left( 5 - \frac{2v^2}{v^2_{th,e}} \right)^2 \right] f_{m,e},
\]

(20)

where we have neglected the second derivative of \( T(x) \) with respect to \( x \) in derivation of the second-order solution.

These equations show that the anisotropy in velocity distribution functions are induced by the existence of the temperature gradient. The first-order solution is an odd function of the velocity. On the other hand, the second-order solution is even in velocity. It is straightforward to make sure that the zero-background electric field assumption is consistent with the zero-electric current condition under the pressure equilibrium plasma condition, by checking \( \langle f_e v \rangle = 0 \), where \( \langle \rangle \) denotes the average over the velocity.
By substituting $m_i$ to $m$ and $v_i$ to $v_i = v_i/\lambda_i$ in equations (19) and (20), the first- and the second-order distribution functions $f^{(1)}$ and $f^{(2)}$ for ions are obtained.

4. DISPERSION RELATION

In the short-wavelength limit, the collisionless Boltzmann equation is a good representative for the evolution of the perturbed distribution functions, and the plain wave exp $[i(k \cdot r - \omega t)]$ is a good representative for the perturbed quantities (e.g., electron distribution function as $\exp [i (k \cdot r - \omega t)]$, electric field as $E = E_k \exp [i(k \cdot r - \omega t)]$, and magnetic field as $B = B_k \exp [i(k \cdot r - \omega t)]$). From the linearized Boltzmann equation in terms of the perturbed quantities of both electrons and ions, the equations for the perturbed distribution functions of electrons and ions are obtained as

$$-i(\omega - k \cdot v) f_{k,e} = \frac{e}{m_e} (E_k + v \times (B_k/c)) \cdot \left( \frac{\partial f_e}{\partial v} \right)$$

$$\tag{21}$$

$$-i(\omega - k \cdot v) f_{k,i} = \frac{e}{m_i} (E_k + v \times (B_k/c)) \cdot \left( \frac{\partial f_i}{\partial v} \right).$$

$$\tag{22}$$

As shown below, the low-frequency condition requires that the charge density be small, implying that the electric current has to be almost transverse. Therefore, only two modes are expected. Without loss of generality, the wavevector $k$ can be taken in the $x$-$y$ plane as $k_x = k \cos \theta$, $k_y = k \sin \theta$, and $k_z = 0$. We expect one mode with the magnetic field along the $z$-direction (mode 1: $B_x = B_y = 0$ and $B_z \neq 0$), and the other one with the magnetic field in the $x$-$y$ plane (mode 2: $B_x, B_y \neq 0$ and $B_z = 0$). We thus denote the component of a vector along the direction of the wavevector by a $\perp$ subscript. The $x$-component of velocity, for example, can be expressed as $v_x = v_x \cos \theta - v_y \sin \theta$ in terms of $v_x$ and $v_y$.

The Faraday’s law in Maxwell’s equations leads to $E_x = 0$ [i.e., $E_x = (E_x, E_y, 0)$] and $kE_x = (\omega/c)B_z$ for mode 1. For mode 2, $\text{div} B = 0$ yields $B_x = 0$ [i.e., $B_x = (0, B_y, 0)$], and the Faraday’s law and $B_z = 0$ lead to $E_z = 0$ and $kE_z = -(\omega/c)B_z$. Then we obtain

$$\begin{align*}
(k \cdot v - \omega) f_{k,e} &= -\frac{ie}{m_e} \left[ E_k + \frac{v \perp}{c} B \right] \frac{\partial f_e}{\partial v_k} + \left( E_x - \frac{v_x}{c} B \right) \frac{\partial f_e}{\partial v_x},

\tag{23}

(k \cdot v - \omega) f_{k,i} &= +\frac{ie}{m_i} \left[ E_k + \frac{v \perp}{c} B \right] \frac{\partial f_i}{\partial v_k} + \left( E_x + \frac{v_x}{c} B \right) \frac{\partial f_i}{\partial v_x},

\tag{24}
\end{align*}$$

for mode 1, and

$$\begin{align*}
(k \cdot v - \omega) f_{k,e} &= -\frac{ie}{m_e} \left[ E_k - \frac{v_x}{c} B \right] \frac{\partial f_e}{\partial v_k} + \left( E_x + \frac{v_x}{c} B \right) \frac{\partial f_e}{\partial v_x},

\tag{25}

(k \cdot v - \omega) f_{k,i} &= +\frac{ie}{m_i} \left[ E_k - \frac{v_x}{c} B \right] \frac{\partial f_i}{\partial v_k} + \left( E_x - \frac{v_x}{c} B \right) \frac{\partial f_i}{\partial v_x},

\tag{26}
\end{align*}$$

for mode 2.

