MAGNETIC FIELD–RELATED HEATING INSTABILITIES
IN THE SURFACE LAYERS OF THE SUN AND STARS

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ABSTRACT

We investigate the stability of a magnetized low-density plasma to current-driven filamentation instabilities and apply our results to the surface layers of stars. Unlike previous studies, the initial (i.e., “precoronal”) state of the stellar surface atmosphere is taken to be a low-density, optically thin magnetized plasma in radiative equilibrium. Our (linear) analysis shows that the surface layers of main-sequence stars (including the Sun) which are “threaded” by magnetic fields are unstable; the instabilities considered here lead to structuring perpendicular to the ambient magnetic fields. These results suggest that relatively modest surface motions, in conjunction with the presence of magnetic fields, suffice to account for the presence of inhomogeneous chromospheric and coronal plasma overlying a star’s surface.

Subject headings: hydromagnetics — plasmas — stars: chromospheres — stars: coronae

I. INTRODUCTION

Our understanding of the physics which underlies the observed highly inhomogeneous state of the solar corona is still very incomplete. If one begins with the observed highly inhomogeneous distribution of magnetic flux on the Sun’s surface as a given condition (and so postpones discussion of a problem which is in its own right unresolved), one must explain how this inhomogeneity survives into the corona: straightforward model potential field calculations show that magnetic field lines which come from network flux elements lose “memory” of the inhomogeneous photospheric flux distribution within ~ 2000 km of the surface (see Gabriel 1976; Chapman 1981). It is therefore commonly supposed that the observed structuring is the result of some instability. Such analyses of the thermal and dynamic stability of stellar atmospheres usually assume that the (unperturbed) equilibrium state of the atmosphere is homogeneous and hot (see Heyvaerts 1974a); that is, that a preexisting coronal state is the appropriate initial condition. In this paper we examine an alternative possibility: we assume that the observed structured solar atmosphere represents the evolved state of a homogeneous atmosphere which is initially cool, magnetized, and in radiative equilibrium with the photospheric photon flux; and we examine the stability properties of such a cool magnetized plasma subject to various classes of “mechanical” heating processes. Our initial state is thus just the radiative equilibrium atmosphere derived by stellar structure calculations; and we regard nonradiative heating as a perturbation which may lead to both atmospheric inhomogeneity and a (hot) coronal state. The aim in our paper is not to describe a specific atmosphere, but is rather to show via a model calculation that the presence of magnetic fields in the surface layers of stars inevitably leads to the formation of an inhomogeneous and hot extended atmosphere.

The specific model problem we address involves a stellar photosphere in which magnetic fields are present (Fig. 1). We assume that the initial (hydrostatic) state of the atmosphere is obtained by solving the transport equations for the medium in the absence of mechanical heating processes, so that the plasma temperature and density are monotonically decreasing functions of radius. If emerged magnetic fields with reentrant field lines are present, then any photospheric motions will excite MHD modes (see Kuperus, Jonson, and Spicer 1981), as well as drive almost field-aligned currents along reentrant field lines (see Heyvaerts 1974b; Rosner et al. 1978); we examine the stability of this atmosphere and show that in the presence of magnetic fields, the atmosphere evolves toward a hot, filamented state.

The detailed model calculations, leading to the derivation of a dispersion relation are given in § II; we apply our calculation to our atmosphere prototype in § III and summarize our conclusions in § IV.

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HEATING INSTABILITIES

II. INSTABILITY ANALYSIS

a) Equilibrium State and Equations of Motion

Our initial cool plasma configuration is in radiative and hydrostatic equilibrium. We assume that the plasma is threaded by magnetic fields which are anchored externally to the medium under study; we shall demonstrate that this magnetized atmosphere is not stable, but is rather subject to MHD dissipative modes which lead to its heating and, therefore, to the formation of a hot "coronal" state. We describe these processes by the standard set of MHD equations, including the energy equation. The gas is assumed to be perfect and partially ionized. The electrical conductivity is considered isotropic; but anisotropy of the thermal conductivity is taken into account because its effect does play a major role.

