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TEMPERATURE DRIFT INSTABILITY OF A PLASMA WITH SHEAR

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A linear theory of temperature drift instability is considered. It is shown that introduction of a sufficiently large shear of the force lines simplifies the classification of the instabilities and makes it possible to separate the temperature drift instability having the largest scale. The localization region, the local increment, and the local frequency of this instability are determined. The case $T_i \ll T_e$ is analyzed in detail. The conditions for stabilization by force-line shear are determined and an estimate is presented for the maximum turbulent transport coefficient.

1. INTRODUCTION

WE consider in this paper a linear theory of drift-temperature instability of a plasma situated in a magnetic field with skewed force lines (shear). This instability was discovered for the case of a straight magnetic field by Rudakov and Sagdeev[1] and was subsequently investigated in[2,3]. Its importance to systems with shear of the force lines was indicated earlier.[4]

Allowance for the shear of the force lines makes it possible to classify in natural fashion the instabilities with respect to the region of their localization.[4] If we assume that the greatest danger lies in large-scale instabilities (with a large localization region), and this conclusion is arrived at from simple dimensional estimates for the diffusion coefficient, then this makes it possible to separate immediately the most dangerous instabilities that lead to the largest transport coefficients. We present below a simple analysis which explains why drift-temperature instability is among the most dangerous.

We introduce first certain relations which we shall find useful later on. We consider a plasma situated in a magnetic field with skewed force lines; then the projection of the wave vector on the magnetic field $k_{\|}(r) = (\mathbf{k}\cdot\mathbf{B})/B$ will be a function of the coordinates. The instabilities develop essentially in regions where $k_{\|} \approx 0$; then, for relatively small-scale perturbations $x \ll a$, where x is the localization region and a is the characteristic transverse dimension of the plasma (all the perturbations considered below are just of this type), we can expand $k_{\|}(r)$ about the point r_0, where $k_{\|}(r_0) = 0$, confining ourselves to the first two terms:

$$k_{\|}(r) = k_{\|}(r_0) + k_{\|}'(r_0)x = k_{\|}'(r_0)x, \quad x = r - r_0. \qquad (1.1)$$

We have here the following connection: $k_{\|}' = k_y\theta(r)r^{-1}$, which serves as the definition of the quantity θ, called the shear of the force lines and representing, in order of magnitude, the angle between the magnetic field force lines that are separated by a distance r. An explicit expression for $\theta(r)$ can be obtained by calculating $k_{\|}'$ for different systems. For example, for a round torus $\theta(r) = q'r^2/Rq^2$, where R is the major radius of the torus, $q = rH_{\varphi}(r)/RH_{\omega}(r)$ is the so called margin coefficient. By k_y we denote the projection of the wave vector on the binormal to the force line, $k_y = m/r$, where m is the azimuthal wave number. Relation (1.1) thus connects $k_{\|}$ with the localization region x of the perturbations.

We proceed now to consider instabilities in a field with shear of the force lines. We note that if the electrons and ions are distributed in the perturbations in accordance with Boltzmann's law, then, naturally, no instability can occur. The condition for this is $\omega/k_{\|}v_T \ll 1$, where v_T is the thermal velocity of the corresponding type of charge. Assume that we are considering an instability whose existence is essentially connected with electron unbalance (e.g., drift instability), i.e.,

$$\omega/k_{\|}v_e \gtrsim 1, \qquad (1.2)$$

then, using (1.1) as well as the fact that $\omega < \omega^*$, where ω^* is the drift frequency, we obtain immediately from (1.2) the characteristic region of localization of these oscillations

$$x_e \lesssim (m_e/m_i)^{1/2}\rho_i/\theta. \qquad (1.3)$$

Here $\rho_i = \sqrt{T/M}\,\Omega_i^{-1}$ is the ion Larmor radius. This localization region is typical of instabilities built up by electrons.

This raises the natural question whether oscillations with a localization region larger than (1.3) exist. If we consider oscillations with $x \gg x_e$ ($\omega/k_{\|}v_e \ll 1$), then the electrons have time to acquire in them a Boltzmann distribution, and the only cause of unbalance would be the ions, for which we could still have $\omega/k_{\|}v_i > 1$. This condition yields another possible localization region

$$x_i \lesssim \rho_i/\theta. \tag{1.4}$$

When $x > x_i$ the ions have also time to enter into equilibrium (if there are no special causes that prevent them from doing so, for example captured particles[5,6]), and the oscillations begin to attenuate. Thus, with such a formulation of the problem, the localization (1.4) is the most feasible and it remains to ascertain whether such oscillations exist, i.e., when the electrons have a Boltzmann distribution and the buildup is produced by the ions. We shall show below that the only instability of this type in the potential region is the drift-temperature instability.

