Linear kinetic Alfvén waves in inhomogeneous plasma: Effects of Landau damping

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Abstract – The turbulent spectrum of kinetic Alfvén waves in inhomogeneous plasma is investigated in the presence of Landau damping. Inhomogeneities in transverse and parallel directions to the ambient magnetic field are incorporated in the dynamics. Numerical solutions of the equations governing kinetic Alfvén waves in the linear regime are obtained while retaining the effects of Landau damping, which have a significant impact on the frequency spectrum generated by propagating kinetic Alfvén waves. A semi-analytical model developed to elucidate the physics of this process is also described.

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Introduction. – Kinetic Alfvén waves (KAWs) play a major role in various space phenomena [1]. The dispersive properties of KAWs make them able to participate in a variety of phenomena [2] in which classic, nondispersive, Alfvén waves play no role. The CLUSTER data [3] shows that mode conversion from surface mode Alfvén waves to KAWs across the magnetopause [4] transports large electromagnetic energy fluxes across geomagnetic field lines from magnetosheath flows into the magnetosphere. KAW’s importance in particle heating and the energy cascade of solar wind turbulence is suggested by theoretical as well as by experimental studies [5,6]. KAWs are also known to play an important role in the stochastic ion heating mechanism in magnetopause [7].

The comparison of the experimental results in helicon plasmas [8] with the theoretically predicted dispersion relation for KAW including full kinetic effects [9] provides strong evidence for the excitation of KAW in inhomogeneous plasma. The kinetic theory may deviate KAW nonlinear processes from an ideal behaviour [10] quantitatively as well as qualitatively. But before going to a fully developed kinetic theory, it is better to study Landau damping in magnetic fluctuations. Landau damping is an important effect associated with KAWs [11,12]. Borgogno et al. [13] have used the Landau fluid model to study the filamentation of Alfvén waves due to density channels. Passot and Sulem [14,15] have derived a Landau fluid model to describe the magnetohydrodynamic waves in collisionless plasma where Landau damping is the main dissipation process. Hasegawa and Chen [16] have described the kinetic process of plasma heating by resonant excitation of shear Alfvén waves.

Motivated by the Houshmandyar and Scime observations [17], Sharma et al. [18] and Goyal et al. [19] examined the localization of KAWs in inhomogeneous plasma. These studies were based on a fluid model and kinetic effects like Landau damping which may play a vital role in the energy decay process were not considered. Landau damping may also affect the localization process and the magnitude of turbulence generated. Here we examine the effects of Landau damping on the linear propagation of KAWs in inhomogeneous plasma. The model equations are solved numerically to study the localization of KAW and resulting turbulent spectrum. A semi-analytical model developed to provide an insight into the physics is also described.

In the next section the dynamical equations for linearly propagating KAW with Landau damping effects are developed and then solved numerically. In the third section the results of the simulation of the dynamics of propagating KAW are presented together with a semi-analytic model of KAW dynamics based on a paraxial approximation. Conclusions and additional results are summarized in the last section.

Dynamical equations. – Consider the ambient magnetic field $\vec{B}_0$ along the $z$-axis. An inhomogeneous plasma
with density inhomogeneity in transverse as well as parallel directions is describable by $n_0 \exp(-x^2/L_z^2 + z^2/L_z^2)$, where $n_0$ is the background plasma density, $L_z$ is the inhomogeneity scale length in transverse (parallel) direction to the ambient magnetic field $B_0$. KAWs with low frequency $\omega$ and finite amplitude are assumed to be propagating in the $x$-$z$ plane, i.e., $k_0 = k_0 x \hat{x} + k_0 z \hat{z}$ where $k_0x$ ($k_0z$) is the component of the wave vector perpendicular (parallel) to the background magnetic field. Following the standard method [20], the dynamical equation for finite-amplitude KAW propagating in the $x$-$z$ plane is obtained:

i) the momentum balance equation:

$$\frac{\partial \vec{v}_j}{\partial t} + 2\gamma_L \vec{v}_j = \frac{q_j}{m_j} \vec{E} + \frac{q_j}{cm_j} \left( \vec{v}_j \times \vec{B} \right) - \frac{k_B}{m_j n_j} \nabla n_j T_j;$$

