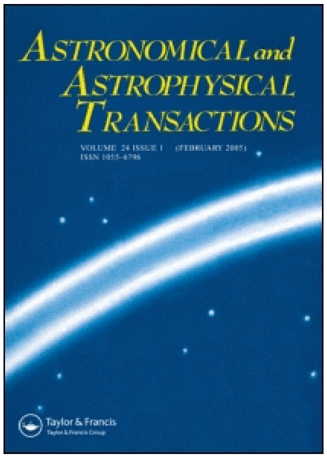

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THERMAL INSTABILITY IN MAGNETIZED INTERSTELLAR CLOUDS

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The effect of a magnetic field on the ionizational-thermal acoustic instability in the cold diffuse ISM is discussed in this paper. It is found that the frozen-in magnetic field leads to new types of instabilities called "magnetoacoustic" instabilities. These instabilities fall into two classes: slow and fast magnetoacouistic instabilities, corresponding to two types of waves in an MHD plasma. The fast modes are found to be effectively suppressed in comparison with the slow modes at the mean Galactic intensities of the magnetic field. Some implications of those results are discussed.

KEY WORDS MHD instabilities, Interstellar medium

1 INTRODUCTION

Thermal instability plays a fundamental role in the formation of interstellar medium (ISM) phases. It was first considered in the pioneering work of Field (1965) and applied to diffuse ISM by Pikel'ner (1967) and Field, Goldsmith and Habing (1969). It has been shown (Field, 1965) that there exist three types of thermally unstable modes: isobaric, isochoric and isentropic. The isobaric and isochoric modes are dynamical but the isentropic mode is oscillatory. The dominating mode in the diffuse ISM is isobaric.

A frozen-in magnetic field leads to the suppression of the thermal instability modes in the case of perturbations with vector $\mathbf{k} \perp \mathbf{B}$ (Field, 1965; Massaglia *et al.*, 1985). Magnetic field diffusion and Joule heating constitute new sources of instability: resistive and current-driven modes (Heyvaerts, 1974; Bodo *et al.*, 1987). In interstellar clouds with density $n > 10$ cm^{-3} the ionization fraction depends on the ionization and recombination processes. Defouw (1970) and Yoneyama (1973) have shown that ionization and recombination processes may have a stabilizing effect on the isobaric thermal mode in some temperature intervals. The influence of time-dependent ionization and recombination on the thermal instability has been extensively investigated by Corbelli and Ferrara (1995) in the case of hot photoionized hydrogen gas.

Flannery and Press (1979) have pointed to the existence of an ionization-coupled oscillatory isentropic thermal instability in diffuse cold ISM heated by cosmic rays and cooled by electron collisional excitations of trace element ions. This instability develops on intermediate length scales with maximum growth rates at $\lambda \sim 1$ pc and characteristic times $t_{\rm growth} \sim 1$–10 mil yrs.

The magnetic field must lead to additional magnetosonic modes of the thermal instability in the case of a variable ionization fraction. This effect is discussed in this paper. The physics of the acoustic instability is the following: if we compress an element of a fluid on intermediate time scales then at the beginning of compression the fluid behaves almost isothermally, i. e. the adiabatic heating is balanced by the increased cooling and thermal balance is very quickly achieved. Then at the end of the compression the recombination process develops and heating increases. This increases the gas pressure and work is done on the decompressing gas thus increasing the amplitude of the acoustic wave. The investigation of the modification of this mode by the magnetic field is the basic goal of our work.