To first order in $\omega/(k v_{th,e})$ and $\omega/(k v_{th,i})$, $(k \cdot v - \omega)^{-1} = (1/k)P(1/v_k) + (\omega/k^2)P(1/v_k^2) + i(\pi/k) \delta(v_k) - i(\pi/k)[\omega/k][d\delta(v_k)/dv_k]$, where $P$ denotes the principal value and the signs in front of the delta functions reflect the causality condition. Then, with the help of the Faraday’s law’s equations (23) and (24) yield

$$\begin{align*}
\hat{f}_{k,e} &= -\frac{ie}{m_e} \left[ \frac{1}{k} P \frac{1}{v_k} + \frac{\omega}{k^2} P \frac{1}{v_k^2} + i \frac{\pi}{k} \delta(v_k) - i \frac{\pi \omega}{k^2} \frac{d\delta(v_k)}{dv_k} \right] \\
&\times \left[ E_k + \frac{v \perp}{c} B \right] \frac{2v_k}{v_{th,e}} + \frac{\epsilon \delta_T}{2v_{th,e}} \cos \theta \left[ 5 - \frac{2v_k^2}{v_{th,e}} \right] - \frac{2v_k}{v_{th,e}} (v_k \cos \theta - v_{th,e} \sin \theta \left[ 7 - \frac{2v_k^2}{v_{th,e}} \right]) \right],

\tag{27}

\hat{f}_{k,i} &= +\frac{ie}{m_i} \left[ \frac{1}{k} P \frac{1}{v_k} + \frac{\omega}{k^2} P \frac{1}{v_k^2} + i \frac{\pi}{k} \delta(v_k) - i \frac{\pi \omega}{k^2} \frac{d\delta(v_k)}{dv_k} \right] \\
&\times \left[ E_k + \frac{v \perp}{c} B \right] \frac{2v_k}{v_{th,i}} + \frac{\epsilon \delta_T}{2v_{th,i}} \cos \theta \left[ 5 - \frac{2v_k^2}{v_{th,i}} \right] - \frac{2v_k}{v_{th,i}} (v_k \cos \theta - v_{th,i} \sin \theta \left[ 7 - \frac{2v_k^2}{v_{th,i}} \right]) \right],

\tag{28}
\end{align*}$$
for mode 1. Similarly, from equations (25) and (26), we obtain

\[ f_{k,e} = \frac{-i e}{m_e} \left( \frac{1}{k} P \frac{1}{v_k} + \frac{\omega}{k^2} \frac{1}{v_e^2} + \frac{i \pi}{k^2} \delta(v_{k}) - \frac{i \pi \omega}{k^2} d\delta(v_{k}) \right) \]

\[ \times \left( -\frac{v_z}{c} B_z \right) \left\{ -\frac{2v_k}{v_{th,e}} + \frac{e \delta_T}{2v_{th,e}} \left[ \cos \theta \left( 5 - 2v_e^2 \right) - \frac{2v_k}{v_{th,e}} (v_k \cos \theta - v_z \sin \theta \left( 7 - \frac{2v_e^2}{v_{th,e}} \right) ) \right] \right\} \]

\[ + \frac{B_z}{ck} \left\{ -\frac{2v_k}{v_{th,e}} + \frac{e \delta_T}{2v_{th,e}} (v_k \cos \theta - v_z \sin \theta \left( 7 - \frac{2v_e^2}{v_{th,e}} \right) ) \right\} f_{m,e} , \]

\[ f_{k,i} = \frac{-i e}{m_i} \left( \frac{1}{k} P \frac{1}{v_k} + \frac{\omega}{k^2} \frac{1}{v_i^2} + \frac{i \pi}{k^2} \delta(v_{k}) - \frac{i \pi \omega}{k^2} d\delta(v_{k}) \right) \]

\[ \times \left( -\frac{v_z}{c} B_z \right) \left\{ -\frac{2v_k}{v_{th,i}} + \frac{e \delta_T}{2v_{th,i}} \left[ \cos \theta \left( 5 - 2v_i^2 \right) - \frac{2v_k}{v_{th,i}} (v_k \cos \theta - v_z \sin \theta \left( 7 - \frac{2v_i^2}{v_{th,i}} \right) ) \right] \right\} \]

\[ + \frac{B_z}{ck} \left\{ -\frac{2v_k}{v_{th,i}} + \frac{e \delta_T}{2v_{th,e}} (v_k \cos \theta - v_z \sin \theta \left( 7 - \frac{2v_i^2}{v_{th,i}} \right) ) \right\} f_{m,i} \]

(29)

(30)

for mode 2.