In the investigation of the instability we shall employ in what follows a quasiclassical approximation, and since the obtained increments turn out to be of the same order as the oscillation frequency, such an approximation can be regarded as sufficient.

2. FUNDAMENTAL EQUATION

The equation describing the oscillations of an inhomogeneous plasma with shear of the force lines differs from the equation for the case of a straight magnetic field essentially in the fact that now $k_{\|} = k_{\|}(x)$ is a function of the coordinate x in accord with (1.1). Using the results of Mikhaĭlovskiĭ[7], we can write out the following integral equation for the potential, from which we get a dispersion equation for the potential oscillations considered below:

$$\int e^{ik_x x}\varphi(k_x)\Big\{\sum_{j=i,e}\frac{1}{T_j}\Big[1-\frac{\omega^*_{Tj}}{\omega}e^{-z_j}I_0(z_j)y_j^2 + i\sqrt{\pi}e^{-z_j}I_0(z_j)y_jW(y_j) \times\Big\langle 1-\frac{\omega^*_{nj}}{\omega}+\frac{\omega^*_{Tj}}{\omega}\Big(z_j\Big(1-\frac{I_1(z_j)}{I_0(z_j)}\Big)+\frac{1}{2}-y_j^2\Big)\Big\rangle\Big] + \frac{k^2}{T_e 4\pi e^2 n}\Big\}dk_x = 0, \quad j=i,e. \tag{2.1}$$

We have introduced here the notation:

$$\varphi(k_x)=\frac{1}{2\pi}\int\varphi(x)e^{-ik_x x}dx \tag{2.2}$$

are the Fourier components of the potential;

$$\omega_{Tj}^* = \frac{k_y}{m_j\Omega_j}\frac{dT_j}{dr_0}, \quad \omega_{nj}^* = \frac{k_yT_j}{m_j\Omega_j n}\frac{dn}{dr_0}$$

are the drift frequencies in the temperature and density;

$$z_j=\frac{k_\perp^2T_j}{m_j\Omega_j^2},\quad k_\perp^2=k_x^2+k_y^2,\quad k_y^2=\Big(\frac{m}{r_0}\Big)^2,$$

$$y_j=\frac{\omega}{k_{\|}v_j},\quad v_j=\sqrt{\frac{2T_j}{m_j}},\quad k_{\|}=k_y\theta\frac{x}{r_0},\quad x=r-r_0,$$

I_0 and I_1 are modified Bessel functions, and

$$w(y)=e^{-y^2}\Big(1+\frac{2i}{\sqrt{\pi}}\int_0^y e^{t^2}dt\Big)$$

is Kramp's function.

If we substitute $\varphi(k_x)$ as given by (2.2) in (2.1) and integrate with respect to dk_x, we obtain a homogeneous integral equation with respect to $\varphi(x)$; similarly, if we multiply (2.12) by $\exp(-ik'_x x)$ and integrate with respect to dx, then we obtain an integral equation in k-space with respect to $\varphi(k_x)$.

For the short-wave oscillations considered below, with localization region $x \sim \rho_i/\theta \ll a$—the characteristic dimension of the inhomogeneity—we can regard the macroscopic quantities n, T, dn/dx, etc. as constants. In the investigation of Eq. (2.1) we shall use a quasiclassical approximation, i.e., we shall assume that several nodes of the potential $\varphi(x)$ are spanned by the localization region $x = \rho_i/\theta$. We can then choose as a solution of (2.1) a quasiclassical "wave function" in the form

$$\varphi(x)\sim\exp\Big\{i\int^x k_x(x)dx\Big\}, \tag{2.3}$$

where $k_x(x)$ varies little over a distance on the order of $\lambda_x \sim 1/k_x$. We then obtain from (2.1) the dispersion equation

