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Research Paper

Collisional Effects on Instabilities in Weakly Ionized Counterstreaming Plasmas

Taghi Mirzaye*[1] · Mohamadreza Taghadosi[2]

[1] Department of Physics Education, Farhangian University, P.O. Box 14665-889, Tehran, Iran;
*E-mail: t.mirzaiee@cfu.ac.ir

[2] Department of Physics, Salman Farsi University of Kazerun, Kazerun, Iran;
Institut für Theoretische Physik, Lehrstuhl IV: Weltraum- und Astrophysik, Ruhr-Universität Bochum, D-44780 Bochum, Germany;
E-mail: mtaghadosi@kazerunsfu.ac.ir

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Abstract. The growth and damping of instabilities in weakly ionized counterstreaming plasmas were investigated using the dispersion relation for electrostatic longitudinal waves. Initially, in the collisionless regime, the frequency spectrum and damping rate were calculated, revealing that energy dissipation is primarily governed by mechanisms such as Landau damping. Subsequently, by including electron–neutral and ion–neutral collision terms, the plasma response and consequently, the growth/damping rates of instabilities were modified. Collisions emerged as effective dissipation mechanisms. A comparison between the collisional and collisionless regimes indicates that the ion–neutral collision term and Landau damping contribute most significantly to instability damping, while the thermal effects of ions play a relatively weaker role in the damping rate. Furthermore, it was determined that an increase in the electron collision frequency reduces the damping rate and enhances the growth of instability.

Keywords: Weakly Ionized Plasma, Counterstreaming Plasma, Plasma Instability, Electron–Neutral Collision, Ion–Neutral Collision, Damping Rate.

1 Introduction

Plasmas play a fundamental role in astrophysical phenomena, ranging from the solar corona to interstellar clouds. Understanding the dynamics of these environments, particularly phenomena related to plasma instabilities, is crucial to interpreting astrophysical observations. Beam-plasma systems, where two or more plasma streams interact with different relative velocities, provide a powerful frameworks for modeling these instabilities. The relative kinetic energy between these streams can excite various waves and instabilities, activating processes such as energy transfer and particle acceleration. A significant portion of plasma characteristics and behaviors is influenced by instabilities; hence, their study is considered a cornerstone of plasma physics [1–4].

In this context, counterstreaming plasmas, a prominent example of two-stream instabilities, are observed in numerous astrophysical environments, including the solar wind, stellar

* *Corresponding author*

winds, and relativistic jets [5–7]. These plasmas also play a significant role in exciting instabilities and even generating magnetic fields in regions such as bow shocks and interplanetary media [6–11]. For instance, non-relativistic electron beams streaming through solar plasma can lead to the generation of various plasma and electromagnetic waves [6,8,10]. Extensive research has been devoted to two-stream instabilities over the past decades, particularly concerning electron beams [12,13] and relativistic beams [5,14,15]. However, comprehensive studies on counterstreaming plasmas containing ion beams [16–18], as well as a detailed investigation of the mutual effects of different species such as electron–neutral and ion–neutral within these systems, especially in weakly ionized regimes, still requires further attention.

On the other hand, understanding how these instabilities grow and damp in the presence of both collisional and collisionless effects is of great importance. Energy dissipation due to collisions can influence the dynamics of instabilities and play a key role in processes such as particle acceleration and plasma heating [4,19]. While in ideal collisionless systems, kinetic instabilities like the Weibel instability and the stream-filament instability dominate [20,21], in realistic environments, especially in weakly ionized plasmas, electron–neutral and ion–neutral collisions are unavoidable. These collisions can act as damping or amplifying mechanisms, ultimately affecting energy transfer between different species and the formation of magnetic fields.

In the present study, these collisional effects are incorporated into the kinetic model to investigate their influence on instability dynamics of counterstreaming plasmas with isotropic Maxwellian velocity distributions assumed for both electrons and ions. This approximation is appropriate for the weakly ionized plasma regime considered here and enables an analytical kinetic treatment based on the plasma dispersion function. The aim is to clarify the impact of these collisions on plasma dynamics and compare it with collisionless regimes.

This paper is organized into four sections. The second section presents the dispersion relation for collisionless counterstreaming plasmas as the baseline. Subsequently, the frequency spectrum and the growth and damping rates associated with the instabilities in these systems are calculated. In the third section, the dispersion relation is derived for counterstreaming plasmas, incorporating the collisional term, followed by the extraction of the excited wave's frequency spectrum and the growth/damping rates of the instabilities. This section also analyzes the role of dissipation from electron and ion collisions and compares it with the collisionless case. Finally, the fourth section provides a summary and conclusion regarding the influence of dissipative and collisional effects on instability dynamics.