The Poisson equation is written as

\[ i k E_k = 4 \pi e \left( \langle f_{k,e} \rangle - \langle f_{k,e} \rangle \right) . \]

(31)

Keeping only the first nonvanishing order in \( \epsilon \) for mode 1, each term on the right-hand side provides

\[-4 \pi e \langle f_{k,e} \rangle = -\frac{i}{k} \left( \frac{2}{\sqrt{\pi} B_z} E_k - \frac{e \delta_T}{c v_{th,e}} B_z \sin \theta \right) , \]

(32)

\[ + 4 \pi e \langle f_{k,i} \rangle = -\frac{i}{k} \left( \frac{2}{\sqrt{\pi} B_z} E_k - \frac{e \delta_T}{c v_{th,i}} B_z \sin \theta \right) , \]

(33)

where \( \omega_{pe} \) and \( \omega_{pi} \) are the electron and ion plasma frequencies, defined by \( \omega_{pe} \equiv (4 \pi n_0 e^2/m_e)^{1/2} \) and \( \omega_{pi} \equiv (4 \pi n_0 e^2/m_i)^{1/2} \), respectively; \( \lambda_D \equiv [k_B T/(4 \pi n_0 e^2)]^{1/2} \) is the Debye length. It is trivial why those equations do not contain \( E_z \) since the transverse component is not constrained from the Poisson equation. Note that the first term in the right-hand side for ions is exactly the same as that for electrons. Therefore, the contribution of ions to the charge density is nonnegligible. On the other hand, the second term for ions is smaller than that for electrons by a factor of \( m_i/m_e \). This is an essential point for in-phase acoustic oscillation between electrons and ions to be possible. Then the spatial charge carried by ion acoustic oscillation is able to be canceled by electrons to keep charge neutrality. In the case of pure electron plasmas examined by Ramani & Laval (1978), electrons have to keep charge neutrality by themselves. Therefore, the amplitude of acoustic oscillation for electron plasmas have to be almost zero. The Poisson equation provides the relation between electric and magnetic fields as

\[ E_k = \frac{e \delta_T}{4} \sin \theta \frac{v_{th,e}}{c} B_z . \]

(34)

The term \((k \lambda_D)^2\) in the denominator comes from the left-hand side of the Poisson equation and is always a very small number since the wavelength of the mode that we are interested in is much larger than the Debye length. Therefore, we neglect this term. This corresponds to the so-called plasma approximation (Chen 1974; Tanaka & Nishikawa 1996) and ensures that the charge neutrality is kept in high accuracy even though the longitudinal oscillation exists. When the wavevector is parallel to the temperature gradient (\( \theta = 0 \)), this equation tells \( E_z = 0 \) and mode 1 becomes pure transverse. Compared with the case of pure electron oscillation as in Ramani & Laval (1978), the right-hand side of equation (34) is a factor of 2 smaller, which makes mode 1 more unstable, as shown below. No constraint comes from the Poisson equation for mode 2 since it is a pure transverse mode.

We are now in a position to derive the dispersion relation. Ampère’s law is written as

\[ -i k B_z = \frac{4 \pi e}{c} (\langle v_z f_{k,i} \rangle - \langle v_z f_{k,e} \rangle) + \frac{\partial E_z}{c \partial t} \]

(35)

for mode 1, and

\[ i k B_z = \frac{4 \pi e}{c} (\langle v_z f_{k,i} \rangle - \langle v_z f_{k,e} \rangle) + \frac{\partial E_z}{c \partial t} \]