2 BASIC EQUATIONS AND PHYSICAL PROCESSES

We consider the linear theory of the thermal instability of a homogeneous medium heated by cosmic rays and grain photoemission and cooled by electron collisional excitations of singly ionized carbon. The matter is ionized by cosmic rays. The initial state of the matter is determined by the equlibrium between heating *and* cooling as well as between ionization and radiative recombinations. The matter is threaded by a homogeneous magnetic field of the mean intensity of the Galactic disk. The development of the thermal instability is described by the following set of MHD equations:

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \cdot \mathbf{v}) = 0, \tag{1}$$

$$\frac{\partial x}{\partial t} + (\mathbf{v} \cdot \nabla) x = S_i, \tag{2}$$

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla)\mathbf{v} = -\frac{1}{4\pi\rho}[\mathbf{B} \times [\nabla \times \mathbf{B}]] - \frac{\nabla p}{\rho}, \tag{3}$$

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times [\mathbf{v} \times \mathbf{B}] + \nu_m \Delta \mathbf{B}, \tag{4}$$

$$\rho c_p \left(\frac{\partial T}{\partial t} + (\mathbf{v} \cdot \nabla) T \right) + p \nabla \cdot \mathbf{v} = -\mathcal{L}, \tag{5}$$

where $x = n_e/n$ is the ionization fraction, S_i is the total ionization source per atom (equal to the ionization minus the recombination functions); $c_p = k_b/(m(\gamma - 1)$ is the specific heat of the gas at constant pressure, the adiabatic index is $\gamma = 5/3$ and $\mathcal{L} = \Lambda - \Gamma$ (in ergs cm^{-3} s^{-1}) is the net cooling function equal to the difference between the cooling and heating functions. All the other variables have their usual meaning.

As one can see from equations (1)–(5) we incorporate the equation of non-stationary ionization (2) into the usual system of equations for the thermal instability. The net cooling function depends on the ionization fraction in this case. Therefore we investigate the instability connected with the behaviour of the ionized component of the gas which is determined by the ionization mechanisms as well as by the magnetic field.

Let us define the following dimensionless variables:

$$\tilde{\rho} = \rho/\rho_0, \quad \tilde{\mathbf{v}} = \mathbf{v}/c_s, \quad \tilde{x} = x/x_0, \quad \tilde{\mathbf{B}} = \mathbf{B}/B_0, \quad \tilde{T} = T/T_0, \quad \tilde{\mathbf{r}} = \mathbf{r}/l_0, \quad \tau = t/\tau_h,$$

$\tilde{\nu}_m = \nu_m \tau_h / l_0^2$ where $\tau_h = \frac{\rho c_p T}{\Gamma}$ is the characteristic heating time, c_s is the isothermal sound speed in the gas, $\mathbf{r}$ is the radius-vector and $l_0 = c_s \tau_h$ is the characteristic thermal lengthscale. The variables with subscript "0" are the characteristic values of the corresponding physical quantities.

It is easy to derive the following system of MHD equations in dimensionless form:

$$\frac{\partial \tilde{\rho}}{\partial \tau} + \nabla(\tilde{\rho} \cdot \tilde{\mathbf{v}}) = 0, \tag{6}$$

$$\frac{\partial \tilde{x}}{\partial \tau} + (\tilde{\mathbf{v}} \cdot \nabla)\tilde{x} = \frac{S_i \tau_h}{x_0} = \tilde{S}_i, \tag{7}$$

$$\frac{\partial \tilde{\mathbf{v}}}{\partial \tau} + (\tilde{\mathbf{v}} \cdot \nabla)\tilde{\mathbf{v}} = -A^2[\tilde{\mathbf{B}} \times [\nabla \times \tilde{\mathbf{B}}]] - \frac{\nabla \tilde{\rho}}{\tilde{\rho}} - \frac{\nabla \tilde{T}}{\tilde{T}}, \tag{8}$$

$$\frac{\partial \tilde{\mathbf{B}}}{\partial \tau} = \nabla \times [\tilde{\mathbf{v}} \times \tilde{\mathbf{B}}] + \tilde{\nu}_m \Delta \tilde{\mathbf{B}}, \tag{9}$$

$$\frac{1}{\tilde{T}}\left(\frac{\partial \tilde{T}}{\partial \tau} + (\tilde{\mathbf{V}} \cdot \nabla)\tilde{T}\right) + (\gamma - 1)\nabla \cdot \tilde{\mathbf{v}} = -\frac{\mathcal{L}\tau_h}{\rho c_p} = -\tilde{\mathcal{L}}. \tag{10}$$

The role of the seperate terms in system (6)–(10) depends on the value of the dimensionless parameter A which denotes the ratio of the Alfvenic speed to the sound speed. This parameter determines the influence of the magnetic field on the thermal instability. The mean value of A is of order unity in diffuse ISM.