$$1+\frac{T_i}{T_e}-\frac{\omega_T^*}{\omega}e^{-z}I_0y^2+i\sqrt{\pi}e^{-z}I_0yW(y)\Big\{1-\frac{\omega_n^*}{\omega}+\frac{\omega_T^*}{\omega}\Big[z\Big(1-\frac{I_1}{I_0}\Big)+\frac{1}{2}-y^2\Big]\Big\}+i\sqrt{\pi}\Big(\frac{m_e}{m_i}\Big)^{1/2}\Big(\frac{T_i}{T_e}\Big)^{3/2} \times y\Big\{1+\frac{T_e}{T_i}\frac{\omega_n^*}{\omega}-\frac{T_e}{T_i}\frac{\omega_T^*}{2\omega}\Big\}=0, \tag{2.4}$$

which we shall investigate by the local method[7]. Equation (2.4) has been written out for perturbations with a localization region $\lambda_x \sim \rho_i/\theta$, with $y_e \ll 1$, and small terms of the order of y_e^2 have been omitted. Equation (2.4) has been rewritten in terms of the ion drift frequencies $\omega^*_{n_i}$ and $\omega^*_{T_i}$ (we have used the relation $\omega^*_e = -\omega^*_i T_e/T_i$, $y_e = (m_eT_i/m_iT_e)^{1/2}y$) and the quantities y_i, from which the index i has been omitted.

In the investigation of (2.4) it is convenient to use for the frequency and localization region the dimensionless ratios

$$\omega/\omega_T^*,\quad k_{\|}v_i/\omega_T^* = x/2l_0 = x\theta/\rho_i,$$

where $l_0 = \frac{1}{2}(\rho_i/\theta) d\ln T/d\ln r_0$.

3. LOCALIZATION REGION, INCREMENT, AND FREQUENCY OF TEMPERATURE DRIFT INSTABILITY

An important parameter characterizing the temperature drift instability is the region of oscillation localization l, which we can choose, in the quasiclassical analysis, to be the distance from the point $x = r - r_0 = 0$ to the point $x = l$ where the local increment $\gamma(x)$ vanishes, $\gamma(l) = 0$. Indeed, when $x > l$ the increment becomes negative and the oscillations attenuate in this region, i.e., they are essentially concentrated where $\gamma > 0$. Recognizing that $\gamma = \operatorname{Im}\omega = 0$ when $x = l$ and equating the real and imaginary parts of (2.4) separately to zero (we shall henceforth omit the small electronic term $\sim (m_e/m_i)^{1/2}$), and then eliminating ω from these two equations, we obtain the following expression for the localization region as a function of the transverse wave number $k_\perp^2$:

$$\left(\frac{l}{l_0}\right)^2 = -\frac{2}{1+T_i/T_e}\,(e^{-z}I_0(z))^2\,\frac{1-2/\eta+2z(1-I_1/I_0)}{1+T_i/T_e-e^{-z}I_0},$$

$$l_0 = \frac{1}{2}\frac{\rho_i}{\theta}\frac{d\ln T_i}{d\ln r_0}, \tag{3.1}$$

where l_0 is the characteristic region of localization of the temperature drift instability.

As follows from (3.1), the instability in question exists ($l^2 > 0$) if the following condition is satisfied[2]

$$\eta \equiv \frac{d\ln T_i}{d\ln n} > \frac{2}{1+2z(1-I_1/I_0)} \quad \text{for} \quad \eta < 0. \tag{3.2}$$

In the region of relatively large-scale perturbations $x \sim \rho_i/\theta$, the criterion (3.2) leads to the condition $\eta > 2$. The minimum value $\eta \approx 0.95$ is reached when $z \approx 1$.

The quantity $(l/l_0)^2$ as a function of $z = (k_x^2 + k_y^2)\rho_i^2$ is shown in Fig. 1 for different values of the parameter η in the case $T_i = T_e$. When T_i/T_e decreases the localization region decreases, and vice versa.

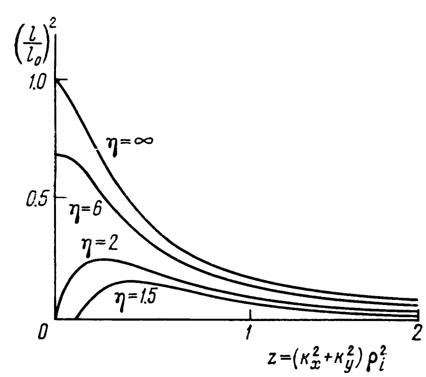

FIG. 1. Dependence of localization region on the wave number.