The set of equations used is given by the standard MHD equations:

\[ \frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \]  
\[ \rho \frac{\partial \mathbf{v}}{\partial t} = -\nabla p + \frac{1}{c} \mathbf{J} \times \mathbf{B} - \rho \nabla \phi, \]  
\[ \frac{\partial p}{\partial t} - \gamma \frac{p}{\rho} \frac{\partial \rho}{\partial t} = (\gamma - 1) \left[ \mathbf{J} \cdot \left( \mathbf{E} + \frac{1}{c} \mathbf{v} \times \mathbf{B} \right) + \nabla \cdot (\kappa \cdot \nabla T) - E_\delta(p, T) + H(p, T) \right], \]  
\[ \nabla \times \mathbf{B} = \frac{4\pi}{c} \mathbf{J}, \]  
\[ \nabla \cdot \mathbf{B} = 0, \]  
\[ \nabla \times \mathbf{E} = -\frac{1}{c} \frac{\partial \mathbf{B}}{\partial t}, \]  
\[ \mathbf{J} = \sigma \left( \mathbf{E} + \frac{1}{c} \mathbf{v} \times \mathbf{B} \right), \]

where \( \rho(=nm) \), \( p \), \( T \), \( \gamma \), and \( \mathbf{v} \) are the plasma density, pressure, temperature, ratio of specific heats, and velocity,
respectively, and the electromagnetic quantities have their usual meaning. In the energy equation (2.1c), the first right-hand term is the Joule heating term; the second represents thermal conduction losses ($\kappa$ is the conductivity tensor); the third, radiative losses; and the last, other possible heating (in particular, radiative heating and MHD wave damping). The two last terms are expressed phenomenologically: the effective radiative losses are assumed to be expressible in the separable form $n^2 P(T)$, where the temperature-dependent term $P(T)$ is given (and discussed) in § III. The heating term can also depend on the local opacity; but, for simplicity, we examine stability in the very small optical depth regime; this restriction is further discussed below. We use the classical electric conductivity, given by Spitzer (1962):

$$\sigma = 9.8 \times 10^{12} T_e^{3/2} \xi^{-1} \text{s}^{-1}, \quad (2.2)$$

with

$$\xi = 1 + 0.305 \log T_e - 0.102 \log n_{e,10}$$

for a completely ionized plasma, and

$$\sigma = 3.69 \times 10^{14} T_e^{3/2} \text{s}^{-1} \quad (2.3)$$

for a weakly ionized plasma (where we follow the notational convention $T_e = T/10^4$ K, $n_{10} = n/10^{10}$ cm$^{-3}$, and similarly for the other variables). The thermal conductivity is anisotropic with respect to the local magnetic field; we use the expressions (see Lang 1978):

$$\kappa_{||} = 1.77 \times 10^5 T_e^{5/2} \xi^{-1} \text{ergs cm}^{-3} \text{s}^{-1} \text{K}^{-1},$$

$$\kappa_{\perp} = 1.70 \times 10^{-2} \xi^{-1/2} B_2^{-2/5} \text{ergs cm}^{-3} \text{s}^{-1} \text{K}^{-1} \quad (2.4a)$$

for a fully ionized plasma, and

$$\kappa_{||} = 1.18 \times 10^6 T_e^{1/2} \text{ergs cm}^{-3} \text{s}^{-1} \text{K}^{-1},$$

$$\kappa_{\perp} = 1.70 \times 10^{-2} X^2 n_{i,10}^{-1/2} T_e^{-1/2} B_2^{-2/5} \text{ergs cm}^{-3} \text{s}^{-1} \text{K}^{-1} \quad (2.4b)$$

for a weakly ionized plasma, where $X$ is the degree of ionization, and $n_i$, the total density of atoms and ions (in the following we assume a pure hydrogen plasma for simplicity, so that $n_i = n_{H} + n_{H^+}$).

