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HEATING BY ALFVÉN WAVE FILAMENTATION IN THE SOLAR CORONA AND THE INTERSTELLAR MEDIUM

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RESUMEN

Se propone un nuevo mecanismo de calentamiento para plasmas intensamente magnetizados, el cual resulta de la generación a pequeña escala, producida por colapso transverso (filamentación) de ondas Alfvén circulares, débiles, no lineales, polarizadas. Esto significa que en hoyos coronales abiertos, las ondas cuya frecuencia no excede unas pocas décimas de Hertz, contribuyen al calentamiento y a la aceleración del viento rápido solar, mientras que este límite inferior se reduce a unos cuantos Hertz en los lazos coronales. También se demuestra la importancia de la filamentación de las ondas de Alfvén en relación al arrastre ambipolar, en la fase tibia ionizada del medio interestelar.

ABSTRACT

A new heating mechanism for strongly magnetized plasmas is proposed. It results from small-scale generation produced by transverse collapse (filamentation) of weakly nonlinear circularly polarized Alfvén waves. It turns out that in open coronal holes, waves whose frequency exceeds a few tens of Hertz contribute to the heating and to the acceleration of the fast solar wind, while this lower bound is reduced to a few Hertz in coronal loops. The importance of Alfvén waves filamentation relatively to ambipolar drift, is also demonstrated in the warm ionized phase of the interstellar medium.

Key words: **ISM: MAGNETIC FIELDS — MHD — PLASMAS — SUN: CORONA**

1. INTRODUCTION

The aim of the present paper is to suggest a new mechanism for the heating of strongly magnetized plasmas ($\beta < 1$) like the quiet regions of the solar corona and certain regions of the Interstellar Medium (ISM). In the lower corona where $\beta \ll 1$, two mechanisms have currently been proposed: a flare type heating in regions of high currents and the dissipation of hydromagnetic waves in quiet regions (see e.g., Narain & Ulmschneider 1996). For the latter process, only large amplitude Alfvén waves dissipation has been considered to be efficient, making this heating mechanism relevant only at large distances (10 to $20R_\odot$) from the sun (Esser et al. 1986; Hollweg 1992). Alfvén wave dissipation is also considered as a heating mechanism for the ISM (Begelman 1995), where these waves appear to be a main mechanism to mediate energy from cosmic rays to the background gas. Note that Alfvén wave dissipation can only be relevant when the involved time and length scales are smaller than those associated with the ambipolar drift, a condition which can only be satisfied outside molecular clouds.

The proposed mechanism results from the collapse (filamentation) of small amplitude Alfvén waves propagating along an external magnetic field in a homogeneous medium. This effect, which has extensively been studied in the context of nonlinear optics, refers to the violent concentration (focusing) of the wave energy in planes transverse to the propagation direction. The resulting formation of small scales allows dissipation processes like ion-cyclotron resonance or Landau damping to act, leading to a heating of the plasma.

2. INDIVIDUAL WAVE COLLAPSE AND DISSIPATION

As it was discussed in Champeaux, Passot, & Sulem 1997a, for $\beta < 1$ the wave focusing develops along the propagation direction (convective instability). A multiple-scale analysis shows that in the simple case where the modulation of the wave-train is assumed stationary, the Alfvén wave amplitude evolves according to the two-dimensional nonlinear Schrödinger equation (NLS)

$$i\partial_\xi B + \frac{1}{4k(1-\beta)}\Delta_\perp B + \frac{1+4\beta}{4\beta}k|B|^2 B = 0, \tag{1}$$

where k represents the wave-number and ξ denotes the coordinate along the direction of propagation. The term involving the laplacian with respect to the transverse variables describes the diffraction of the wave. The nonlinear term reflects the interactions of the Alfvén wave with the low frequency hydromagnetic fields, which are stirred by the amplitude modulation and modify the "refractive index" of the medium. These fields are given in terms of the wave amplitude by

$$\bar{\rho} = \frac{|B|^2}{2\beta}, \qquad \bar{u}_x = 0, \qquad \bar{b}_x = -|B|^2. \tag{2}$$

The coefficients in eq. (1) have been evaluated in the dispersionless limit, which corresponds to situations often encountered in astrophysical plasmas, where the Alfvén wavelength is much larger than the gyromagnetic radius of the ions. In this context, the convective instability can be understood in the following way. A local increase of the plasma density reduces the Alfvén velocity

$$v_A = \frac{\mathcal{B} + \bar{b}_x}{\sqrt{4\pi(\rho_0 + \bar{\rho})}}, \tag{3}$$

where $\mathcal{B}$ is the magnitude of the ambient magnetic field and ρ_0 the unperturbed density. This produces a bending of the wave front which in turn leads to an enhancement of the transverse magnetic field. The resulting density increase together with the reduction of the longitudinal magnetic field imply a further decrease of the Alfvén velocity. The plasma being magnetically dominated, the counter effect of thermal pressure is insufficient to balance the resulting transverse Lorentz force, which confines the particles within a filament around which the magnetic field spirals. In the longitudinal direction, the Lorentz force is canceled by the longitudinal pressure gradient, preventing the development of a longitudinal velocity.