2 Formulation of the Problem

2.1 Dispersion Relation and Wave Spectrum in Counterstreaming Collisionless Plasmas

In this study, two counterstreaming collisionless plasmas with equal and opposite velocities (u) are considered. Their drift velocity (u) is assumed to be much smaller than the electron thermal velocity (V_{Te}) and greater than the ion thermal velocity (V_{Ti}). Given that the plasma drift velocity is significantly less than the speed of light, only electrostatic perturbations are investigated. Furthermore, under conditions where the phase velocity of the waves is much lower than the speed of light, the waves in the plasma will exhibit longitudinal characteristics [22]. Accordingly, since the electron thermal velocity is less than the speed of light, the plasma dispersion relation is considered to be longitudinal. By applying the approximations $V_{Ti} \ll u \ll V_{Te}$ and assuming the frequency of fluctuations falls within the range $kV_{Ti} \ll \omega \ll kV_{Te}$, the dispersion relation is obtained, based on the relationships presented in the book by Alexandrov et al. [22], for two counterstreaming plasmas, by substituting $\omega \to \omega \pm k \cdot u$, is as follows:

$$\varepsilon^L(\omega, k) = 0 \Rightarrow 1 + \sum_{\alpha=e,i} \frac{\omega_{p\alpha}^2}{k^2 V_{T\alpha}^2} \left[2 - I_+\left(\frac{\omega - \mathbf{k}\cdot\mathbf{u}}{kV_{T\alpha}}\right) - I_+\left(\frac{\omega + \mathbf{k}\cdot\mathbf{u}}{kV_{T\alpha}}\right)\right] = 0. \quad (1)$$

Here, $\varepsilon^L(\omega, k)$ denotes the longitudinal permittivity of the plasma, where the superscript L indicates the longitudinal component, k is the wavenumber and ω is the angular frequency. Furthermore, $I_+(x)$ denotes the integral plasma dispersion function introduced by Alexandrov *et al.* [22], whose explicit integral representation is

$$I_+(x) = xe^{-x^2/2} \int d\tau\, e^{\tau^2/2}.$$

Here I_+ represents the plasma dispersion integral corresponding to an isotropic Maxwellian velocity distribution. Therefore, the influence of the particle velocity distributions is implicitly contained in the dielectric function through this kinetic response function. Unlike the two-fluid approach of Ref. [23], in which the effects of the particle distributions are represented through the fluid moments of the distribution functions, the present work adopts a kinetic description in which the equilibrium Maxwellian distributions are incorporated through the plasma dispersion integral, I_+.

The asymptotic expansions employed in the present work are those given in Ref. [22]. By expanding the relations and applying the assumptions, the following is obtained as:

$$\begin{aligned} 1 + \frac{2\omega_{pe}^2}{k^2 V_{Te}^2}\left(1 + i\sqrt{\frac{\pi}{2}}\frac{\omega}{kV_{Te}}\right) - \omega_{pi}^2 \left[\left(\frac{1}{(\omega - \mathbf{k}\cdot\mathbf{u})^2}\right) + \left(\frac{1}{(\omega + \mathbf{k}\cdot\mathbf{u})^2}\right)\right] \\ - 3\,\omega_{pi}^2 k^2 V_{Ti}^2 \left[\left(\frac{1}{(\omega - \mathbf{k}\cdot\mathbf{u})^4}\right) + \left(\frac{1}{(\omega + \mathbf{k}\cdot\mathbf{u})^4}\right)\right] \\ + i\sqrt{2\pi}\,\frac{\omega_{pi}^2}{k^3 V_{Ti}^3}(\mathbf{k}\cdot\mathbf{u}\sinh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right) + \omega\cosh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right))\exp\left[-\frac{1}{2}\left(\frac{\omega}{kV_{Ti}}\right)^2\right] = 0, \end{aligned} \quad (2)$$

where ω_{pe} and ω_{pi} are the plasma frequencies for the electron and ion, respectively defined by $\omega_{pe} = \sqrt{\frac{n_e e^2}{\varepsilon_0 m_e}}$ and $\omega_{pi} = \sqrt{\frac{n_i e^2}{\varepsilon_0 m_i}}$ where n_e and n_i are the electron and ion number densities, m_e and m_i denote the electron and ion masses, e is the elementary charge, and ε_0 is the vacuum permittivity. By solving the real part of the dispersion equation $\mathrm{Re}\,\varepsilon^L = 0$ and applying the conditions $k \cdot u \ll \omega$ and also $u < V_s$ (where $V_s = \sqrt{T_e/m_i}$ is the ion-acoustic speed, T_e is the electron temperature, and r_{De} is the electron Debye length defined by $r_{De} = \sqrt{\frac{\varepsilon_0 T_e}{n_e e^2}}$), the two roots of the frequency modes are obtained as follows:

$$\omega_1^2 = \frac{2\omega_{pi}^2}{1 + \frac{1}{k^2 r_{De}^2}} + 5(\mathbf{k}\cdot\mathbf{u})^2 + 3\,k^2 V_{Ti}^2. \quad (3)$$

This mode behaves as follows at two limits:

$$\omega_1^2 = 2\,\omega_{pi}^2 + 5(\mathbf{k}\cdot\mathbf{u})^2 + 3\,k^2 V_{Ti}^2\,; \qquad k^2 r_{De}^2 \gg 1, \quad (4)$$

$$\omega_1^2 = 2\,\omega_{pi}^2 k^2 r_{De}^2 + 5(\mathbf{k}\cdot\mathbf{u})^2 + 3\,k^2 V_{Ti}^2; \quad k^2 r_{De}^2 \ll 1, \quad (5)$$

and the second root:

$$\omega_2^2 = -(\mathbf{k}\cdot\mathbf{u})^2 - 3\,k^2 V_{Ti}^2 < 0. \tag{6}$$

The negativity of ω_2^2 indicates the occurrence of an aperiodic instability. The terms including V_{Ti} represent the role of temperature and thermal effects of ions in the frequency modes that are included in the last term of Equations (3) and (6).

In the mode ω_1^2 and in the short-wavelength limit $k^2 r_{De}^2 \gg 1$, it is observed that the excited mode frequency explicitly depends on the ion plasma frequency and the streaming velocity. The condition $u < V_s$ leads to the excitation of a mode with a frequency approximately $\omega_1^2 \simeq 2\omega_{pi}^2$. In contrast, in the long-wavelength limit $k^2 r_{De}^2 \ll 1$, due to the condition $\omega_{pi}^2 \ll \omega_{pe}^2$, the frequency value will be very small.

Regarding the mode ω_2^2, its negativity indicates a strong interaction between the counterstreaming flows and the occurrence of an aperiodic instability, with the ion thermal term playing an amplifying role in the instability. In this regime, due to their much lower frequency compared to the plasma streaming motion, ions are unable to respond rapidly to fluctuations and thus create a screening effect. Finally, comparing the absolute values, $\mid \omega_2^2 \mid \ll \mid \omega_1^2 \mid$, demonstrates that the contribution of the ω_2 mode to the primary wave excitation process is limited.

2.2 Analysis of Growth and Damping in Collisionless Counterstreaming Plasma

To investigate the evolution of instabilities and analyze the dissipative effects in this system, the frequency is assumed be complex, $\omega = \omega_r + i\delta$, where δ represents the growth rate when positive and the damping rate when negative. The value of δ is calculated using the relation

$$\delta = -\frac{\operatorname{Im}\varepsilon^L(\omega_r, k)}{\frac{\partial}{\partial\omega}\operatorname{Re}\varepsilon^L(\omega,k)\big|_{\omega=\omega_r}}.$$

By substituting the imaginary part of the dispersion relation Equation (2) into this formula, the frequency change rate is derived as follows:

$$\delta = -\sqrt{\frac{\pi}{8}}\,\frac{\omega_{pe}^2}{\omega_{pi}^2}\,\frac{{\omega_r}^4}{k^3 V_{Te}^3} - \sqrt{\frac{\pi}{8}}\,\frac{{\omega_r}^3}{k^3 V_{Ti}^3}\,(\omega_r + \mathbf{k}\cdot\mathbf{u})\exp\left[-\frac{1}{2}\left(\frac{\omega_r}{kV_{Ti}}\right)^2\right]. \tag{7}$$

The value of δ is always negative. This result indicates that the waves are damped in this system, where wave energy is dissipated through mechanisms such as Landau damping for plasma heating. The analysis of Equation (7) shows that in the short-wavelength limit, the contribution of ion thermal effects and counterstreaming velocities (u) to the damping process is significantly greater than in the long-wavelength limit. In other words, at short wavelength scales, direct interaction between ion particles and the plasma flow plays a decisive role in the energy dissipation rate, while at long wavelengths, damping is modulated by electron processes.