The equilibrium state of the atmosphere is assumed to be static and isothermal ($v = 0$, $T = \text{constant}$) with a current flowing along the uniform magnetic field. Consistency with Ampere's law requires that the scale lengths in the problem be smaller than the field gradient scale:

$$L \ll L_0 \approx \frac{c B_0}{4 \pi j_0} \quad (2.5)$$

Finally, we assume that, in equilibrium, the net radiative heating rate is balanced by the divergence of the outward conductive flux:

$$h(\rho, T) = H(\rho, T) - E_R(\rho, T) = -\text{div} \kappa \nabla T. \quad (2.6)$$

### b) Dispersion Relation

A detailed study of the local dispersion relations for MHD thermally unstable modes was carried out by Heyvaerts (1974a); related discussions can be found in Kadomtsev (1966) and, more recently, in Biskamp and Horton (1975). Following the same procedure of analysis in the linear regime, and taking into account the presence of gravity, we derive a somewhat more general dispersion relation and then restrict our attention to the so-called Joule modes discussed by Heyvaerts (1974a); these lead to the possibility of current filamentation and enhanced Joule dissipation.

We Fourier-analyze all perturbed quantities:

$$f(r, t) = \int \frac{d^3k}{(2\pi)^3} \int \frac{d\omega}{2\pi} f(k, \omega) e^{i(\omega t - k \cdot r)};$$
thus, unstable modes have \( \text{Im } \omega < 0 \). We also choose \( \mathbf{k} \) in the \((x,z)\)-plane and \( \mathbf{B} \) along the \( z \)-axis. We linearize equations (2.1) and reduce them to a linear algebraic system in terms of the velocity components. The dispersion relation is obtained by setting the determinant equal to zero. We then define the following frequencies:

\[
\omega_s = \frac{c^2 k^2}{4 \pi \sigma_0}, \quad \omega_r = \frac{J^2}{p_0 \sigma_0}, \quad \omega_c = (\mathbf{k} \cdot \mathbf{k}) \frac{T_0}{p_0},
\]

\[
\omega_T = -\left\{ \frac{\partial}{\partial T} \left[ H(\rho,T) - E_R(\rho,T) \right] \right\} \frac{T_0}{p_0 \sigma_0}, \quad \omega_p = -\left\{ \frac{\partial}{\partial p} \left[ H(\rho,T) - E_R(\rho,T) \right] \right\} \frac{T_0}{p_0 \sigma_0},
\]

\[
\omega_1 = (\gamma - 1) \omega_c + \omega_T + \omega_f \left( \frac{\partial \log \sigma}{\partial \log T} \right)_0, \quad \omega_2 = -\frac{\gamma - 1}{\gamma} \left[ \omega_c - \omega_p + \omega_T + \omega_f \left( \frac{\partial \log \sigma}{\partial \log T} \right)_0 \right]
\]

(a subscript zero indicates a quantity evaluated at equilibrium). In addition, the following dimensionless quantities are also used:

\[
A = 2(\gamma - 1) \frac{\omega_f \frac{c^2 k^2}{\omega} \left( \frac{\partial \log \sigma}{\partial \log T} \right)_0}{4 \pi \sigma_0},
\]

\[
\Delta = \frac{\omega - i \omega_S}{\omega + i \omega_1}, \quad D = 1 - i \frac{\omega_S}{\omega}, \quad f = \frac{1 + i (\omega_1/\omega)}{1 + i (\omega_2/\omega)}, \quad S = \frac{D + f \Delta}{D + \Delta}.
\]

The general dispersion relation expanded in terms of \( k_0/k \) at zero order has two branches. The first branch allows only \( v_y = 0 \), while \( v_x \neq v_z \neq 0 \):

\[
\omega^2 = 0, \quad 1 + i \frac{\omega_1}{\omega} = 0, \quad \omega^2 (D + \Delta) - c_s^2 k_1^2 = 0.
\]