Linearizing the system (6)–(10) and substituting perturbations in the form of normal modes $f' = f_1 \exp(i\mathbf{K} \cdot \mathbf{R} + n\tau)$ we obtain the following system of linear algebraic equations:

$$n\rho_1 + i(\mathbf{K} \cdot \mathbf{v}_1) = 0, \tag{11}$$

$$nx_1 = S_x x_1 + S_\rho \rho_1 + S_T T_1, \tag{12}$$

$$n\mathbf{v} = -iA^2[\mathbf{b} \times [\mathbf{K} \times \mathbf{b}_1]] - i\mathbf{K}(\rho_1 + T_1), \tag{13}$$

$$n\mathbf{b}_1 = i\mathbf{K} \times [\mathbf{v}_1 \times \mathbf{b}] - \nu_m K^2 \mathbf{b}_1, \tag{14}$$

$$nT_1 + (\gamma - 1)(\mathbf{K} \cdot \mathbf{v}_1) = -\mathcal{L}_x x_1 - \mathcal{L}_\rho \rho_1 - \mathcal{L}_T T_1, \tag{15}$$

where n is the increment of the wave, $\mathbf{K}$ is the wave vector, and the subscripts "1" denote Fourier amplitudes of corresponding perturbations. The subscripts x, ρ, T

denote corresponding partial derivatives of the cooling function and the ionization source function. The system of equations (11)–(15) describes waves with different relative orientation of vectors $\mathbf{K}$, $\mathbf{v}_1$, and $\mathbf{b}$. The case with $\mathbf{v}_1 \perp \mathbf{K}$ and $\mathbf{v}_1 \perp \mathbf{b}$ corresponds to Alfvenic waves which are not considered in this paper. Hence the system (11)–(15) written in coordinate form can be split up into two systems, one for Alfvenic waves and the other for compression waves. Alfvenic waves will dissipate on a time scale determined by the magnetic viscosity. We are not interested here in damped modes, so we retain only the system describing compression modes.

This system will have a non-trivial solution if its determinant equals zero. Finding the determinant and equating it to zero we obtain a 6th degree dispersion equation:

$$a_0 n^6 + a_1 n^5 + a_2 n^4 + a_3 n^3 + a_4 n^2 + a_5 n + a_6 = 0 \qquad (16)$$

with the following coefficients:

$$a_0 = 1,$$

$$a_1 = \mathcal{L}_T - S_x + \nu_m K^2,$$

$$a_2 = (\gamma + A^2 + (\mathcal{L}_T - S_x)\nu_m)K^2 + \mathcal{L}_x S_T - \mathcal{L}_T S_x,$$

$$a_3 = \gamma \nu_m K^4 + \left((\mathcal{L}_T - S_x)A^2 + (\mathcal{L}_x S_T - \mathcal{L}_T S_x)\nu_m + \mathcal{L}_T - \mathcal{L}_\rho - \gamma S_x\right) K^2,$$

$$\begin{aligned} a_4 = {} & (A^2 \cos(\phi)^2 \gamma + (\mathcal{L}_T - \mathcal{L}_\rho - \gamma S_x)\nu_m)K^4 \\ & + ((\mathcal{L}_x S_T - \mathcal{L}_T S_x)A^2 + \mathcal{L}_x S_T - \mathcal{L}_T S_x + \mathcal{L}_\rho S_x - \mathcal{L}_x S_\rho)K^2, \end{aligned}$$

$$a_5 = \left((\mathcal{L}_T - \mathcal{L}_\rho - \gamma S_x)A^2 \cos(\phi)^2 + (\mathcal{L}_x S_T - \mathcal{L}_T S_x + \mathcal{L}_\rho S_x - \mathcal{L}_x S_\rho)\nu_m\right) K^4,$$

$$a_6 = (\mathcal{L}_x S_T - \mathcal{L}_T S_x + \mathcal{L}_\rho S_x - \mathcal{L}_x S_\rho)A^2 K^4 \cos(\phi)^2,$$

In order to study the dependence of the growth rates on the wavenumber and temperature we must solve the dispersion equation numerically.