(1)

ii) the continuity equation:

$$\frac{\partial n_j}{\partial t} + \nabla \cdot (n_j \vec{v}_j) = 0;$$

(2)

iii) Faraday’s law

$$\nabla \times \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}}{\partial t};$$

(3)

iv) Ampère’s law

$$\nabla \times \vec{B} = \frac{4\pi}{c} \vec{J} + \frac{1}{c} \frac{\partial \vec{E}}{\partial t},$$

(4)

where $m_j$, $\vec{v}_j$, $n_j$, $T_j$, respectively, define the mass, fluid velocity, density and temperature of the species $j$ ($j = $ electrons $e$) or ions ($i$), $k_B$ is the Boltzmann constant, $\gamma_L$ is the phenomenologically incorporated Landau damping factor given by

$$\frac{\gamma_L}{\omega_0} \approx -\sqrt{\pi} \frac{V_A^2 k^2 \rho_i^2}{4},$$

(5)

which may be obtained by using the dispersion relation obtained by Lysak and Lotko [11] and Hasegawa and Chen [12]. $J = n_0 (v_i - v_e)$ is the current density. Here $\omega_0$ is the real part of the frequency $\omega$. From eq. (4):

$$\frac{\partial \vec{B}_y}{\partial t} = \frac{c}{e} \frac{\partial E_z}{\partial x} - \frac{c}{e} \frac{\partial E_z}{\partial z}.$$

(6)

Using eq. (2), we get

$$v_{iz} = i \omega n_i \left( ik_z n_0 + \frac{\partial}{\partial x} \delta n \right),$$

(7)

where $\delta n$ is the density perturbation.

Using electron dynamics of eq. (1), we obtain

$$E_z = \left[ ik_z \frac{k_B T_e}{e} + m_e \left( \frac{i \omega^2 + 2\gamma_L \omega}{k_z} \right) \left( \delta n \right) n_0 \right].$$

(8)

Using the velocity components and eq. (8) in the current density conservation equation (\nabla \cdot J = 0), we get

$$\frac{\partial E_z}{\partial t} = -\left\{ \left( \frac{\omega_i^2 + \frac{\partial^2}{\partial t^2}}{\omega_{ci}^2} \right) V_A \left( 1 - \frac{\delta n}{n_0} \right) \frac{\partial \vec{B}_y}{\partial z} \right\}.$$  

(9)

The $\partial^2 / \partial t^2$ in eq. (9) comes due to finite-frequency effects and in the context of low frequencies, this equation (9) reduces to the usual equation [20]. Combining eqs. (6), (8) and (9), we obtain the dynamical equation for KAW:

$$\frac{\partial^2 \vec{B}_y}{\partial t^2} + 2\gamma_L \frac{\partial \vec{B}_y}{\partial t} = \left( \frac{\partial^2}{\partial t^2} + 2\gamma_L \frac{\partial}{\partial t} \right) \lambda_e^2 \frac{\partial^2 \vec{B}_y}{\partial x^2} - V_A^2 k_z^2 \frac{\partial^4 \vec{B}_y}{\partial x^2 \partial z^2} + V_A^2 \left( \frac{\partial^2}{\partial t^2} + \frac{\partial^2}{\partial t^2} \right) \exp \left[ \frac{x^2}{L_x^2} + \frac{z^2}{L_z^2} \right] \frac{\partial^2 \vec{B}_y}{\partial x^2 \partial z^2}.$$  

(10)

where $\lambda_e = \sqrt{c^2 \rho_e / 4 \pi n_0 e^2}$ is the inertial length of the electrons, $V_A^2 = B_0^2 / 4 \pi n_0 m_e$ is the Alfven speed, $\rho_i = c_s / \omega_i$ is the ion sound gyroradius at the electron temperature, $\rho_i = \sqrt{T_i / m_i}$ is the ion sound speed, $v_{ti} = \sqrt{T_i / m_i}$ is the electron thermal speed, $v_{ti} = \sqrt{T_i / m_i}$ is the ion thermal speed. $\rho_i = v_{ti} / \omega_{ci}$ is the ion gyroradius and $\omega_{ci}$ is the ion cyclotron frequency. The inhomogeneities used in eq. (10) are just like perturbations and are very small so their spatial derivatives may be omitted in comparison to other terms in the equation. To study the localization process, we substitute the envelope solution