3 Oscillatory Spectrum Excited and Dissipative Effects in Counterstreaming Plasmas

In this section, we investigate the effects of collisions. Since the considered state pertains to weakly ionized plasmas, only electron–neutral collisions with frequency ν_{en} and ion–neutral collisions with frequency ν_{in} are taken into account. For counterstreaming plasmas, by substituting $\omega \to \omega \pm \mathbf{k}\cdot\mathbf{u} + i\nu_{\alpha n}$ into the dispersion relation presented by Alexandrov *et al.* [22], the longitudinal dispersion relation is written as follows:

$$\varepsilon^L(\omega,k)=0 \quad \Rightarrow \quad 1+\sum_{e,i}\frac{\omega_{p\alpha}^2}{k^2V_{T\alpha}^2}\left[\frac{1-\ I_+(\frac{\omega+i\nu_{\alpha n}-\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}{1-i\frac{\nu_{\alpha n}}{\omega+i\nu_{\alpha n}-\mathbf{k}\cdot\mathbf{u}}I_+(\frac{\omega+i\nu_{\alpha n}-\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}\right]$$
$$+\sum_{e,i}\frac{\omega_{p\alpha}^2}{k^2V_{T\alpha}^2}\left[\frac{1-\ I_+(\frac{\omega+i\nu_{\alpha n}+\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}{1-i\frac{\nu_{\alpha n}}{\omega+i\nu_{\alpha n}+\mathbf{k}\cdot\mathbf{u}}I_+(\frac{\omega+i\nu_{\alpha n}+\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}\right]=0. \quad (8)$$

Using the conditions $\nu_{en} \gg \nu_{in},\ \ \nu_{in} \ll \omega$ and $\nu_{in} < kV_{Ti}$ the dispersion equation becomes the following as

$$1+\frac{2\omega_{pe}^2}{k^2V_{Te}^2}-\omega_{pi}^2\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^2}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^2}\right)\right]$$
$$-3\omega_{pi}^2\ k^2V_{Ti}^2\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^4}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^4}\right)\right]-\sqrt{\frac{\pi}{2}}\frac{\omega_{pi}^2}{k^3V_{Ti}^3}\nu_{in}\exp\left[-\frac{1}{2}\left(\frac{\omega}{kV_{Ti}}\right)^2\right]$$
$$+i\sqrt{\frac{\pi}{2}}\frac{2\omega_{pe}^2}{k^2V_{Te}^2}\frac{\omega}{kV_{Te}}\frac{1}{1-\sqrt{\frac{\pi}{2}}\frac{\nu_{en}}{kV_{Te}}}+i\ \omega_{pi}^2\nu_{in}\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^3}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^3}\right)\right]$$
$$+i\ 9\ \omega_{pi}^2\ k^2V_{Ti}^2\ \nu_{in}\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^5}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^5}\right)\right]$$
$$+i\sqrt{2\pi}\frac{\omega_{pi}^2}{k^3V_{Ti}^3}(\mathbf{k}\cdot\mathbf{u}\sinh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right)+\omega\cosh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right))\exp\left[-\frac{1}{2}\left(\frac{\omega}{kV_{Ti}}\right)^2\right]=0. \quad (9)$$

Similar to the non-collision plasma state, solving the dispersion relation for one of the excited modes yields:

$$\omega_1^2=\frac{2\omega_{pi}^2}{1+\frac{1}{k^2r_{De}^2}}+5(\mathbf{k}\cdot\mathbf{u})^2+3\ k^2V_{Ti}^2$$
$$+\sqrt{\frac{\pi}{8}}\frac{1}{k^3V_{Ti}^3}\left(\frac{2\omega_{pi}^2}{1+\frac{1}{k^2r_{De}^2}}+5(\mathbf{k}\cdot\mathbf{u})^2\right)^2\nu_{in}\exp\left[-\frac{1}{2}\left(\frac{\frac{2\omega_{pi}^2}{1+\frac{1}{k^2r_{De}^2}}+5(\mathbf{k}\cdot\mathbf{u})^2}{k^2V_{Ti}^2}\right)\right]. \quad (10)$$

In the limit of short wavelengths $k^2r_{De}^2 \gg 1$, the above relationship simplifies to the following:

$$\Rightarrow \omega_1^2=2\omega_{pi}^2+5(\mathbf{k}\cdot\mathbf{u})^2+3\ k^2V_{Ti}^2+\sqrt{2\pi}\frac{\omega_{pi}^4}{k^3V_{Ti}^3}\ \nu_{in}\exp\left[-\frac{\omega_{pi}^2}{k^2V_{Ti}^2}\right]. \quad (11)$$

Equation (11) indicates that the excited mode remains primarily governed by the ion plasma frequency and streaming velocity, while the contribution of ion-neutral collisions appears only as a small correction. Therefore, collisions modify the wave dynamics without substantially altering the real-frequency spectrum.

