Heating the Solar Atmosphere by the Self-Enhanced Thermal Waves Caused by the Dynamo Processes

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ABSTRACT

We discuss a possible mechanism for heating the solar atmosphere by the ensemble of thermal waves, generated by the photospheric dynamo and propagating upwards with increasing magnitudes. These waves are self-sustained and amplified due to the specific dependence of the efficiency of heat release by Ohmic dissipation on the ratio of the collisional to gyro-frequencies, which in its turn is determined by the temperature profile formed in the wave. In the case of sufficiently strong driving, such a mechanism can increase the plasma temperature by a few times, i.e. it may be responsible for heating the chromosphere and the base of the transition region.

Subject headings: Magnetohydrodynamics — Sun: chromosphere — Sun: transition region

1. INTRODUCTION

Seeking for the mechanism of heating the solar atmosphere with height from a few thousand to a million Kelvin is one of long-standing problems in astrophysics. The approaches proposed by now, roughly speaking, can be separated into two groups (e.g. reviews by Walsh & Ireland 2003; Erdélyi & Ballai 2007, and references therein). The first group deals with a generation of some kinds of (magneto-) hydrodynamic waves or pulses in the base of the solar atmosphere and their subsequent propagation and dissipation in the upper layers. The mechanisms of the second group assume that the heating is due to the ensemble of small-scale flare-like events, caused by the reconnection processes in the specific magnetic field configurations. It is commonly believed now that no unique mechanism can provide the entire heating of the solar atmosphere, and a few of them are acting simultaneously in the Sun.

Besides, as was proposed by Aschwanden et al. (2007), the inverted (increasing with height) temperature profile might be formed “dynamically”— due to the fluxes of plasma heated in the lower layers and propagating upwards (which is often called evaporation). This is, in fact, the third kind of the heating mechanisms.

The aim of the present work is to consider one more mechanism, which, strictly speaking, belongs to none of the three above-mentioned groups but inherits some features from all of them. Like mechanisms from group II, it assumes an ensemble of the small-scale flare-like events; but such events require no special field topology for their occurrence (as in the case of reconnection) and are more similar to the specific type of wave pulses, as in the mechanisms of group I. At last, our scenario reminds the Aschwanden’s approach (group III), because the heating process propagates dynamically from the lower to upper layers.

2. FORMULATION OF THE MODEL

The dynamo processes in weakly-ionized plasmas, considered in this section, closely remind the so-called $S_q$-dynamo, well known in the ionospheric physics (e.g. review by Richmond 1989).
Fig. 1.— Sketch of the dynamo processes in the partially-ionized plasma (motion of the neutrals, electrons, and ions are colored green, blue, and red, respectively, in the online version of this figure) and a schematic structure of the associated thermal wave, propagating upwards (colored magenta in the online version).

and references therein). Similar processes in the solar physics are called photospheric dynamo, and they were employed in a number of papers to describe generation of field-aligned currents and perturbations of the magnetic fields (Hénoux & Somov 1987). The associated thermal effects were usually neglected; an exception is the old work by Sen & White (1972).

2.1. Qualitative Description

Let us consider a column of partially-ionized plasma stratified by the gravitational field \( g \) and permeated by the magnetic field \( B \), which for simplicity is taken to be vertical and constant. Besides, we assume that this plasma is driven radially in the horizontal plane by the motion of a neutral component \( v_n \) (which can be associated, for example, with convective flows at the photospheric level); see Figure 1.

Due to the gravitational stratification, a neutral gas density sharply decreases with height. Respectively, the free-path length and the time of inter-particle collisions sharply increase, while the collisional frequency decreases. Then, at the bottom of the column, where the frequencies of collisions of both electrons and ions with neutrals, \( \nu_{en} \) and \( \nu_{in} \) respectively, are much greater than their gyrofrequencies \( \omega_e \) and \( \omega_i \), the charged particles are completely dragged by the neutrals and move with them; so that no local electric fields and currents are generated.

In the opposite case, at the top of the column, where the collisional frequencies are much less than the gyrofrequencies, the electrons and ions are unaffected by the radial motion of the neutral component and, therefore, will remain at rest. No electric fields and currents are produced again.

The most nontrivial situation arises at the intermediate heights, where \( \nu_{in} \approx \omega_i \) and \( \nu_{en} < \omega_e \). Under these circumstances, the ions are dragged by neutrals, while the electrons remain approximately at rest. Due to the separation of the electrons and ions, a radial electric field \( E \) should develop, resulting in its turn in the azimuthal drift of the charged particles. However, while the electrons will experience such a drift almost without resistance, the ions will be decelerated by the collisions with neutrals. Therefore, the azimuthal (Hall) electric currents \( j \) are generated, as shown in Figure 1.