A main property of the two-dimensional NLS equation is the possible blow-up of the wave amplitude (collapse) at a finite propagation distance (Vlasov, Petrischev, & Talanov 1971; Glassey 1977). This collapse requires that the wave energy $N = \frac{1}{2}\int |B|^2 d\mathbf{x}_\perp$ in each transverse plane exceeds a critical value N_T associated to the energy of the so called Townes soliton (Weinstein 1983). When the wave collapses, N_T is indeed the amount of energy captured in the focus (Landman et al. 1988; Lemesurier et al. 1988). Physically, N_T appears to be an energy per unit length which when taking into account the coefficients entering eq. (1), is given (in c.g.s. units) by

$$N_T = \frac{5.84\beta\mathcal{B}^2}{4\pi k^2}. \tag{4}$$

In reality, this singularity is never reached because dissipation processes become relevant near the collapse. Indeed, when the filament diameter (transverse characteristic scale) becomes of the order of the length scale $2\pi/k_d$ required for dissipation to be relevant, the energy carried by Fourier of wavenumbers exceeding k_d is dissipated. As a consequence, the energy becomes smaller than N_T, the collapse is arrested and the remaining energy dispersed. The amount of energy δN dissipated in one such event has been estimated in Dyachenko et al. 1992, where it is found that for very large k_d, the energy loss scales like

$$\delta N \sim (\ln\ln k_d)^{-1}. \tag{5}$$

This dependence, although very weak, appears to be relevant when dealing with scale separations such as those present in the solar corona or in the ISM. From the numerical data presented in Dyachenko et al. 1992, Champeaux et al. 1997b, estimated at about 5% the relative energy loss $r = \frac{\delta N}{N}$ in one collapse in regions of the solar corona or of the ISM considered in sections 4 or 5 respectively.

3. GLOBAL ENERGY DISSIPATION

Two different processes dissipate energy in a turbulent regime: the resonant wave-wave interactions which lead to a gradual energy transfer towards small scales and wave collapse which produces a sudden energy transfer. The latter process is dominant when dealing with low amplitude waves, but a comparable contribution can originate from the former phenomenon in a strongly turbulent regime (Newell, Rand, & Russell 1988). Dissipation due to collapses thus provides a lower bound for the total dissipation.

In the large k_d limit, the amount of energy dissipated in the mean by one Alfvén wave collapse is $\langle r\rangle N_T$ (where $\langle r\rangle$ refers to an average over all collapses). The associated power is thus given by $P = \langle r\rangle N_T v_A$. Assuming that the turbulent fluctuations are space filling, the number of collapses contained in a volume limited by an area S and a distance L along the ambient magnetic field is estimated by $\nu = SL/L_\perp^2 L_x$, where $L_\perp$ denotes the transverse scale of the most unstable mode for the convective instability and L_x its e-folding length. The number of individual collapses required to significantly reduce the Alfvén wave energy flux $F_e = \frac{(\delta B)^2}{4\pi} v_A$ being large, one is lead to write the differential equation

$$\frac{dF_e}{dL} = -\frac{4\pi}{\mathcal{B}^2 l_0 v_A} F_e^2 \ , \tag{6}$$

with $l_0 = \frac{4\pi^2\beta\lambda}{5.84<r>}$. Here δB is the dimensional Alfvén wave amplitude and λ its wavelength. Assuming constant values of β and $\mathcal{B}$ over the propagation distance of the Alfvén waves it follows that $F_e(L) = \frac{F_e(0)}{1+(L/L_d)}$, where

$$L_d = \frac{4\pi^2\beta\lambda}{5.84 < r >}\frac{1}{(\frac{\delta B}{\mathcal{B}})^2}, \tag{7}$$

defines the dissipation length associated to Alfvén wave filamentation.

4. ALFVÉN WAVE DISSIPATION IN THE SOLAR CORONA

At the base of open coronal holes where $n \approx 10^8 \mathrm{cm}^{-3}$, $\mathcal{B} \approx 10$ G, and $T \approx 10^6 K$, the β of the plasma is $\beta = c_s^2/v_A^2 \approx 5.79 \times 10^{-3}$. For Alfvén waves having a typical wavelength $\lambda = 2\pi v_A/\omega \approx 1.37 \times 10^4$ km, and a relative amplitude $B_0 = \sqrt{2}\delta B/\mathcal{B} \approx 1.95 \times 10^{-2}$, the transverse wavelength of the most unstable mode for the modulational instability and its e-folding length are $L_\perp = \lambda\sqrt{\beta}/B_0 \approx 5.36 \times 10^4$ km and $L_x = 2\pi/K_x = 4\beta\lambda/B_0^2 \approx 8.38 \times 10^5$ km, respectively. Requiring the equality between the magnetic energy per unit length and the critical energy for collapse N_T, the critical transverse length scale L_c for Alfvén wave filamentation is defined in the form $L_c^2 = \frac{5.84\beta\lambda^2\mathcal{B}^2}{4\pi^2(\delta B)^2}$. The value of L_c at the base of coronal holes, $L_c \approx 2.91 \times 10^4$ km, implies that the energy contained in one wavelength of the most linearly unstable mode is enough to trigger a nonlinear collapse. The corresponding dissipation length in such a region becomes