(36)

for mode 2. Under the low-frequency condition, the last terms in the right-hand side of equations (35) and (36) (the displacement current) are negligibly small. We thus neglect these terms. Furthermore, the contribution of ions to the current density is an order of \((m_i/m_e)\) smaller than that of electrons since the current density carried by each particle is proportional to the thermal velocity of each particle. Therefore, the current density carried by ions introduces only a negligible contribution to
the dispersion relation. The dispersion relation for the real parts $\omega_r$ and the imaginary parts $\omega_i$ in the leading order of $\epsilon$ is then obtained as follows:

$$\omega_r = \frac{\epsilon \delta_T}{4} k v_{th,e} \cos \theta,$$

$$\omega_i = \frac{\epsilon^2 \delta_T^2}{4 \sqrt{\pi}} k v_{th,e} (2 \cos^2 \theta - \sin^2 \theta) - \frac{1}{\sqrt{\pi}} \left( \frac{c}{\omega_{pe}} \right)^2 k^3 v_{th,e},$$

for mode 1, and

$$\omega_r = \frac{\epsilon \delta_T}{4} v_{th,e} k \cos \theta,$$

$$\omega_i = \frac{\epsilon^2 \delta_T^2}{2 \sqrt{\pi}} v_{th,e} k \cos^2 \theta - \frac{1}{\sqrt{\pi}} \left( \frac{c}{\omega_{pe}} \right)^2 v_{th,e} k^3,$$

for mode 2.

Since the imaginary part $\omega_i$ is an order of $\epsilon^2$, we have checked whether the leading order of the dispersion relation is changed when the second-order distribution functions $f(2)$ and $f(3)$ are taken into consideration. We have confirmed that these second-order distribution functions only introduce one-order higher terms in both real and imaginary parts and, thus, main results shown above are not changed.

5. INSTABILITY AND MODE CHARACTERISTICS

The dispersion relations (37), (38), (39), and (40) show that the low-frequency mode, for which the phase velocity of the wave is slower than ion thermal velocity, can exist as long as $\epsilon < (m_i/m_e)^{1/2} \sim 0.025$ is satisfied. This is the first confirmation that the Ramani & Laval (1978) type instability can exist beyond the wall of square root of the mass ratio. The dispersion relations are almost the same as those obtained for electron plasmas by Ramani & Laval (1978). However, there is a slight difference in the imaginary part of mode 1. The difference comes from the difference in the charge neutrality condition in equation (34). As explained in § 4, in our case the in-phase acoustic oscillation between electrons and ions is possible. Hence, the spatial charge carried by ion acoustic oscillation is able to be canceled by electrons to keep charge neutrality. However, in the case of pure electron plasmas examined by Ramani & Laval (1978), electrons have to keep charge neutrality by themselves. Therefore, the amplitude of acoustic oscillation for electron plasmas has to be almost zero.

The characteristics of the instability are summarized as follows. For mode 1, the imaginary part of the wave frequency is positive, and the instability sets in when the direction of the wavevector is within the double cone spanned by $\theta \in (-\theta_c, \theta_c)$ and $\theta \in (\pi - \theta_c, \pi + \theta_c)$, where $\theta_c \equiv \arccos (1/\sqrt{3}) (0 \leq \theta_c \leq \pi/2)$. For comparison, we recalculated the dispersion relations for pure electron plasmas. In the case of pure electron plasmas, a factor of 2 appears in front of the $\sin^2 \theta$ term in the imaginary part (this result is slightly different from Ramani’s & Laval’s 1978 result). Therefore, the unstable region in $k$ space in the case of pure electron plasmas is somewhat narrower than in our case. It shows that acoustic oscillation of ions is assisting the instability. For mode 2, the unstable mode exists for all directions of the wavevector, although the growth rate decreases as $\theta$ increases from 0 to $\pi/2$ for fixed $k$. Two modes become identical when $\theta = 0$. For both modes, the growth rate is at maximum when $\theta = 0$. Since the imaginary part of the wave frequency is a third-order polynomial of $k$ with a negative coefficient for $k^3$, there exists a maximum growth rate $\omega_{i,max}$. For mode 1, $\omega_{i,max}$ is given as

$$\omega_{i,max} \sim \frac{\epsilon^3 \delta_T^3}{12 \sqrt{3 \pi}} v_{th,e} \omega_{pe} |3 \cos^2 \theta - 1|^{3/2},$$

when

$$k = k_{max} = \frac{\epsilon \delta_T \omega_{pe}}{2 \sqrt{3} \ c} \left|3 \cos^2 \theta - 1\right|^{1/2},$$