The above-mentioned electric fields and currents are concentrated in a quite narrow height range, where \( \nu_{in} \approx \omega_i \); and they should lead, firstly, to the magnetic-field perturbations (considered by Hénoux & Somov 1987) and, secondly, to the dissipative effects and heat release, which are just the subject of our study. (The dissipation exists due to the azimuthal electric field appearing in the coordinate frame co-moving with plasma.)

Once a heat release has occurred at the height where \( \nu_{in} \approx \omega_i \) (we shall call it \( z_{eff} \)), this spot is no longer efficient for the subsequent heat release, because the plasma is heated and the thermal velocity of its particles \( v_i \) as well as the collisional frequency \( \nu_{in} = v_i/\lambda_{in} \) increased (under assumption of time-independent neutral density \( n_n \), constant ion-neutral cross-section \( \sigma_{in} \) and, consequently,
the unchanged mean free path $\lambda_{in} = 1/(\sigma_{in} n_i)$. Therefore, the condition $\nu_{in} \approx \omega_i$ will no longer be satisfied.

On the other hand, due to the heat conduction in space, the ion thermal velocity $v_i$ will increase also at a greater height (where the mean free path $\lambda_{in}$ is larger due to the initial gas stratification) and, therefore, the spot of effective heat release, $\nu_{in} \equiv v_i/\lambda_{in} \approx \omega_i$, will move there. As a result, we should expect formation of the specific self-sustained thermal wave propagating upwards, as depicted in the right-hand panel of Figure 1. Moreover, the wave amplitude will increase with height, just because of decreasing a heat capacity per unit volume in the stratified gas.

It should be also mentioned that heat release in the dynamo-region may be further facilitated by the onset of two-stream instability (Sen & White 1972 and references therein); but we shall not consider here in more detail the corresponding microphysical mechanisms.

Finally, let us discuss more carefully the scope of applicability of our scenario. In fact, the ion’s collisional frequency is composed of two parts: $\nu_i = \nu_{in} + \nu_{ii}$, where $\nu_{in}$ and $\nu_{ii}$ are the frequencies of ion-neutral and ion-ion collisions. It was assumed everywhere above that the second term can be ignored. On the other hand, the degree of ionization increases with height; so that the ion-ion collisions will inevitably become dominant starting from some altitude. In this case, $\lambda_{ii} = 1/(\sigma_{ii} n_i) \sim v_i^4/n_i$, since $\sigma_{ii} \sim v_i^{-4}$. So, the ion’s collisional frequency $\nu_{ii} = v_i/\lambda_{ii} \sim n_i/v_i^3$ decreases with temperature, as distinct from the case of ion-neutral collisions. However, this fact is unrelated to the criterion of the dynamo-layer development, $\nu_{in} \sim \omega_i$, because collisions between the charged particles of the same kind cannot affect the generation of electric currents. The dominant role of ion-ion collisions can change only the coefficient of heat conductivity.

Therefore, the scope of applicability of the presented model is somewhat wider than seems at the first sight. At the same time, the plasma must be weakly ionized in the entire height range under consideration, since otherwise it would be meaningless to speak about dragging the charged particles by the neutral gas.

### 2.2. Basic Equations

In the simplest one-dimensional approximation, the process of heat transfer along the magnetic flux tube can be described by the equation:

$$\frac{\partial}{\partial t}[c n T] - \frac{\partial}{\partial z} \left[ c n \chi \frac{\partial}{\partial z} T \right] = q, \quad (1)$$

where $T$ is the temperature of the heavy particles, i.e., neutrals and ions (the electron temperature is not of interest here), $t$ is the time, $z$ is the vertical coordinate, $n$ is the number density of the heavy particles, $\chi \approx (1/3) \lambda v$ is the temperature conductivity, $c = \text{const}$ is the heat capacity per a heavy particle, and $q$ is the volume density of the heat release due to Ohmic dissipation. (The basic parameters of neutrals and ions are assumed to be approximately the same and, therefore, written without additional subscripts.)

Leaving aside a microscopic theory of the Ohmic heating, let us use a phenomenological approximation of the heat-release profile, based on its qualitative behavior discussed above. Namely, we take Taylor expansion in the vicinity of the heat-release maximum, achieved at $\nu_{in} \approx \omega_i$:

$$q(z,T) = q_0 \left\{ 1 - \alpha \left( 1 - \frac{\nu_{in}(z,T)}{\omega_i} \right)^2 + \ldots \right\}, \quad (2)$$

where $q_0$ is the amplitude of the heat release, and $\alpha$ is the characteristic width of the profile.

Next, we shall use for simplicity the time-independent exponential (i.e., actually, isothermal) height profile of the gas density:

$$n(z) = n_0 \exp[-z/H], \quad (3)$$

where $H = T_0/mg \equiv \text{const}$ is the height scale. (In other words, it is assumed that the initial gas profile is not changed during the heat propagation; of course, redistribution of the gas density should be taken into account in a more accurate treatment.)