3 NUMERICAL SOLUTION OF THE DISPERSION RELATION

We consider a medium which is heated by cosmic rays and photoelectrons from grains and ionized by cosmic rays. The corresponding rates are $h = 5 \times 10^{-26}$ ergs s^{-1} and $\zeta = 1 \times 10^{-17} s^{-1}$. The cooling is produced by excitations of fine structure lines of C^+ ions in collisions with H atoms and electrons. For the cooling function on the C^+ (158μ) transition we use the expresion from Wolfire *et al.* (1995).

For $n \ll n_{\rm cr} \approx 3 \times 10^3$ cm^{-3} the cooling per hydrogen nucleus is given by

$$n\Lambda_{CII} = 2.54 \times 10^{-14} \mathcal{A}_C f_{CII} [\gamma^{H_0} n_{H_0} + \gamma^e n_e] \exp(-92/T) \text{ ergs s}^{-1} \text{ H}^{-1},$$

where γ^{H_0} and γ^e are the collisional de-excitation rate coefficients with neutral hydrogen and electrons respectively, and f_{CII} is the fraction of C^+ in C and $\mathcal{A}_C = n_C/n$ is the carbon abundance. We use the rates

$$\gamma^{H_0} = 8.86 \times 10^{-10} \text{ cm}^3 \text{ s}^{-1},$$

and

$$\gamma^e = 2.1 \times 10^{-7} T_2^{-0.5} \Omega(T) \text{ cm}^3 \text{ s}^{-1},$$

where $T_n \equiv T/10^n$ K and the collision strength $\Omega(T)$ is given by

$$\Omega(T) = 1.80 + 0.484 T_4 + 4.01 T_4^2 - 3.339 T_4^3.$$

We take $\mathcal{A}_C = 3 \times 10^{-3}$ and $f_C = 1$. For the ionization source S_i we take

$$S_i = \zeta(1 - x) - x n^2 \alpha(T),$$

where $\alpha(T) = 2.6 \times 10^{-13} T_4^{-0.6}$ cm s^{-1} is the recombination coefficient.

We do not take the magnetic viscosity into account because we expect it to work appreciably only on very small spatial scales or on very long time scales.

The results of calculations for typical conditions of the Galactic disk are demonstrated in the figures.

Figure 1(a) shows the dependence of the growth times of instabilities on the wavelength for slow and fast magnetosonic waves (SMSW and FMSW) in the temperature interval 70–90 K. The graphs are plotted for $\phi = \pi/4$. The fast magnetosonic modes (the curves in the upper right conner) are effectively suppressed at the mean Galactic intensities of the magnetic field ($B = 2 \times 10^{-6}$ Gauss). FMSW have a maximum growth rate at $\lambda = 10$–70 pc and the corresponding growth time approximately equals 700 million years.

On the contrary the instability of SMSW is not absolutely suppressed in comparison with its FM counterpart and have the same maximum growth rates as the instability of sonic waves in the absence of a magnetic field. It is interesting to note that the wavelengths of the most rapidly growing perturbations are shifted towards the smaller values when the angle ϕ between the wavevector and the direction of the magnetic field increases towards $\pi/2$. This behaviour can be clearly seen in Figure 2. This result can be easily explained if we take into account the fact that the oscillatory mode reaches its maximum growth rate when the period of the wave approximately equals the characteristic ionization–recombination time (Corbelli and Ferrara, 1995). Hence if the phase velocity of the wave decreases at the increasing angle ϕ (as in the case of slow magnetosonic waves), then the wave number must increase at the same frequency.