The second branch allows \( v_x = 0, v_z = 0 \), but \( v_y \neq 0 \):

\[
\omega^2 = 0, \quad 1 + i \frac{\omega_1}{\omega} = 0, \quad \Delta + D = 0, \quad D = 0,
\]

\[
\omega^4 (\omega + i \omega_1) - \omega^2 k^2 \left[ \frac{c_T^4}{D} (\omega + i \omega_1) + c_s^2 S (\omega + i \omega_2) \right] + \frac{c_A^2 c_T^2}{D} S k^2 k_\parallel^2 (\omega + i \omega_2) - i g k_\parallel \left( \omega^2 - c_A^2 k^2 \right) (\omega + i \omega_1) = 0.
\]

The general behavior of these modes follows the same lines proposed by Heyvaerts (1974a), although the dispersion relation derived here is somewhat more general (and corrects several errors in signs); from his discussion, we note that the most interesting modes in our context are the so-called Joule mode and antidiffusion mode:

\[
(D + \Delta) - \frac{c_s^2 k_\parallel^2}{\omega^2} = 0 \quad \text{(first branch)},
\]

\[
(D + \Delta) = 0 \quad \text{(second branch)}.
\]

Actually, these two modes are physically very similar and, since we shall later use \( k_\parallel \to 0 \), any polarization will be allowed. Written explicitly, this mode becomes:

\[
1 - i \frac{\omega_1}{\omega} + 2(\gamma - 1) \frac{\omega_f}{\omega (\omega + i \omega_1)} \frac{c_s^2 k_1^2}{4 \pi \sigma_0} \left( \frac{\partial \log \sigma}{\partial \log T} \right)_0 = 0.
\]
In the limit $\omega_c \gg \omega$, corresponding to short diffusion time scales (field not frozen in the plasma), one obtains the simpler relation:

$$\omega = -i(\gamma - 1) \left[ -\left( \kappa_i k_i^2 + \kappa_\perp k_\perp^2 \right) \frac{T_0}{p_0} + \frac{T_0}{p_0} \left( \frac{d h}{d T} \right)_0 + \frac{J_0^2}{\sigma_0 p_0} \left( \frac{d \log \sigma}{d \log T} \right)_0 \left( \sin^2 \theta - \cos^2 \theta \right) \right], \quad (2.13)$$

where $\theta$ is the angle between $k$ and $B$. For instability, the imaginary part of the frequency must be negative; this means that the last term within the square brackets must be positive and exceed the two preceding terms in absolute value. As $\theta$ is the angle of the mode's wave vector with respect to the magnetic field (and currents), unstable modes must be essentially transverse to the field, i.e., $\sin^2 \theta > \cos^2 \theta$. In addition, a numerical evaluation of the conduction term (see eq. [2.4]) shows that parallel conduction damps longitudinal modes most effectively; we shall therefore require that $k_i \sim 0$ and consider only purely transverse modes, for which thermal conduction is impeded by the presence of the magnetic field. We thus obtain as the general condition for instability

$$\left( \frac{J_0^2}{\sigma_0} \right) \left[ \frac{d \log \sigma}{d T} \right]_0 > -\left[ \frac{d h}{d T} \right]_0. \quad (2.14)$$

Note that, under the conditions of the "precoronal" equilibrium atmosphere discussed in § III, thermal conduction may be neglected altogether; in that case, longitudinal modes may be unstable as well. In the regime in which local current heating is predominantly balanced by radiative losses, we have instead of equation (2.14) [using the relation for energy balance at equilibrium, i.e., $J_0^2/\sigma_0 \sim \kappa_0 P(T_0)$]:

$$\left( \frac{d \log \sigma}{d \log T} \right)_0 > \left( \frac{d \log P}{d \log T} \right)_0. \quad (2.15)$$

It is crucial to note here that, because our analysis is local, it cannot answer the question whether the inhomogeneities in temperature that develop along $B$ ultimately can suppress the instability.