$$\vec{B}_y = \vec{B}_0(x, z, t) e^{i(k_0 x + k_0 z + \omega_{ti} t)}$$

(11)

into eq. (10) and obtain

$$- \frac{2i k_0 \lambda_e}{V_A^2 k_z^2} \left( 1 + \frac{i \gamma_L}{\omega_0} \right) \frac{\partial \vec{B}_0}{\partial t} + \frac{2i \gamma_L}{V_A^2 k_z^2} \vec{B}_0 =$$

$$\left[ -2i k_0 \frac{\lambda_e^2 \omega_0^2}{V_A^2 k_z^2} + 4 \gamma_L \frac{k_0 \omega_0}{V_A^2 k_z^2} + 2i k_0 \rho_i^2 \frac{\partial \vec{B}_0}{\partial x} \right] +$$

$$\left[ k_0 \frac{\lambda_e^2 \omega_0^2}{V_A^2 k_z^2} + \frac{\partial^2 \vec{B}_0}{\partial x^2} + \left( \frac{\rho_i}{V_A^2 k_z^2} + \frac{2i \gamma_L}{\omega_0} \right) \frac{\partial^2 \vec{B}_0}{\partial x \partial z} \right]$$

$$+ k_0^2 \xi \frac{\partial^2 \vec{B}_0}{\partial z^2} - \vec{B}_0 \exp \left( \frac{x^2}{L_x^2} + \frac{z^2}{L_z^2} \right) + \vec{B}_0.$$  

(12)

Equation (12) after normalization becomes

$$\left( 1 + \frac{i \gamma_L}{\omega_0} \right) \frac{\partial B}{\partial t} + 2c_{4 \xi_1} \xi_2 \left[ 1 - c_5 - 2i \xi_5 c_6 \frac{\partial^2}{\partial x^2} \right] \frac{\partial B}{\partial x} =$$

$$+ 2 \xi_1 \frac{\partial B}{\partial x} + i \xi_1 \left[ 1 - c_5 - 2i \xi_5 c_6 \frac{\partial^2}{\partial x^2} \right] \frac{\partial^2 B}{\partial x^2}$$

$$- \xi_1 \frac{\partial B}{\partial x} + i \xi_1 \left( \frac{\partial^2 B}{\partial x^2} + \frac{x^2}{\lambda_x^2} - 1 \right) B = 0.$$  

(13)

where $c_4$, $c_5$, $c_6$, $\xi_1$, $\xi_2$, $\xi_3$ are constants given by $c_4 = k_0 \rho_s$, $c_5 = \lambda_e^2 (1 + k_0^2 \rho_s^2) / \rho_i^2 (1 + k_0^2 \lambda_e^2)$, $c_6 = \sqrt{\xi_1 \lambda_x^2}$, $\xi_1 \approx 650$, $\xi_2 = \xi_3 \approx 50$ and the normalizing parameters are $x_n = \xi_2 \rho_s$, $z_n = 2 \xi_3 / k_0$. 

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$\tau_n = 2\omega_0 k_\perp V_A^2 \zeta x / \xi_\perp, \zeta x = L_x / \rho_s$ and $\zeta z = L_z / \rho_n$ are the normalized inhomogeneity scale lengths of the plasma channel. As our model is motivated from the experiment conducted in helium plasma, we have used the following laboratory plasma parameters [17] in our numerical study as well as in the semi-analytical analysis: $B_0 \approx 560$ G, plasma density, $n_0 \approx 6 \times 10^{12}$ cm$^{-3}$, $T_e \approx 81200$ K, $T_i \approx 2900$ K. These values yield characteristic plasma parameters: $V_A \approx 2.5 \times 10^7$ cm/s, $c_s \approx 1.3 \times 10^6$ cm/s, $\beta=\zeta^2 / V_A^2 \approx 0.002$, $\nu_{ce} \approx 1.1 \times 10^5$ cm/s, $\nu_{ci} \approx 2.445 \times 10^5$ cm/s, $\omega_{ci} \approx 2.68 \times 10^8$ rad/s, $\lambda_c \approx 0.1085$ cm, $\rho_i \approx 0.09$ cm, $\rho_s \approx 0.49$ cm, $k_{0x} \approx 0.85$ cm$^{-1}$, $k_{0z} \approx 0.01$ cm$^{-1}$ for $\omega_0 / \omega_{ci} = f_0 / f_{ci} = 0.1$ and $\omega_0 = 2.68 \times 10^8$ rad/s. For the parameters defined above, $x_n \approx 24.5$ cm for $k_{0x} \approx 0.85$ cm$^{-1}$, $z_n \approx 10^4$ cm for $k_{0z} \approx 0.01$ cm$^{-1}$, $t_n \approx 0.005$ s for $\omega_0 = 2.68 \times 10^8$ rad/s.