Knowing the localization region l we can find out how many nodes the solution has at a given value of θ. The approximate number of nodes is

$$n \lesssim k_x l\alpha/\pi \quad (n = 1.2\ldots), \tag{3.3}$$

where n is the nearest integer smaller than $k_x l\alpha/\pi$, and α is a numerical factor of order of unity, which cannot be determined in the local quasiclassical approach. If $k_x l\alpha/\pi < 1$, then not a single node exists. This is precisely the condition for stabilization by the shear of the force lines. The critical value of θ_i can be obtained by substituting in the relation $k_x l\alpha/\pi = 1$ the value of l from (3.1). As a result we obtain θ_i as a function of k_y. The maximum of this quantity with respect to k_y will be denoted by θ_{mc}. An approximate expression for this quantity is

$$\theta_{mc} \approx \frac{0.3}{\pi}\alpha\left(1-\frac{0.95}{\eta}\right)^{1/2} \quad (T_i = T_e). \tag{3.4}$$

If $T_e \ll T_i$, then

$$\theta_{mc} \approx \frac{0.5}{\pi}\alpha\frac{T_e}{T_i}\left(1-\frac{0.95}{\eta}\right)^{1/2}. \tag{3.5}$$

The inverse limiting case, $T_i \ll T_e$, will be considered in the next section.

We now proceed to determine the increments and frequencies of the oscillations. The dispersion equation for the temperature drift instability at $T_e = T_i$ contains no small parameters, and therefore the expression for the local increment $\gamma(x)$ and frequency $\omega(x)$ can be obtained only numerically. Some results of the calculations are shown in Fig. 2 for different values of the parameters

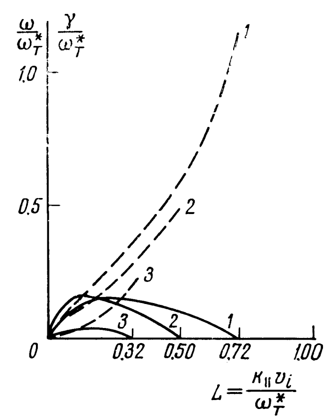

FIG. 2. Local increment γ (solid line) and frequency ω(dashed) for different parameters a and b (3.6): a) a–0.5, b = 1.6; 2) a = 0.5, b = 2; 3) a = 0.2, b = 2.

$$a=\left[\frac{1}{2}-\frac{1}{\eta}+z\left(1-\frac{I_1(z)}{I_0(z)}\right)\right],$$

$$b=\left(1+\frac{T_i}{T_e}\right)e^z I_0^{-1}(z). \quad (3.6)$$

It follows from Fig. 2 that typical values of γ are of the order of 0.1 ω_T^*, and the corresponding ω are of the order of 0.3 ω_T^*. In the region $x \ll 1$ and $|x - l| \ll 1$ we can obtain analytic expressions for γ and ω, but we shall not present here the corresponding cumbersome equations. As follows from the conditions (3.2), the instability in question can develop also at negative values of η ($\eta < 0$). However, if $|\eta| \ll 1$, then the corresponding increment turns out to be exponentially small, $\gamma \sim \exp(-1/|\eta|)$. On the other hand, if $|\eta| > 1$, then this case does not differ in practice from $\eta > 0$. One can advance somewhat farther in the calculations in the important particular case $T_i \ll T_e$.

4. TEMPERATURE DRIFT INSTABILITY WHEN $T_i \ll T_e$

Such a situation is usually realized in Joule heating of a plasma[8], when the entire energy is first taken up by the electrons, and then spreads to the ions via heat exchange, within a time $\tau \sim \nu_{ie}^{-1} \sim \nu_{ee}^{-1} m_i/m_e$.