From the expressions for the conductivity (2.2) and (2.3), we find that Joule modes are excited when the temperature dependence of the net local radiation losses is less steep than $T^{3/2}$. In § III, we discuss one astrophysical context in which this condition obtains. Here we point out that the transverse Joule modes perturb the current distribution perpendicular to the flow; this can be interpreted as the linear stage of current filamentation. A limit to the size of these filaments can be obtained from an order of magnitude estimate based on the Maxwell equations:

$$\Delta x \lesssim \frac{c}{4\pi} \frac{\Delta B}{J} \sim \frac{c}{4\pi} \frac{B_0}{J_0} = \frac{1}{k_0}, \quad (2.16)$$

which coincides with the consistency condition of our model equations.

Finally, we note that the magnetic field and current buildup in filaments is not hindered in the linear regime by pressure forces as long as the plasma is not frozen in on the time scale $\omega_c^{-1}$ of the instability, i.e., such that $\omega_c \gg \omega$; this constraint imposes an upper bound on the instability wavelength of $\lambda \sim (4\pi p_0)^{1/2} c/J_0$. In addition, as just discussed, classical thermal conduction between adjacent filaments is insufficient to smear out local temperature gradients perpendicular to $B$ (although thermal conduction effects acting along $B$ in the nonlinear limit cannot be excluded). A further mechanism which may saturate the instability is increased current dissipation because of enhanced electron scattering by plasma turbulence; such dissipative processes limit the filamentation process by increasing the diffusion rates (which will tend to broaden spatially confined currents). The instability with lowest threshold commonly considered is ion-cyclotron turbulence (see Papadopoulos 1977 and references therein; more recently, see Hinata 1980 and Morrison and Ionson 1981). As long as the ratio of the ion gyrofrequency $\Omega_i$ to the ion plasma frequency $\omega_{pi}$ is small (e.g., $n \approx 5.4 \times 10^3 B_2^2$ cm$^{-3}$), one finds that, for $T_e/T_i \sim O(1)$, the enhanced resistivity is (Ionson 1976)

$$\eta_i \approx 0.06 (\Omega_i/\omega_{pi}) (1 - 12 v_i/v_D) \omega_{pe}^{-1}, \quad (2.17)$$

where $v_D$ is the current drift velocity and $v_i$ is the ion thermal speed, provided that $v_D \geq 13 v_i$. The turn-on of this instability acts precisely in the opposite direction to the Joule mode: whereas the classical resistivity decreases with increasing temperature in the current-carrying channels (and so provides positive feedback for further current channelization), ion-cyclotron turbulence acts precisely within the current-carrying channels (where the current drift speeds are largest) to increase the local resistivity and thus prevent further channelization. An estimate of the filament
III. APPLICATION OF INSTABILITY CALCULATION

The results of our stability analysis given in § II above are now applied to a model atmosphere which approximates conditions in a stellar photosphere. Unlike Heyvaerts (1974a), who studied the stability of (preexisting) chromospheric and coronal plasmas, we assume a “precoronal” atmosphere as our unperturbed state. This atmosphere is assumed to be in radiative and hydrostatic equilibrium, so that both temperature and density are monotonically decreasing functions of radius. We suppose that magnetic “loops” protrude into this atmosphere from the convection zone under the action of buoyancy, leading to a surface magnetic field whose field lines are largely “reentrant” (e.g., field lines traverse the stellar surface at least twice), and we further suppose that fluid motions at the stellar surface or below induce (almost force-free) current flow (see Heyvaerts 1974b; Rosner et al. 1978). The resulting ohmic heating is easily shown to be negligible near the field line “footpoints” (viz., in the region of unit optical depth in the continuum), where radiative heating and cooling predominate by far. However, the equilibrium atmosphere experiences an exponential density decrease characterized by very small scale heights; thus, sufficiently far out, it may no longer be true that ohmic heating is negligible. For example, in the solar case, the density scale height at 5000 K is \( \sim 2.5 \times 10^7 \) cm; in that case, the plasma densities within magnetic field structures of typical solar dimensions (\( \sim 10^9 - 10^{10} \) cm) immersed in a “precoronal” atmosphere would lie well below \( 10^5 \text{ cm}^{-3} \) throughout most of the structures, so that ohmic heating could easily compete with radiative heating (viz., photoionization) for temperatures below \( \sim 10^4 \) K. It is this (optically thin) region whose stability we investigate.