for $\theta \in (-\theta_c, \theta_c)$ or $\theta \in (\pi - \theta_c, \pi + \theta_c)$. For mode 2,

$$\omega_{i,max} \sim \frac{\epsilon^3 \delta_T^3}{3 \sqrt{6 \pi}} v_{th,e} \omega_{pe} |\cos \theta|,$$

when

$$k = k_{max} = \frac{\epsilon \delta_T \omega_{pe}}{\sqrt{6} \ c} |\cos \theta|,$$

for arbitrary $\theta$. For the wavevector that satisfies the above conditions, the real part $\omega_r$ of the wave frequency is obtained as

$$\omega_r \sim \frac{\epsilon^2 \delta_T^2}{8 \sqrt{3}} v_{th,e} \omega_{pe} |3 \cos^2 \theta - 1|^{1/2} \cos \theta.$$
for mode 1, and

$$\omega_r \sim \frac{1}{4\sqrt{6}} \epsilon^2 \delta_T \frac{v_{th,e}}{c} \omega_{pe} \cos^2 \theta$$  \hspace{1cm} (46)$$

for mode 2.

The electric field strength can be related to the magnetic field strength. For mode 1,

$$|E_k| \sim \frac{\epsilon \delta_T}{4} \frac{v_{th,e}}{c} |B_k| \sin \theta.$$  \hspace{1cm} (47)$$

When $k = k_{max}$, the electric field perpendicular to the $k$-direction is related to the magnetic field as

$$|E_z| \sim \frac{\epsilon \delta_T}{4} \frac{v_{th,e}}{c} |B_k| \cos \theta.$$  \hspace{1cm} (48)$$

For mode 2,

$$|E_z| \sim \frac{\epsilon \delta_T}{4} \frac{v_{th,e}}{c} |B_k| \cos \theta,$$  \hspace{1cm} (49)$$

when $k = k_{max}$.

The electric field strength is then an order of $\epsilon$ smaller than the magnetic field strength. This nature combined with the low-frequency nature of the mode shows that the mode is similar to the magnetohydrodynamical mode. However, the magnetohydrodynamical treatment cannot identify the mode and the instability. For example, the instability is microscopical, in which case the resonance of particles with waves is essential.

The possible reduction factor of the mean free path, in other words the lowest possible value of $\epsilon$, is set by the application limits and assumptions summarized in § 2. In the following discussion, we treat the fastest growing mode, that is the $\theta = 0$ case. First, the condition in equation (6) provides

$$\epsilon > 4 \left( \frac{L}{\omega_{pe}} \right)^{1/2} \left( \frac{\delta_T}{1.0} \right)^{-3/2}.$$  \hspace{1cm} (50)$$

Second, the age condition in equation (7) provides

$$\epsilon > 2 \left( \frac{H_0}{v_{th,e}} \frac{c}{\omega_{pe}} \right)^{1/3} \left( \frac{\delta_T}{1.0} \right)^{-1}.$$  \hspace{1cm} (51)$$

Third, the collisionless condition is satisfied when $2\pi/k_{max} < \epsilon \times L$, which yields

$$\epsilon > 4 \left( \frac{L}{\omega_{pe}} \right)^{1/2} \left( \frac{\delta_T}{1.0} \right)^{-1/2}.$$  \hspace{1cm} (52)$$

Finally, the condition that the lifetime of the interface given by equation (1) must be longer than the age of the cool region, $t_{age(\text{cool})}$, provides the upper limit on $\epsilon$ as

$$\epsilon < 10^{-5} \left[ \frac{t_{age(\text{cool})}}{10^{10} \text{ yr}} \right]^{-1} \left( \frac{L}{4 \text{ kpc}} \right) \left( \frac{k_B T_c}{4 \text{ keV}} \right)^{-0.5}.$$  \hspace{1cm} (53)$$