The characteristic growth times for SMSW are $\sim 10^6$ years at T less than 90 K. At such low temperatures the instability of SMSW develops only when the wavelength is greater than some critical value $\lambda_{\rm cr}$. The FMSW have a similar behaviour in this respect but for a wider temperature interval $T = 50$–500 K.

At $T \approx 80$ K there begins a transition from oscillatory to dynamical instability for SMSW. The critical wavelength $\lambda_{\rm cr}$ quckly decreases with increasing temperature and reaches zero at $T \approx 87$ K. $\lambda_{\rm max}$ remains non-zero until $T = 90$ K.

In the transition region ($T = 85$–95 K) the instability of SMSW has the character of standing growing oscillating waves because the phase velocity is zero here while the frequency remains non-zero. At $T = 90$–100 K (see Figure 1(b)) the instability has a dynamical character on small wavelengths, while for $\lambda > 1$ pc it becomes

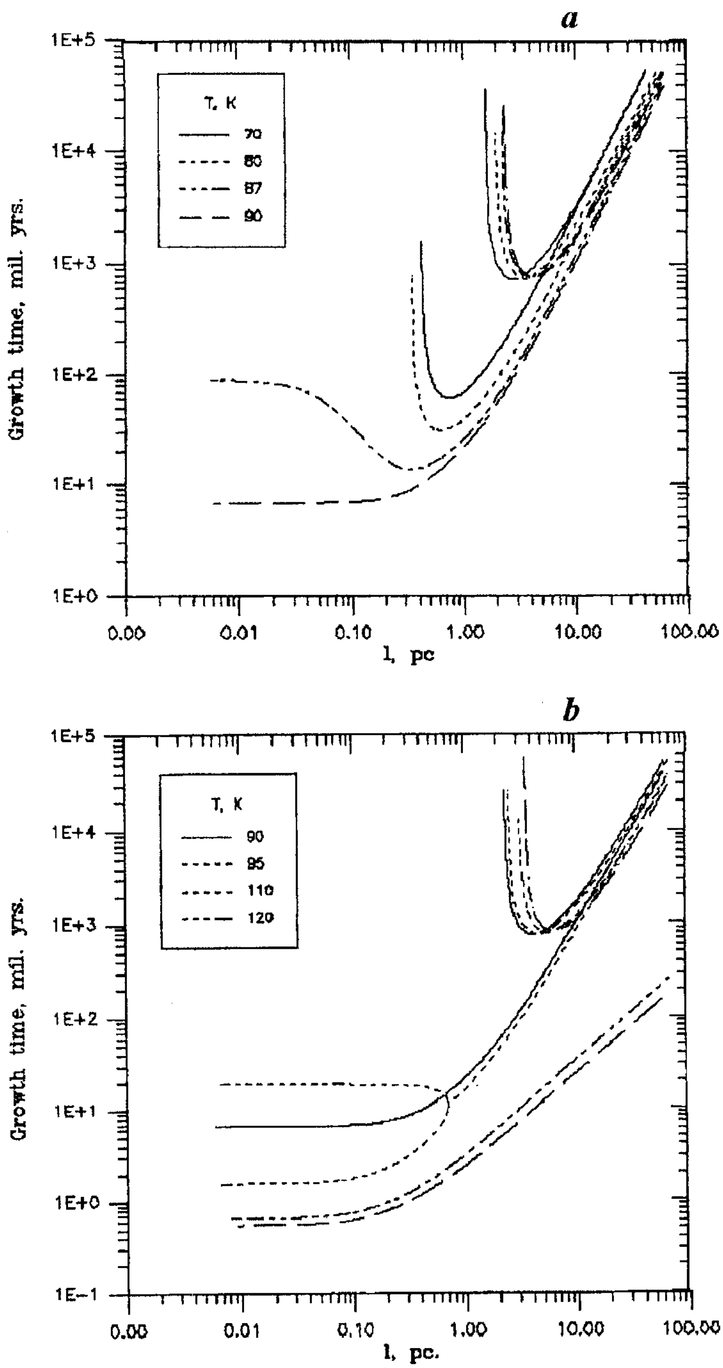

Figure 1 The dependence of the growth time of unstable modes on the wavelength for different temperatures.