In the case when $T_i \ll T_e$ the dispersion equation greatly simplifies. Indeed, as will be shown below, the characteristic frequencies of the oscillations in question turn out to be of the order of $\omega \sim k_{||}c_s = k_{||}\sqrt{T_e/m_i}$, and therefore the argument of the ionic W-function turns out to be larger than unity, $\omega/k_{||}v_i \sim \sqrt{T_e/T_i} > 1$, and we can use its asymptotic expansion. Leaving the exponentially small imaginary terms, we obtain from (2.4) the following hydrodynamic dispersion equation (compare with[1]):

$$1+\frac{k_\perp^2 T_e}{m_i\Omega_i^2}\left(1-\frac{\omega^*_{p_i}}{\omega}\right)-\frac{\omega^*_{n_e}}{\omega}-\frac{k_\parallel^2 T_e}{m_i\omega^2}\left(1-\frac{\omega^*_{p_i}}{\omega}\right)=0,$$

$$\omega_{p_i}=\omega_{n_i}+\omega_{T_i}. \quad (4.1)$$

The same equation can be obtained also from the hydrodynamic equations, i.e., the plasma oscillations in question can be described hydrodynamically when $T_i \ll T_e$. We have retained in (4.1) the small term with the transverse ion inertia (the second on the left in (4.1)), which turns out to be of importance in the determination of the stability limit.

We consider first the most unstable case, $\nabla n = 0$. Then the equation simplifies somewhat:

$$1+\frac{k_\perp^2 T_e}{m_i\Omega_i^2}\left(1-\frac{\omega^*_{T_i}}{\omega}\right)-\frac{k_\parallel^2 T_e}{m_i\omega^2}\left(1-\frac{\omega^*_{T_i}}{\omega}\right)=0. \quad (4.2)$$

Let us obtain the condition under which the shear of the force lines stabilizes the oscillations in question. To this end we make in (4.2) the substitution $k_\perp^2 \to -d^2/dx^2 + k_y^2$ and obtain a differential equation for φ. Substituting $k_{||} = k_y\theta x/a$, we get an equation that coincides with that of a quantum oscillator. From this equation it is easy to see that when the condition

$$\theta \gtrsim {}^{27}/_4\, T_i/T_e \quad (4.3)$$

is satisfied there are no localized solutions, i.e., the instability is stabilized.

Far from the stability limit, when the increment becomes of the same order as the frequency, i.e., when the inequality inverse to (4.3) $\theta < T_i/T_e$ is satisfied, we can omit from (4.2) the term with the transverse ion inertia, and obtain the increment and the frequency from the obtained local dispersion equation. The explicit expressions for them are

$$\frac{\gamma}{\omega_{T_i}^*}=\frac{\sqrt{3}}{2}\left(\frac{1}{2}\frac{T_e}{T_i}L^2\right)^{1/3}(R_+-R_-),$$

$$\frac{\omega}{\omega_{T_i}^*}=\frac{1}{2}\left(\frac{1}{2}\frac{T_e}{T_i}L^2\right)^{1/3}(R_++R_-),$$

$$R_\pm=\left(1\pm\sqrt{1-\frac{4T_e}{27T_i}L^2}\right)^{1/3},\quad L=\frac{k_\parallel v_i}{\omega_{T_i}^*}=\frac{x}{2l_0}. \quad (4.4)$$

Plots of these functions for the case $T_e/T_i = 10$ are shown in Fig. 3. Comparing Figs. 2 and 3 we see that the temperature drift instability with $T_i \ll T_e$ considered in this section has a much larger increment (for $T_e \sim 10T_i$) than the instability with $T_e = T_i$, i.e., when the ion temperature is lowered the instability becomes intensified. The reason is as follows: the main mechanism for

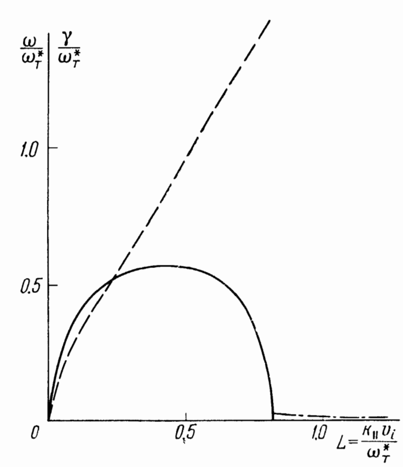

FIG. 3. Local increment γ (solid line) and frequency ω (dashed) in the case $T_i = T_e/10$, $n = 0$.

building up the instability under consideration is hydrodynamic, i.e., the phase relations in the inhomogeneous plasma vary in such a way that if, say, the perturbed temperature rises in a given volume element, then the heat fluxes cause it to increase further (the last term of (4.2)). On the other hand, when $T_e \sim T_i$ the main damping mechanism is the Landau damping by the ions. When T_i decreases the damping drops exponentially, whereas the buildup mechanism depends quite weakly on the ion temperature. This intensifies the instability. Now the localization region begins to be determined only by the pure Boltzmann equalization of the ions along the force lines when $k_{\|}c_S > \omega^*_{T_i}$ (expression (4.5)). With further decrease of T_i, the weakening of the oscillation buildup comes into play, and the oscillations become stabilized.