6. APPLICATION TO THE CLUSTER OF GALAXIES

6.1. A Jam Bun Model

For the several nearby clusters of galaxies, the existence of at least two different temperatures of gas in the central region is spectroscopically confirmed (Ikebe 1995). If the cool gas found in the cluster central region is a single block of low-temperature gas located at the bottom of the cluster gravitational potential, the cool gas is much like jam in a jam bun; the interface between the cool and hot gas can be well described by the pressure equilibrium condition, and thus the results obtained in this paper can be applied. Among these well-studied clusters is the Centaurus cluster (Ikebe et al. 1999). In the following discussion, the cool gas found in the Centaurus cluster is assumed to be well described by the jam bun model. The temperature of the cool region is about $k_B T_c \sim 1 \text{ keV}$ and that of the hot region is about $k_B T_h \sim 4 \text{ keV}$ in the Centaurus cluster (Ikebe et al. 1999). The factor $\delta_T$ is obtained as $\delta_T \sim (T_{\text{hot}} - T_c)/(T_{\text{hot}} + T_c)/2 \sim 1.2$. The size of the region, where the cool component exists, is about the size of the X-ray halo around the elliptical galaxies, say, 50–100 $h^{-1}$ kpc (Forman, Jones, & Tucker 1985; Matsushita et al. 1998; Ikebe et al. 1999). Using the best-fit values for the electron density profile for the hot-phase gas obtained by ASCA results (Ikebe et al. 1999), the electron number density of the hot-phase gas at around 100 kpc can be estimated as $2.5 \times 10^{-3}$ cm$^{-3}$. Then the electron mean free path due to Coulomb collisions, $\lambda_C$, is about 2 kpc. Hence, the condition in equation (5) implies $L > 2$ kpc, at least when the low-temperature region was formed.
The collisionless condition in equation (52) also provides a lower limit on $\epsilon$ as

$$\epsilon > 10^{-7} \left( \frac{n_0}{2.5 \times 10^{-3} \text{ cm}^{-3}} \right)^{1/4} \left( \frac{L}{4 \text{ kpc}} \right)^{1/2} \left( \frac{\delta_T}{1.2} \right)^{-1/2}.$$  

(55)

Therefore, the age condition in equation (51) yields the lower limit on $\epsilon$ when $L > 0.25$ kpc and is compatible with the lifetime condition in equation (53). On the other hand, the condition in equation (55) provides the lower limit on $\epsilon$ when $L < 0.25$ kpc. In this case, for this condition to be compatible with the lifetime condition in equation (53), the lower limit of the size of the interface region is set as

$$L > 0.2 \left[ \frac{t_{\text{age(cool)}}}{10^{10} \text{ yr}} \right]^{2/3} \left( \frac{n_0}{2.5 \times 10^{-3} \text{ cm}^{-3}} \right)^{-1/6} \left( \frac{k_B T_e}{4 \text{ keV}} \right)^{1/3} \left( \frac{\delta_T}{1.2} \right)^{-1/3} \text{ kpc}.$$  

(56)

Compared with the cluster pressure scale height, the obtained small values of the lower limit on $L$ from both the age condition and the condition in equation (56) guarantee the application of the pressure equilibrium condition. Therefore, we conclude that the instability found in this paper can be a possible mechanism to inhibit the heat conduction between cool and hot phases in the cluster central region, if the cool gas is a single block of the low-temperature gas located at the bottom of the cluster gravitational potential well.

### 6.2. The Raisin Bread Model

If the cool gas is contained in small blobs and the blobs distribute in the entire cluster central region like raisins in raisin bread, the obtained results cannot be simply applied. The cooling flow model, which is one of the widely accepted models to describe the cluster central region (Fabian 1994), assumes that the intracluster medium in the cluster central region is in a multiphase and is composed of many small cool clumps (raisins) that have different temperatures from clump to clump. Even in the two-temperature phase model (Ikebe et al. 1999), the raisin bread model is thought to be more plausible to describe the situation in the cluster central region. It is well known that these cool clumps, which are in pressure equilibrium with the surrounding hot medium in the cluster central region, oscillate owing to buoyancy force (Balbus 1986; Malagoli, Rosner, & Bodo 1987; Balbus & Soker 1989) and that they are fragmented (Nulsen 1986; Balbus 1986) and merge into the surrounding medium in the cluster central region. It is well known that these cool clumps, which are in pressure equilibrium with the surrounding hot medium in the cluster central region, oscillate owing to buoyancy force (Balbus 1986; Malagoli, Rosner, & Bodo 1987; Balbus & Soker 1989) and that they are fragmented (Nulsen 1986; Balbus 1986) and merge into the surrounding medium in the cluster central region. It is well known that these cool clumps, which are in pressure equilibrium with the surrounding hot medium in the cluster central region, oscillate owing to buoyancy force (Balbus 1986; Malagoli, Rosner, & Bodo 1987; Balbus & Soker 1989) and that they are fragmented (Nulsen 1986; Balbus 1986) and merge into the surrounding medium in the cluster central region.