Equation (13) is mathematically similar to the nonlinear Schrödinger (NLS) equation if we neglect the second and fourth terms and replace the last term on the left-hand side by $(BB^*)B$. To solve eq. (13) numerically, we have adopted the algorithm developed for solving the standard NLS equation. Therefore, after testing the invariants of NLS up to $10^{-5}$ accuracy; we modified the algorithm for solving the NLS to accommodate eq. (13). For the initial hollow Gaussian profile:

$$B(x, z, 0) = a_0 (x^2 / r_{10}^2 + z^2 / r_{20}^2) \exp(-x^2 / r_{10}^2 - z^2 / r_{20}^2).$$

Equation (13) is solved numerically, where $a_0$ is the initial value of the amplitude of the KAW and $r_{10}$ ($r_{20}$) is the scale size of the initial wave field in the transverse (parallel) direction to the background magnetic field, where $r_{10} = k_{x1}^{-1}$ and $r_{20} = k_{z1}^{-1}$ with $k_x$ and $k_z$ the respective normalized wave numbers. The specific algorithm used to solve eq. (13) was developed by Sharma and Singh [21] and is based on a pseudo spectral method for space integration with periodic lengths $l_x = r_{10}, l_z = r_{20}$ and $64 \times 64$ grid points. A finite-difference method with predictor corrector scheme was employed for the evolution in time. The time step was taken to be of the order of $\Delta t \approx 10^{-4}$. To control the dynamics of KAW evolution, we choose $a_0 = 0.02$ and $k_0 = k_s = 0.2$ as parameters where the value of $a_0$ is based on the experimental conditions [17] and it defines the wave launching field normalized by the background magnetic field $B_0$. $r_{10}$ and $r_{20}$ respectively define the system lengths that extend beyond the size of a hollow Gaussian profile in transverse and parallel directions to $B_0$.

Simulation results and simplified model. – The normalized magnetic-field intensities as a function of time for fixed values of transverse and parallel inhomogeneity scale lengths are shown in figs. 1(a) and (b). The effect of Landau damping is clearly observed. When Landau damping effects are included, the corresponding peaks have the same spatial structures but have lower values of peak magnetic-field intensity (normalized). As expected, Landau damping clearly affects the linear behavior of KAWs. The magnetic-field fluctuations vs. time are shown in fig. 2 for a fixed spatial location $(x, z)$ with and without Landau damping effects. Looking carefully at the initial magnetic field fluctuations in fig. 2, the fluctuations with Landau damping effects have the same initial behavior then as time progresses the Landau damping effects contribute and the magnetic-field amplitude decays.

The variation of $|B_x|^2$ with $\omega$, the power spectra, obtained from Fourier transforms of the total magnetic-field fluctuations, are shown in fig. 3. The values shown in fig. 3 were obtained from the magnetic-field fluctuation data generated by the simulation for a time window $t \approx 300 \mu s$. The spectra show that the levels of turbulence vs. frequency are quite different with and without Landau damping for fixed values of inhomogeneity scale lengths. The nature (spectral structure) of the turbulence in fig. 3 is...
\[ \tilde{A}_0(\eta, z) = \frac{\tilde{B}_0}{\sqrt{f_0}} \left( (\sqrt{2} + \eta)^2 / 2 \right) \exp \left( - (\sqrt{2} + \eta)^2 / 2 \right) \exp(-k_1 z / 2), \]  

(20)

Fig. 2: (Colour online) Time series of magnetic field $B$ with and without Landau damping.