$$L_d = 5.35 \times 10^7 \ (\frac{T}{10^6\ \mathrm{K}})(\frac{n}{10^8\ \mathrm{cm}^{-3}})^{-1/2}(\frac{\mathcal{B}}{10\ \mathrm{G}})(\frac{\omega}{\mathrm{rad\ s}^{-1}})^{-1}(\frac{v_{\mathrm{r.m.s.}}}{30\ \mathrm{km\ s}^{-1}})^{-2}\ \mathrm{km}.$$

The above expression provides a condition on the Alfvén waves frequency to contribute to coronal heating. For example, for a value of L_d of the order of $0.5R_\odot$, required to accelerate the fast solar wind up to its observed velocity (McKenzie, Banaszkiewicz, & Axford 1995), we obtain that frequency $f = \frac{\omega}{2\pi} > 27\ Hz$ is needed in open coronal holes. In regions of higher density, as the coronal loops, the condition is less stringent. For example, taking $\mathcal{B} = 15$ G and $N = 5 \times 10^{10}$ cm^{-3}, the minimum frequency becomes of order $f = 1.7\ Hz$. Such frequencies appear realistic if the Alfvén waves are generated by the supergranular network through small-scale reconnection events (McKenzie et al. 1995).

5. ALFVÉN WAVE DISSIPATION IN THE INTERSTELLAR MEDIUM

In the warm intercloud medium, with a temperature $T \approx 8000$ K, a number density $n \approx 0.2$ cm^{-3} and an ambient magnetic field $\mathcal{B} \approx 4\ \mu$G (see e.g., Kulkarni & Heiles 1987; Cox 1991), the β of the plasma is $\beta \approx 0.23$. Assuming velocity fluctuations $v_{\mathrm{r.m.s.}} \approx 0.5\ c_s$, the relative amplitude is $B_0 \approx 0.34$. The relevant length scales for the filamentation of Alfvén waves with typical wavelength $\lambda \approx 1.25 \times 10^{-4}(\omega/\mathrm{rad\ yr}^{-1})^{-1}$ pc, are $L_x \approx 5.19 \times 10^{-4}(\omega/\mathrm{rad\ yr}^{-1})^{-1}$ pc and $L_\perp \approx 1.45 \times 10^{-4}(\omega/\mathrm{rad\ yr}^{-1})^{-1}$ pc.

For the hot phase, in the solar neighborhood $T \approx 10^6$ K, $n \approx 10^{-3}$ cm^{-3} and $B \approx 3\ \mu$G (Ferrière 1995), which implies $\beta \approx 0.64$ and $B_0 \approx 0.40$. For Alfvén waves with $\lambda \approx 1.33 \times 10^{-3}(\omega/\mathrm{rad\ yr}^{-1})^{-1}$ pc, we get $L_x \approx 5.97 \times 10^{-3}(\omega/\mathrm{rad\ yr}^{-1})^{-1}$ pc and $L_\perp \approx 2.35 \times 10^{-3}(\omega/\mathrm{rad\,yr}^{-1})^{-1}$ pc. Since in the ISM the amplitude of the waves is significant, the dissipation due to the resonant wave interactions comes into play and its contribution becomes comparable to the burn-out due to collapses. Taking this effect into account, we obtain a dissipation length $L_d \approx 5 \times 10^2\ \lambda$ in the warm phase (resp. $L_d \approx 10^3\ \lambda$ in the hot phase).

Ambipolar drift prevents the propagation of Alfvén waves with wavelengths smaller than $\lambda_A = \pi v_A/\nu_{ni}$, where ν_{ni} is the neutral-ion collision frequency. The associated dissipation length scale is given by $L_{Ad} = \lambda^2/\lambda_A$. In the warm medium with low ionization ($x \approx 6 \times 10^{-3}$), $\lambda_A \approx 1.2$ pc and the ambipolar diffusion dominates for wavelengths smaller than ≈ 618 pc. In the hot medium, $\lambda_A \approx 93$ pc because of the small value of the density, making the ambipolar diffusion again dominant.

In contrast, in the fully ionized warm medium, $\lambda_A = 2.36 \times 10^{-4}$ pc. For a wave with $\lambda \approx 0.1$ pc, we get $\frac{L_{Ad}}{\lambda} \approx 424$. The wavelength at which the dissipation due to filamentation are comparable to that of the ambipolar diffusion is thus $\lambda_c \approx 0.12$ pc. Alfvén wave filamentation can thus significantly contribute to the heating of the warm ionized medium.

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