$\frac{1}{\beta} > \frac{\delta_T}{R}.$  

(60)
In the equilibrium state, the total pressure in the clump must be the same as the total pressure of the surrounding hot medium. This condition provides

$$|\delta P_g| \approx |\delta_b P_m|,$$  \hspace{1cm} (61)

in the clump. Therefore, the deviation of the gas pressure relative to the surrounding gas pressure can be estimated as

$$\frac{|\delta P_g|}{P_g} \approx \frac{\delta_r r}{R}.$$  \hspace{1cm} (62)

From equations (60) and (62), we obtain the constraint on the relative deviation of the gas pressure:

$$\frac{1}{\beta} > \frac{|\delta P_g|}{P_g}.$$  \hspace{1cm} (63)

Since the hot gas in the cluster is a high $\beta$ plasma, say, $\beta \sim 10-100$ (Kim et al. 1989; Makino 1997), the relative deviation of the gas pressure is at most 10% and very small.

The arguments presented above ensure that the model of the interface adopted in this paper is applicable to the case when the raisin bread model is a good representative for the cluster central region. Therefore, the results obtained in this paper can provide some hint for the reduction of heat conductivity in the cooling flow. However, it is required that further analysis including the background magnetic field applied to the raisin bread model, since the background magnetic field introduces another resonance related to cyclotron frequencies of electrons and ions.

7. DISCUSSION

We have shown that there exist low-frequency growing modes driven by a global temperature gradient in electron and ion plasmas. The instability provides a hint for a reduction mechanism of thermal conductivity in electron-ion plasmas. However, to answer whether the instability really reduces heat conductivity down to the required level, the nonlinear saturation level of the instability and the effect of the background magnetic field must be examined. Although Levinson & Eichler (1992) have studied the effect of the magnetic field on the Ramani & Laval (1978) type instability for pure electron plasmas, there is no study which includes ions with the background magnetic field.

The evolutionary nature of the modes in the interface region can be briefly described as follows. When the low-temperature gas is injected into the intracluster medium, any wave is not excited in the interface region yet. Hence, $\epsilon$ is defined by the mean free path due to Coulomb collisions and is of the order of $1-0.1$, as shown in § 1. At this initial time period, the Ramani & Laval (1978) type instability excites the waves only for electron plasmas, since the phase velocity of excited waves is about the electron thermal velocity and, thus, much faster than the ion thermal velocity. As waves grow, the mean free paths of electrons and ions become shorter and $\epsilon$ gets smaller. When $\epsilon$ goes down to $(m_e/m_i)^{1/2}$, ions start to respond to the waves, and the modes switch to those described in this paper. Since the growth rate of the instability is also proportional to $\epsilon^2$ and very large at the initial stage, the condition in equation (50) is satisfied. We thus do not have to worry about the disappearance of the interface due to the thermal conduction before the waves grow. The waves must be scattered back at the edge of the interface and must be confined in the interface region, because the waves driven by the instability cannot propagate in the uniform temperature region.

Since the instability creates and grows the magnetic field without any seed field, it provides a new possibility for the origin of the cosmic magnetic field. The Biermann battery effect (Biermann 1950) is the mechanism that has been considered as one of the central mechanisms to create the seed magnetic field in our universe (Kulsrud et al. 1997). The Biermann battery effect requires that the direction of the density gradient of the gas is not parallel to the direction of the temperature gradient. However, the instability found in this paper requires only the existence of the temperature gradient. In the astrophysical situation, two different temperature plasmas contacting and keeping a pressure equilibrium can be commonly found compared with the situation in which the Biermann battery effect may act. Therefore, it may be interesting to examine how the instability found in this paper plays a role on the origin of the cosmic magnetic field.

Finally, it is interesting to note that our results suggest that the realization of the global thermal equilibrium is postponed by the local instability, which is induced for quicker realization of local thermal equilibrium state in plasmas. This result is, of course, not trivial. Once the global thermal equilibrium is achieved, the local thermal equilibrium is realized simultaneously. Therefore, letting electrons be free to conduct thermal energy without any instability is the fastest way to achieve thermal equilibrium. However, the plasma prefers restoring the local deviation of the velocity distribution function from the Maxwell-Boltzmann and excites waves by the instability. As a result, the realization of the local thermal equilibrium is also postponed. The plasma be hurried into an error.

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