Fig. 3: (Colour online) Variation of $|B_z|^2$ vs. $\omega$ for KAW showing the Fourier spectrum for the data of fig. 2.

unchanged when the Landau damping effects are included, demonstrating the linear behavior of KAW propagation.

Because of the Landau damping factor introduced in the KAW dynamics, the localization changes with time for fixed values of inhomogeneity scale lengths, as described by eq. (13) and shown in figs. 1(a) and (b). As the time progresses, the spectral components of these spatially localized structures may evolve at different rates. Therefore, the magnetic-field fluctuation amplitude is expected to vary with time at a particular spatial location, as shown in fig. 2.

To better understand the results of the numerical simulation, we developed a simplified model using a semi-analytical approach and the paraxial approximation. The plasma is assumed to be homogeneous along the background magnetic field and the only inhomogeneity is in the transverse direction ($L_z \gg L_z$). Then eq. (10) can be approximated as

\[ \frac{\partial^2 \tilde{B}_y}{\partial t^2} + 2\gamma L \lambda^2 \frac{\partial \tilde{B}_y}{\partial t} = \left( \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} \right) \lambda^2 \frac{\partial^2 \tilde{B}_y}{\partial x^2} - \frac{V_A^2 \lambda^2}{\omega_0^2} \frac{\partial^2 \tilde{B}_y}{\partial x^2} \exp \left( \frac{x^2}{L_x^2} \right) \frac{\partial^2 \tilde{B}_y}{\partial z^2} \]  

(15)

Equation (15) is then solved assuming the paraxial approximation ($x \ll r_0 f_0$) [22], here $r_0$ is the scale size of the KAW in the transverse direction and $f_0$ is a dimensionless parameter that defines the beam width of the KAW. Substituting the envelope solution

\[ \tilde{B}_y = B_0(x, z) e^{i(k_0 x + k_0 z - \omega_0 t)} \]  

(16)

into eq. (15), the steady-state KAW equation becomes

\[ -2i k_0 z \frac{\partial B_0}{\partial z} - \frac{k_0^2}{V_A^2} \left( \frac{\lambda^2}{\omega_0^2} - \frac{\lambda^2}{k_0^2} \right) \frac{\partial^2 B_0}{\partial x^2} = \frac{2\gamma L \omega_0}{V_A^2} B_0 + k_0^2 \left( \frac{x^2}{L_x^2} \right) B_0 = 0. \]  

(17)

An additional eikonal $S_0$ is introduced to separate eq. (17) into real and imaginary parts by assuming a solution of the form

\[ B_0 = \tilde{A}_0(x, z) \exp \{ i k_0 z S_0(x, z) \}. \]  

(18)

The solution of eq. (17) is accomplished with a technique [23] analogous to the paraxial approximation in which eqs. (17) and (18) are expressed in terms of the parameters $\eta$ and $z$, i.e.,

\[ \eta = \left( x / r_0 f_0 - \sqrt{2} \right). \]  

(19)

$r_0 f_0$ is the beam width and is the position of maximum irradiance.

The solution of the real part is then [24]

see eq. (20) above

\[ S_0(\eta, z) = \left( (\sqrt{2} + \eta)^2 / 2 \right) \zeta(z) + \phi_0 \]  

(21)

\[ \zeta(z) = a \frac{f_0}{\sqrt{2}} \frac{d f_0}{d z}, \]  

(22)

where $a = V_A^2 \left( (\sqrt{2} + \eta)^2 \right) \frac{\lambda^2}{k_0^2} + k_0^2 \left( \frac{x^2}{L_x^2} \right)$ and $k_i = \omega_0 \gamma L / k_0 V_A^2$.

The governing differential equation for the KAW beam width parameter $f_0$ is

\[ \frac{d^2 f_0}{d z^2} = \frac{7}{12a^2 R_d^2 f_0} - \frac{f_0}{a L_x^2}, \]  

(23)
The right-hand side of eq. (23) plays a crucial role in the convergence and divergence of the wave as it contains the terms arising from diffraction and the inhomogeneous plasma profile. If these terms balance each other, self-trapping of the KAW occurs. In that case, the convergence or divergence of the wave does not occur (i.e., $f_0 = 1$) as it propagates along $z$. This gives us a critical transverse inhomogeneity scale length ($L_{ci}^{2}(Cr) \sim (12/7)\alpha R_d^2 \sim $), where $R_d \approx 4$ cm for $r_0 \sim 20$ cm.