In the long wavelength limit $k^2r_{De}^2 \ll 1$ is also obtained as follows:

$$\Rightarrow \omega_1^2 = 2\ \omega_{pi}^2 k^2 r_{De}^2 + 5(\mathbf{k}\cdot\mathbf{u})^2 + 3\ k^2 V_{Ti}^2 + \sqrt{\frac{\pi}{8}}\,\frac{\left(2\ \omega_{pi}^2 k^2 r_{De}^2 + 5(\mathbf{k}\cdot\mathbf{u})^2\right)^2}{k^3 V_{Ti}^3}\ \nu_{in}\exp\left[-\frac{1}{2}\frac{\left(2\ \omega_{pi}^2 k^2 r_{De}^2 + 5(\mathbf{k}\cdot\mathbf{u})^2\right)^2}{k^2 V_{Ti}^2}\right]. \quad (12)$$

Considering $\omega_{pi}^2 \ll \omega_{pe}^2$, this mode has a very low frequency and the excited mode's dependence is more on the wavelength. Thus, as the wavelength increases, all terms decrease, and consequently, the frequency of the excited mode decreases. Conversely, decreasing wavelength increases the frequency. In this case as well, the role of the term arising from ion–neutral collisions is less significant than the other terms. Furthermore, in both wavelength limits, considering the applied conditions, the ion thermal term is small compared to the other terms, and its contribution to the frequency of the excited mode is negligible.

The other excited mode in these plasmas will be as follows:

$$\omega_2^2 = -(\mathbf{k}\cdot\mathbf{u})^2 - 3\ k^2 V_{Ti}^2 - \sqrt{\frac{\pi}{8}}\ \frac{(\mathbf{k}\cdot\mathbf{u})^4}{k^3 V_{Ti}^3}\nu_{in}\exp\left[+\frac{1}{2}\left(\frac{\mathbf{k}\cdot\mathbf{u}}{kV_{Ti}}\right)^2\right]. \quad (13)$$

The negativity of all terms in the above equation indicates an aperiodic instability and strong interaction between counterstreaming plasmas. This situation is similar to the results from the previous section when collisions were not considered. The ion thermal effect also appears in the second and third terms. By comparing this relation with Equation (6), it is evident that the third term in this equation corresponds to the contribution of ion–neutral collisions. This equation also shows that the contribution of the collisional term is considerably greater than the effects of counterstreaming plasma velocities and ion thermal velocity in the first and second terms, leading to an increase in the aperiodic instability. Comparing the relations for the plasma frequency modes reveals a dominant linear relationship between frequency and wavenumber in all equations. However, comparing Equations (11) and (5) shows that in the short wavelength limit, there is a deviation from this behavior for the collisional plasma regime.

To investigate the dissipation effects in plasma, it is necessary to calculate the value of δ, which is obtained as follows, considering the imaginary value of the dispersion equation (Equation 12):

$$\delta = -\frac{1}{2}\nu_{in} - \sqrt{\frac{\pi}{8}}\ \frac{\omega_{pe}^2}{\omega_{pi}^2}\frac{{\omega_r}^4}{k^3 V_{Te}^3}\frac{1}{1-\sqrt{\frac{\pi}{2}}\frac{\nu_{en}}{kV_{Te}}} - \sqrt{\frac{\pi}{8}}\ \frac{{\omega_r}^3}{k^3 V_{Ti}^3}\,(\mathbf{k}\cdot\mathbf{u}+\omega_r)\exp\left[-\frac{1}{2}\left(\frac{\omega_r}{kV_{Ti}}.\right)^2\right] \quad (14)$$

From the above equation, it can be concluded that if the electron-neutral collision frequency ν_{en} is less than the electron thermal frequency scale, all terms of δ will be negative. In this case, dissipation and damping occur in the instability process. In contrast, based on Equation (14), if the electron–neutral collision frequency is greater than the electron thermal frequency, δ can become positive, and therefore, instability grows in the plasma. This situation is particularly more likely at long wavelengths. In this regime, the higher the electron–neutral collision frequency, the less the dissipation and damping, and consequently, the growth rate of the instability increases. Each of the δ terms represents the contribution of different dissipation/growth mechanisms. The first term in Equation (14) shows the contribution of ion–neutral collisions to the damping and heating of the plasma. The second term corresponds to the contribution of electron–neutral collisions, which, by comparing it with the electron thermal velocity/frequency scale, can either lead to damping or cause instability growth. The third term also specifies the contribution of counterstreaming plasma velocities and ion thermal velocity to the dissipation of instabilities and heating of counterstreaming plasmas.