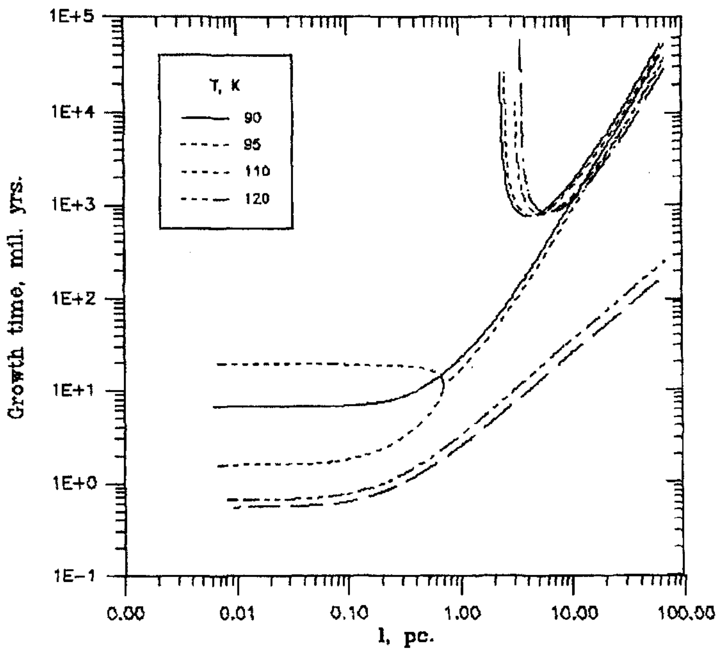

Figure 2 The typical dependence of the growth time on the wavelength at $T = 80$ K for different angle ϕ between the wave vector **K** and the magnetic field vector **b** (the case of slow magnetoacoustic instability).

oscillatory. At $T > 100$ K the instability of slow magnetosonic waves becomes purely dynamical. The effect of the magnetic field on this instability is to decrease its growth rate at the increasing angle ϕ between the wave vector and the magnetic field vector. The suppression is greatest at the angle $\phi = \pi/2$.

4 CONCLUSION

It has been shown in this paper that the maximum growth rate of the slow magnetosonic instability depends on the relative orientation of the wavevector to the magnetic field vector. The wavelength of the most rapidly growing perturbation is greatest when these vectors are parallel and is 1 pc by an order of magnitude. The transverse dimensions of the condensations are an order of magnitude smaller. The formed filamentary condensations can then be gravitationally contracted along the magnetic field lines and become oblate spheroids.

It should be noted that the mass of magnetosonic perturbations with maximal increment is close to the mediate clump mass in molecular clouds that contain $\approx 100M_{\odot}$ with spread from $10M_{\odot}$ to $1000M_{\odot}$. The typical density of clumps is $n \approx 10^3$ cm^3; the typical temperature $T \approx 40$ K (see Blitz, 1993). The clumps have a 3D structure similar to an oblate ellipsoid. In the case of L1688 the minor axis is directed along the magnetic field lines (Goodman and Heiles, 1996)

It is interesting to note that the form of the magnetosonic condensed clumps is close to thick disks connected with class O protostars (see Andre, 1995). The minor axes of this type of disk are directed along magnetic field lines. It is very attractive from our point of view to consider that magnetosonic instability causes the formation of an oblate spheroid which looks like a thick disk. It is possible that we have this type of situation in the case of young stellar objects (YSO) of class "zero" (see Pudritz *et al.*, 1996). It is usually supposed that this type of object is similar to other types of YSOs. In our proposition the oblate thick disk is first formed by the ionizational instability, and then a protostar is formed in this disk.

It is very important to understand that the instability of fast magnetosonic waves implies that superalfvenic turbulence cannot exist in the ISM. Any cloud turbulence in ISM must therefore be subalfvenic, i.e. practically non-compressible MHD turbulence without intermittency and fractal structure (see Dudorov, 1991).

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