The dominant terms in the energy balance of an optically thin plasma for temperatures below \( \sim 10^4 \) K are heating by photoionization (see Osterbrock 1974) and cooling by continuum and line radiation; we approximate the latter by using the low-density approximation of Dalgarno and McCray (1972), given by

\[
E_R(n, T) \approx X n^2 \begin{cases} 
7.1 \times 10^{-29} T^{0.88} \text{ergs cm}^{-3} \text{s}^{-1} & (10^2 \text{ K} \leq T \leq 10^3 \text{ K}), \\
2.4 \times 10^{-27} T^{0.32} \text{ergs cm}^{-3} \text{s}^{-1} & (10^3 \text{ K} \leq T \leq 10^4 \text{ K}),
\end{cases}
\]

where \( X \) is the fractional ionization (\( n_e/n_H \)), \( T \) is the local plasma temperature, and where we have fit the results given by Dalgarno and McCray for \( X = 10^{-3} \) to a power law in the two temperature ranges shown (for \( X < 10^{-3} \), the dependence of \( E_R \) on \( X \) is weaker than linear, and becomes independent of \( X \) for \( X < 10^{-4} \)). At the temperatures of interest, thermal conduction is entirely negligible.

We may easily estimate the current density \( J_0 \) at which ohmic heating can be important by simply equating the ohmic heating gains with the radiative losses. Using the Spitzer conductivity, we obtain

\[
J_{\text{min}} \approx \begin{cases} 
0.029 X^{1/2} n_5 T_4^{1/2} \text{cgs} & (10^2 \text{ K} < T < 10^3 \text{ K}), \\
0.013 X^{1/2} n_5 T_4^{3/2} \text{cgs} & (10^3 \text{ K} < T < 10^4 \text{ K}).
\end{cases}
\]

These current density values are very small (e.g., “nonpotential” magnetic fields of only \( \sim 10^{-2} R_9 \) gauss suffice; \( R_9 \) is the loop mirror radius in units of \( 10^9 \) cm), confirming our supposition that ohmic heating can easily dominate other local heating processes. We note, upon comparing equation (3.1) with the conductivity's temperature dependence (\( \sigma \propto T^{3/2} \)), that relation (2.14) above predicts instability. Therefore we expect the transverse Joule mode to be operative; a lower bound to the (near) current filamentation time scale is \( t \sim \sigma / \nu \) (where \( \sigma \approx \sigma_0 J_0^2 \leq 1.3 \times 10^{-4} X_3 n_5 T_4^{5/2} \varepsilon B_2^{-1} T_5^{3/2} / v \text{s under solar conditions}^3 \) with a limiting filamentation scale length (from eq. [2.16]) of less than \( 2 \times 10^7 \varepsilon T_5^{3/2} / v \text{ cm.} \)

Within the resulting filaments, the increasing current density may then excite plasma instabilities. For example, ion-cyclotron electrostatic modes will be excited as \( v_D/v_B \rightarrow 1 \) (where \( v_B = J_0/\rho_n e \)); the typical time scale of

\[v \sim 1 \text{ cm s}^{-1}.\]