4 Summary and Conclusion

The dispersion relation for counterstreaming plasmas in both collisionless and collisional regimes was derived based on the longitudinal dispersion relation, and the frequency spectrum of the instabilities was calculated for both cases. In both regimes, the nature of the dispersion relation roots indicates that one root appears predominantly as a pure imaginary number, expressing an aperiodic instability, while the other root can play a more prominent role in the overall instability behavior.

Table 1: Summary of the principal findings and physical implications of the present study.

Physical aspect	**Collisionless plasma**	**Collisional plasma**	**Main conclusion**
Dispersion spectrum	Two frequency modes are obtained. One corresponds to a propagating mode and the other represents an aperiodic instability.	The same two characteristic modes remain.	Collisions do not significantly modify the overall structure of the dispersion spectrum.
Frequency spectrum	Mainly determined by plasma parameters and streaming velocity.	Slightly affected by collisional terms.	The overall spectral behavior remains nearly unchanged.
Ion thermal effect	Provides only a minor contribution to the frequency and damping.	Remains relatively weak compared with collisional effects.	Ion thermal effects are secondary.
Ion-neutral collisions	Not present.	Introduce additional dissipation.	Ion-neutral collisions constitute one of the dominant damping mechanisms.
Electron-neutral collisions	Not present.	May either reduce damping or promote instability growth depending on the collision frequency.	Increasing electron-neutral collision frequency decreases damping and may enhance instability growth.
Dominant damping mechanism	Landau damping.	Combined effect of Landau damping and ion-neutral collisions.	Dissipation is mainly governed by Landau damping together with ion-neutral collisions.
Instability behavior	Instability evolution controlled by kinetic effects.	Growth/damping strongly depends on collisional frequencies.	Collisions mainly regulate the instability growth rate rather than the wave spectrum itself.

A comparison between the collisionless and collisional regimes demonstrates that collisions have only a minor influence on the real-frequency spectrum, whereas they substantially modify the imaginary part of the dispersion relation and therefore the instability growth and damping rates. The analysis of dissipative effects revealed that the excited waves undergo damping, and the dominant damping mechanism in this framework is a combination of Landau damping and the ion-collision term. Therefore, the contribution of damping from collisional/Landau mechanisms was found to be more significant than that from purely thermal ion effects alone. Consequently, within the parameter regime considered in this work, the damping is primarily governed by Landau damping together with the ion-neutral collisional contribution. Although collisions do not cause significant changes in the overall spectral shape, they can adjust the growth/damping rates through the sign and magnitude of the effective dissipation terms: when the electron–neutral collision frequency is small compared to the electron thermal scale, damping primarily dominates. Conversely, an increase in collisions can, within certain parameter ranges (especially at long wavelengths), modulate the terms and facilitate the amplification and growth of instabilities. In general, increasing the electron collision frequency reduces the damping rate and enhances the instability growth. To facilitate a clear comparison between the collisionless and collisional regimes, the principal findings of the present study are summarized in Table 1.

Appendix

A Derivation of the Longitudinal Dispersion Relation

A.1 Collisionless Regime

In this section, we derive the longitudinal dispersion relation for a collisionless counterstreaming plasma. For vacuum electrodynamics, it is well-established that monochromatic electromagnetic waves of the form $exp(-i\omega t+ik\cdot r)$ can propagate in the absence of external sources. In vacuum, the frequency ω and wave vector k are real and related by $\omega = kc$, where $\mathbf{c}$ is the speed of light. The relation between ω and k is referred to as the dispersion relation. In non-absorbing media, the dispersion relation involves real quantities; however, in absorbing media, these quantities become complex. Since plasmas are generally absorbing media, we consider non-trivial solutions to the field equations in the absence of external sources. Assuming plane monochromatic wave dependence, the field equations are given by:

$$\mathbf{k}\times\mathbf{B} = -\frac{\omega}{c}\varepsilon_{i,j}(\omega,\mathbf{k})E_j, \mathbf{k}\cdot\mathbf{B} = \mathbf{0}, \mathbf{k}\times\mathbf{E} = \frac{\omega}{c}\mathbf{B}, k_i\varepsilon_{i,j}(\omega,\mathbf{k})E_j = \mathbf{0}. \tag{A.1}$$