To solve eq. (23) numerically, we assume boundary conditions corresponding to a plane-wave front, i.e., $f_0 = 1$ and $df_0/dz = 0$ at $z = 0$. The results are presented in figs. 4(a) and (b) for a fixed value of the inhomogeneity scale length $L_{ci}$ but for $\omega_0 = 0.1\omega_{ci}$ and $0.4\omega_{ci}$, respectively. By increasing the wave frequency from $0.1\omega_{ci}$ to $0.4\omega_{ci}$ the beam width parameter $f_0$ attains its minimum value faster than that at lower frequency. Therefore, the localization process becomes faster. This may be understood from eq. (23): the first term on the right-hand side viz. the diffraction term is inversely proportional to $a^2 R_d^2$. So, the diffraction term as a whole is inversely proportional to $\omega_0^2$. On the other hand, the second term on the right-hand side of eq. (23) viz. the inhomogeneous term is inversely proportional to $a$. It is obvious from the expression of $a$ that $a$ is proportional to $\omega_0^2$; therefore the inhomogeneous term is inversely proportional to $\omega_0^2$. Hence, as the value of $\omega_0$ increases, the converging effect dominates over the diverging effect and $f_0$ attains its minimum value faster than that at lower frequency (figs. 4(a) and (b)).

In addition to providing an insight into the localization process, the semi-analytical model also helps to explain the effect of Landau damping on the propagation of linear KAW in the steady state (eq. (20)). Figures 5(a) and (b) show the variation of wave magnetic-field amplitude vs. the distance of propagation ($z$) at wave frequencies $\omega_0 = 0.1\omega_{ci}$ and $0.4\omega_{ci}$, respectively. With the increase in wave frequency, the magnetic-field amplitude decreases faster with increasing distance of propagation. This effect can be understood from eq. (20) in which the wave magnetic field $A_0$ depends on $f_0$ and the Landau damping factor $k_i$. From the expression of $k_i$, it is clear that $k_i$ linearly depends upon $\omega_0$. Therefore, at higher frequencies, the Landau damping factor $k_i$ increases. As discussed above, at higher frequency, $f_0$ attains its minimum values faster than at lower frequency. Also, from the expression of $A_0$ in eq. (20), it is clear that the wave magnetic field $A_0$ decays exponentially faster with increasing wave frequency. The net effect is also seen in figs. 5(a) and (b).

The semi-analytical model also shows the Landau damping effect for steady-state systems also observed in the numerical simulations. Just as figs. 1(a) and (b) clearly indicate the role of Landau damping in KAW propagation in the spatio-temporal domain (the wave intensity is attenuated and the localization of KAW occurs due to inhomogeneous plasma effects), figs. 5(a) and (b)...
demonstrate the same Landau damping effect on wave propagation in inhomogeneous plasma in the spatial domain. Thus, the results obtained using the semi-analytical model are consistent with those obtained using numerical simulations.

**Conclusion.** – With Landau damping, the model equations have been derived and solved using the numerical simulation technique as well as a semi-analytical method. With the inclusion of the term arising due to linear Landau damping, the wave dynamics is shown to follow modified NLS equation.

Landau damping is the main dissipation process due to electrons in collisionless plasma. Collisionless heating depends on the parallel electric field and owing to its small perpendicular wavelength, KAW accompanies this field. It is identified as one of the dissipation mechanisms of KAW which appears because of this parallel electric field. The inhomogeneities included in the dynamics help in forming wave packets with varying wave numbers resulting into the localization of KAW. As Landau damping is also wave number dependent, so it affects the amplitude of localization, level of turbulence etc. Landau damping has a profound effect on wave propagation and energy decay as indicated by the wave amplitudes and the resultant frequency spectra. This analysis suggests that Landau damping of KAWs in inhomogeneous plasmas could play an important role in particle heating in space and laboratory plasmas. Further, it may also be the probable mechanism for the damping of interplanetary magnetic field.

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