\[^3\text{We estimate an upper bound to} J_0 \text{ by adopting the analysis of Heyvaerts (1974b); assuming canonical photospheric plasma parameters of} T = 5000 \text{ K,} n = 10^{16} \text{ cm}^{-3}, \text{ and} B = 10^4 \text{ gauss, we find that} J_0 \approx 2 \times 10^4 \varepsilon^{-1} \frac{1}{\epsilon} (B/10^2 \text{ gauss}) (T/5000 \text{ K})^{5/2} (v/1 \text{ cm s}^{-1}) \text{ cgs where} \epsilon \sim O(1) \text{ is the ratio of the horizontal to vertical scale lengths at the photospheric field line footpoints, and} \xi = O(1) \text{ is a very weak (logarithmic) function of} n \text{ and} T. \text{ Note that the amplitude of the required surface motion is quite modest} (v \sim 1 \text{ cm s}^{-1}).\]
growth for this process is \( r \approx \Omega^{-1} \sim 10^{-5} R_\odot^{-1} \) s and can thus occur with rapid onset. The process of filamentation can therefore saturate by the formation of a turbulent state of electrostatic fluctuations, as discussed above (§ II). It is further noteworthy that, under the assumed plasma conditions, the instability growth rate is somewhat smaller than the ion gyrofrequency, so that the transverse current heating effects discussed by Biskamp and Horton (1975) are probably absent.

Because the analysis of § II deals with the linear instability regime, we cannot predict the ultimate evolution of the Joule mode. However, our analysis does show that substantial plasma temperature and density inhomogeneities evolve within magnetically confined surface plasma structures and suggests that this inhomogeneity allows the onset of other quasi-stationary heating processes (such as intermittent tearing mode reconnection, see Galeev et al. 1981; surface wave dissipation, see Ionson 1978; fast-mode damping, see Habbal, Leer, and Holzer 1981) which require the preexistence of spatially confined currents or of temperature or density contrasts, or both, in order to produce hot coronal “loops.” For example, the Alfvén surface wave model of Ionson (1978) presupposes the existence of discontinuities in Alfvén speed at coronal heights (\( \gtrsim 0.01 R_\odot \)); however, as emphasized in § I, straightforward potential field extrapolations of inhomogeneous photospheric magnetic fields for a “precoronal” cool atmosphere show very uniform spatial structure at coronal heights. We have, however, shown here that such a uniform state is unstable and that the instability discussed here leads to a variation in Alfvén speed perpendicular to the ambient magnetic field, just the condition required for the efficient propagation and damping of surface modes. Similar arguments apply to the tearing mode heating model of Galeev et al. (1981) and the fast-mode damping model of Habbal and Leer (1981, see also references therein).

IV. DISCUSSION

We have investigated the local linear stability of magnetized plasma in the small optical depth regime, focusing particularly on the so-called Joule mode; our interest lies in answering the question under what conditions (if any) an initially diffuse current system can undergo filamentation (and thereby drive local thermal instability). Our results show that, under the conditions appropriate to a “precoronal” stellar atmosphere (i.e., above the surface of a star which is assumed to have neither chromosphere nor corona), the low-density atmosphere is unstable to filamentation. In fact, we have shown that the presence of surface magnetic fields, together with very modest surface flows (whose amplitude is far less than that associated with the vigorous surface convection of late-type dwarf stars), suffice to produce on time scales consistent with linear theory an inhomogeneous atmosphere characterized by relatively hot (almost force-free) current filaments. A variety of other plasma heating processes have been suggested which require an initially inhomogeneous atmosphere in order to function; it is likely that these processes will determine the ultimate plasma state of the atmosphere.

Because one expects surface magnetic fields in virtually all stars (either dynamo-produced or primordial fields), and because the required surface flow amplitudes are so modest (§ IIb), our results support the argument of Athay (1976) and others that the solar chromosphere and corona are manifestations of (thermal) instability, with the additional element that the instability is by nature filamentary and necessarily involves magnetic fields. In addition, we find that this instability is a general attribute of the outer layers of stars and hence not restricted to solar-type stars. Unfortunately, the present analysis does not allow determination of the large-amplitude behavior of the instability (which, as just mentioned, is likely to involve other heating processes), and so does not establish the ultimate (possibly stationary) state of the atmosphere; nevertheless, our results do suggest that the presence of hot and spatially structured X-ray emitting gas above a star’s surface is a phenomenon not restricted to solar-type stars alone, and thus supports the view that observations of X-ray emission from all main-sequence (as well as most other) stars (Vaiana et al. 1981) can be understood on the basis of the solar analogy.

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