By eliminating $\mathbf{B}$, we derive a system of three homogeneous equations for the electric field components $\mathbf{E}$:

$$\left[k^2\delta_{i,j} - k_ik_j - \frac{\omega^2}{c^2}\varepsilon_{i,j}(\omega,\mathbf{k})\right]E_j = 0. \tag{A.2}$$

The condition for a non-trivial solution is that the determinant of the operator vanishes. For an isotropic medium, this dispersion equation decomposes into two independent branches:

$$\varepsilon^L(\omega,\mathbf{k}) = 0, \qquad \text{and} \qquad k^2 - \frac{\omega^2}{c^2}\varepsilon^{tr}(\omega,\mathbf{k}) = 0,$$

corresponding to longitudinal and transverse waves, respectively. Given that the thermal velocity of electrons is much less than the speed of light, the longitudinal mode is of primary interest.

The dielectric tensor $\varepsilon_{i,j}(\omega,k)$ is related to the complex conductivity tensor $\sigma_{i,j}(\omega,k)$ as:

$$\varepsilon_{i,j}(\omega,\mathbf{k}) = \delta_{i,j} + \frac{4\pi i}{\omega}\sigma_{i,j}(\omega,\mathbf{k}) = \delta_{i,j} + \sum_\alpha \frac{4\pi e_\alpha^2}{\omega}\int d\mathbf{p}\frac{v_i}{\omega-\mathbf{k}\cdot\mathbf{v}}\frac{\partial f_{0\alpha}}{\partial p_j}. \tag{A.3}$$

Consequently, the longitudinal permittivity $\varepsilon^L(\omega,k)$ is expressed as:

$$\varepsilon^L(\omega,\mathbf{k}) = 1 + \sum_\alpha \frac{4\pi e_\alpha^2}{\omega k^2}\int d\mathbf{p}\frac{(\mathbf{k}\cdot\mathbf{v})}{\omega-\mathbf{k}\cdot\mathbf{v}}(\mathbf{k}\cdot\nabla_\mathbf{p})f_{0\alpha}. \tag{A.4}$$

Assuming a non-relativistic Maxwellian distribution function

$$f_{0\alpha}(p) = \frac{N_\alpha}{(2\pi m_\alpha T_\alpha)^{3/2}}\exp\left(-\frac{p^2}{2m_\alpha T_\alpha}\right).$$

The integral yields the well-known plasma dispersion function $I_+(z)$ [22]:

$$1 + \sum_{\alpha=e,i}\frac{\omega_{p\alpha}^2}{k^2V_{T\alpha}^2}\left[1 - I_+\left(\frac{\omega}{kV_{T\alpha}}\right)\right] = 0. \tag{A.5}$$

For two counterstreaming plasma beams with equal drift velocities $\mathbf{u}$ and $-\mathbf{u}$, the dispersion relation is modified to:

$$1+\sum_{\alpha=e,i}\frac{\omega_{p\alpha}^2}{k^2V_{T\alpha}^2}\left[2-I_+\left(\frac{\omega-\mathbf{k}\cdot\mathbf{u}}{kV_{T\alpha}}\right)-I_+\left(\frac{\omega+\mathbf{k}\cdot\mathbf{u}}{kV_{T\alpha}}\right)\right]=0. \tag{A.6}$$

The function $I_+(x)$ is defined as [22]:

$$I_+(x)=xe^{-x^2/2}\int_{i\infty}^{x}d\tau e^{\tau^2/2}, \tag{A.7}$$

The asymptotic expansions of $I_+(x)$ for different limits are given as follows:

- For $|x|\gg 1$ and $|\mathrm{Re}\{x\}|\gg|\mathrm{Im}\{x\}|$ (with $\mathrm{Im}\{x\}<0$):

$$I_+(x)\approx 1+\frac{1}{x^2}+\frac{3}{x^4}+\cdots-i\sqrt{\frac{\pi}{2}}xe^{-x^2/2}.$$