The thermal velocity is evidently related to the temperature as

$$v = v_0 \sqrt{T/T_0}. \quad (4)$$

So, assuming that the cross section $\sigma$ is constant (which is a reasonable approximation for the collisions of ions with neutrals and neutrals with each other), we get:

$$\nu_{in} = \sigma n v_0 \sqrt{T/T_0} = \nu_{in0} (n/n_0) \sqrt{T/T_0}. \quad (5)$$
Besides, it is convenient to choose the origin of z-axis (z = 0) in the point where \( \nu_{in0} = \omega_i \).

At last, substituting expressions (2)–(5) into equation (1), we obtain the equation of heat transfer in the following form:

\[
e^{-z^*H^*} \frac{\partial T^*}{\partial t^*} - \frac{1}{3} \frac{\partial}{\partial z^*} \left[ T^{*+1/2} \frac{\partial T^*}{\partial z^*} \right] = q_0^* \left\{ 1 - \alpha \left( 1 - e^{-z^*/H^*} T^{*+1/2} \right)^2 + \ldots \right\},
\]

where the dimensionless quantities (marked by asterisks) were introduced as

\[
T^* = T/T_0, \quad t^* = t/\tau_0, \quad z^* = z/\lambda_0, \quad H^* = H/\lambda_0, \quad q_0^* = \frac{q_0 \tau_0}{c \rho_0 T_0},
\]

(7)

(8)

(9)

(10)

\( T_0 \) is the initial temperature of the gas, \( \tau_0 \equiv \nu_{in0}^{-1} \) and \( \lambda_0 \equiv \lambda_{in0} \).

3. RESULTS OF THE COMPUTATION

Expression (6) represents a quite specific type of the nonlinear parabolic partial differential equation. A convenient tool for solving just this type of problems was implemented in the NAG numerical library—this is the subroutine D03PCF (Numerical Algorithms Group 2001; Berzins 1990).

We restrict our consideration here by the simplest numerical solution, just to illustrate how the proposed mechanism works. (A detailed treatment of more realistic situations will be published elsewhere.) So, let us take the initial condition in the form:

\[
T^* = 1 \quad \text{at} \quad t^* = 0
\]

(i.e. the entire gas had initially the same temperature \( T_0 \)).

The boundary conditions will be specified, somewhat arbitrarily, as vanishing heat fluxes at the top and bottom of the computational region:

\[
\frac{\partial T^*}{\partial z^*} = 0 \quad \text{at} \quad z^* = z_d^* \quad \text{and} \quad z^* = z_u^*.
\]

(10)

(However, we shall discuss below only behaviour of the solution far from the boundaries, which is insensitive to the boundary conditions.)

At last, let us take, for example, the following parameters of the problem: \( H^* = 10, \quad q_0^* = 0.1, \quad \alpha = 1, \quad z_d^* = -100, \quad \text{and} \quad z_u^* = 100 \).

The obtained numerical solution is presented in Figure 2 where height profiles of the temperature are drawn for the successive instants of time with increments \( \Delta t^* = 10 \).

As is seen, a propagation of the self-sustained thermal wave really results in the formation of temperature profiles monotonically increasing with height above the spot of the initial heat release at \( z = 0 \). This is essentially related to the dynamical nature of the process; otherwise, at the fixed heating source, the temperature profile would be strongly asymmetric (because of the substantial height dependence of the heat conductivity), but its maximum would always take place at the fixed site of the heat release.

Next, the series of profiles in Figure 2 monotonically increases with time, without any clear evidence for saturation. This is just because our simplified model does not take into account the radiative loss as well as exhausting efficiency of the dynamo processes in the course of time. It would be desirable, of course, to include into the future numerical simulations a realistic function of radiative loss (which should be especially important at the enhanced degree of ionization, associated with the
increased plasma temperature). However, the role of radiative cooling should not be overestimated in advance: when the plasma temperature is well below the ionization thresholds of the plasma elements, then the establishment of higher ionization density requires a very large number of interparticle collisions and, therefore, may take a considerable time. So, if the discussed thermal waves are impulsive phenomena, then there may be just insufficient time for the increase in ionization, and the radiative losses will be not so large.

4. DISCUSSION AND CONCLUSIONS

As follows from the above consideration, the presented model may be a mechanism for heating the lower part of the solar atmosphere, i.e. formation of the temperature profile increasing with height. At the same time, it should be borne in mind that it works only in weakly-ionized plasma. To provide plasma heating up to the coronal temperatures, it is necessary to invoke other heat-release mechanisms, commonly based on the magnetic reconnection (for the review of recent results, see Shibata et al. 2007; Aulanier et al. 2007).

Let us mention also that some qualitative ideas why heating of the solar plasma should occur just at the spot of approximate equality between the collisional and gyro- frequencies were put forward in our paper about a decade ago (Dumin 2002). Unfortunately, it was not clear from that work how is it possible to get the temperature profile monotonically increasing with height rather than peaking at a fixed altitude of the maximum heat release. The mechanism of the propagating thermal wave, described in the present paper, gives a possible answer to that question.

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