As follows from (4.4), the characteristic region of localization of the hydrodynamic perturbations is determined by the condition

$$L_0^h = ({}^{27}/_4 T_i/T_e)^{1/2}. \qquad (4.5)$$

We note that the exact kinetic expression (3.1) would have yielded in lieu of (4.5) (in the case when $T_i \ll T_e$, $\eta = \infty$)

$$L_0^k = ({}^1/_2 T_e/T_i)^{1/2}, \qquad (4.6)$$

which is much larger than (4.5) when $T_e \gg T_i$. The difference is attributed to the fact that although the increment outside the region $x < L_0^h$ is not equal to zero (it vanishes when $x = L_0^k$), it is exponentially small. It is shown schematically in Fig. 3 by the dash-dot line to the right of the point L_0^h.

We have considered above the case $\eta = d \ln T_i/d \ln n = \infty$. If $\eta \neq \infty$, then the foregoing analysis is valid when $\eta \gg {}^4/_{27} T_e/T_i$, but if the opposite inequality $\eta \ll {}^4/_{27} T_e/T_i$ is satisfied, then we can leave out of (4.1) the unity in addition to the inertial term, and cancel T_e out. The resultant equation, which does not contain T_e at all, yields the following expression for the oscillation frequencies:

$$\omega/\omega_{T_i}^* = {}^1/_2[\eta L^2 \pm i2\sqrt{\eta}\, L(1 - {}^1/_4 \eta L^2)^{1/2}], \qquad (4.7)$$

where $L = k_{\|}v_i/\omega^*_{T_i}$ and $\eta = d \ln T_i/d \ln n$. As follows from (4.7), when $4T_e/27T_i \gg \eta$, the region of localization ceases to depend on the electron temperature and becomes equal to

$$L_0 = 2/\sqrt{\eta}, \qquad (4.8)$$

and the expression for the characteristic increment takes the form

$$\gamma/\omega_{T_i}^* \sim \sqrt{\eta}\, k_{\|}v_i. \qquad (4.9)$$

Both (4.7) and its corollaries (4.8) and (4.9) are valid when $\eta \gtrsim 4$. In the case just considered $\eta < {}^4/_{27} T_e/T_i$, when the density gradient must be taken into account, it is easy to obtain also the critical shear of the force lines. The oscillations become stabilized when the following inequality is satisfied (compare with (4.3))

$$\theta > 1/\eta \qquad (1 \ll \eta \ll {}^4/_{27} T_e/T_i). \qquad (4.10)$$

Using the results of the foregoing analysis, we can write as an estimate of the turbulent temperature conductivity coefficient[4]:

$$\chi \sim \alpha\rho_i^2 v_i/\theta a, \qquad (4.11)$$

where $\alpha < 1$ is a numerical factor. It is assumed that $\theta > \rho/a$ and $T_i \sim T_e$. From considerations advanced in the introduction, a transport coefficient of the order of $D \sim \rho_i^2 v_i/\theta\, a$ can be regarded as maximal for drift instabilities in a plasma with a sufficiently large shear of the force lines ($\theta \gg (m_e/m_i)^{1/2}$).

5. CONCLUSION

The results of the foregoing analysis can be summarized as follows: In a plasma with shear of the force lines, the maximum region of localization for the drift instabilities does not exceed the order of magnitude of $x_i \sim \rho_i/\theta$; this localization region is attained in the collisionless case only for one instability—the temperature drift instability; the most favorable conditions for its existence are when $T_i \sim T_e$ and $\eta = d \ln T/d \ln n \gg 1$, in which case $\omega \sim \gamma \sim k_{\|}v_i \sim \omega_{T_i}$. If the plasma is not isothermal, then the instability in question becomes stabilized by the shear of the force lines upon satisfaction of a condition that can be written approximately as $\theta \gg T_i T_e/(T_i^2 + T_e^2)$. If the shear of the force lines is not very large, then the decrease of T_i intensifies the instability, and finally, the transport coefficient $D \sim \varphi_i^2 v_i^2/\theta$ a can be regarded as the maximum possible for drift instabilities in plasma with shear.

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