- For $|x|\ll 1$:

$$I_+(x)\approx -i\sqrt{\frac{\pi}{2}}x.$$

- For $|x|\gg 1$ and $|\mathrm{Re}\{x\}|\ll|\mathrm{Im}\{x\}|$ (with $\mathrm{Im}\{x\}<0$):

$$I_+(x)\approx -i\sqrt{2\pi}xe^{-x^2/2}.$$

Applying the approximations $V_{Ti}\ll u\ll V_{Te}$ and assuming the fluctuation frequency falls within the range $kV_{Ti}\ll\omega\ll kV_{Te}$, we expand the dispersion relation. Substituting the respective asymptotic limits for electrons and ions, the equation takes the following form:

$$\begin{aligned}1+&\frac{2\omega_{pe}^2}{k^2V_{Te}^2}\left(1+i\sqrt{\frac{\pi}{2}}\frac{\omega}{kV_{Te}}\right)-\omega_{pi}^2\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^2}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^2}\right)\right]\\&-3\,\omega_{pi}^2k^2V_{Ti}^2\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^4}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^4}\right)\right]+i\sqrt{2\pi}\,\frac{\omega_{pi}^2}{k^3V_{Ti}^3}(\mathbf{k}\cdot\mathbf{u}\sinh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right)\\&+\omega\cosh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right))\exp\left[-\frac{1}{2}\left(\frac{\omega}{kV_{Ti}}\right)^2\right]=0.\end{aligned} \tag{A.8}$$

A.2 Collisional Regime

Incorporating collisional effects into the kinetic theory, the general dispersion relation is modified as follows:

$$\begin{aligned}\varepsilon^L(\omega,\mathbf{k})=0\quad\Rightarrow\quad 1+\sum_{e,i}\frac{\omega_{p\alpha}^2}{k^2V_{T\alpha}^2}&\left[\frac{1-\ I_+(\frac{\omega+i\nu_{\alpha n}-\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}{1-i\frac{\nu_{\alpha n}}{\omega+i\nu_{\alpha n}-\mathbf{k}\cdot\mathbf{u}}I_+(\frac{\omega+i\nu_{\alpha n}-\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}\right]\\+\sum_{e,i}\frac{\omega_{p\alpha}^2}{k^2V_{T\alpha}^2}&\left[\frac{1-\ I_+(\frac{\omega+i\nu_{\alpha n}+\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}{1-i\frac{\nu_{\alpha n}}{\omega+i\nu_{\alpha n}+\mathbf{k}\cdot\mathbf{u}}I_+(\frac{\omega+i\nu_{\alpha n}+\mathbf{k}\cdot\mathbf{u}}{kV_{Te}})}\right]=0.\end{aligned} \tag{A.9}$$

Applying the physical approximations $\nu_{en} \gg \nu_{in}$, $\nu_{in} \ll \omega$, and $\nu_{in} < kV_{Ti}$, the dispersion equation simplifies to:

$$\begin{aligned}
&1+\frac{2\omega_{pe}^2}{k^2V_{Te}^2}-\omega_{pi}^2\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^2}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^2}\right)\right]\\
&-3\,\omega_{pi}^2\,k^2V_{Ti}^2\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^4}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^4}\right)\right]-\sqrt{\frac{\pi}{2}}\,\frac{\omega_{pi}^2}{k^3V_{Ti}^3}\nu_{in}\exp\left[-\frac{1}{2}\left(\frac{\omega}{kV_{Ti}}\right)^2\right]\\
&+i\sqrt{\frac{\pi}{2}}\,\frac{2\omega_{pe}^2}{k^2V_{Te}^2}\,\frac{\omega}{kV_{Te}}\,\frac{1}{1-\sqrt{\frac{\pi}{2}}\frac{\nu_{en}}{kV_{Te}}}+i\,\omega_{pi}^2\nu_{in}\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^3}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^3}\right)\right]\\
&+i\,9\,\omega_{pi}^2\,k^2V_{Ti}^2\,\nu_{in}\left[\left(\frac{1}{(\omega-\mathbf{k}\cdot\mathbf{u})^5}\right)+\left(\frac{1}{(\omega+\mathbf{k}\cdot\mathbf{u})^5}\right)\right]+i\sqrt{2\pi}\,\frac{\omega_{pi}^2}{k^3V_{Ti}^3}(\mathbf{k}\cdot\mathbf{u}\sinh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right)\\
&\qquad+\omega\cosh\left(\omega\frac{\mathbf{k}\cdot\mathbf{u}}{(kV_{Ti})^2}\right))\exp\left[-\frac{1}{2}\left(\frac{\omega}{kV_{Ti}}\right)^2\right]=0. \qquad (A.10)
\end{aligned}$$

Authors' Contributions

All authors contributed equally to this work.

Data Availability

No additional data are available.

Conflicts of Interest

The authors declare that there is no conflict of interest.

Ethical Considerations

The authors have diligently addressed ethical concerns, such as informed consent, plagiarism, data fabrication, misconduct, falsification, double publication, redundancy, submission, and other related matters.

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