A new paradigm for solar coronal heating

J. Vranjes (a) and S. Poedts

K. U. Leuven, Center for Plasma Astrophysics - Celestijnenlaan 200B, 3001 Leuven, Belgium, EU
and Leuven Mathematical Modeling and Computational Science Center (LMCC) - Leuven, Belgium, EU

received 31 March 2009; accepted 20 April 2009
published online 18 May 2009

PACS 96.60.-j – Solar physics
PACS 96.60.P – – Corona
PACS 52.35.Kt – Drift waves

Abstract – The solar coronal heating problem refers to the question why the temperature of the Sun’s corona is more than two orders of magnitude higher than that of its surface. Almost 70 years after the discovery, this puzzle is still one of the major challenges in astrophysics. The current basic paradigm of coronal heating is unable to explain all the observational features of heating. Here we argue that a new paradigm is required to solve the puzzle in a self-consistent manner. The alternative approach is based on the kinetic theory of drift waves. We show that, with qualitative and quantitative arguments, the drift waves have the potential to satisfy all coronal heating requirements.

Copyright © EPLA, 2009

A self-consistent coronal heating model must fulfill a lot of requirements imposed by observational facts. First of all, it should be consistent with the measured energy losses in the solar corona due to conduction and radiation, i.e., it should a) not only provide the right amount of energy but b) do so at the right time scales, e.g., \( \approx 10^{-4} \text{J/(m}^3\text{s)} \) in active regions. Moreover, it should c) include the source of the required energy [1], and d) work everywhere in the corona, i.e., for all different magnetic structures (with different heating requirements). Furthermore, it should be able e) to explain the observed temperature anisotropy \([2,3] (T_\perp > T_\parallel)\), f) be more effective on ions than on electrons \((T_i > T_e)\), and g) heat heavier ions more efficiently than lighter ions [4]. None of the proposed heating mechanisms so far even claimed to fulfill all these model requirements.

The current paradigm of coronal heating states that the required energy source is provided by the plasma flows below the solar surface and that this energy is transferred to the corona through the motions of the “footpoints” of the coronal magnetic field lines that are “anchored” in this zone. Depending on the ratio \( \tau \) of the time scale of these “driving” motions to the dynamic (Alfvén) time scale, the current models are classified as “wave heating” \((\tau < 1)\) or “magnetic reconnection (or nanoflare)” \((\tau > 1)\) models [5]. The main challenge, however, is to explain how the energy is dissipated in the highly conductive corona with a Lundquist number (the ratio of the dissipation time to the Alfvén time scale) around \(10^{13}\). All proposed mechanisms have a problem either with the energy transport to the corona (the observed wave fluxes are too small) or with the dissipation (too little and/or too slow, or only sporadic) [5].

Most current models rely on the continuum or fluid approximation (magnetohydrodynamics, MHD). However, these models cannot really explain coronal heating completely because i) it is clear that the actual heating takes place at length scales much smaller than those on which the (macroscopic) MHD model is justified; and ii) it is obvious that the observed discrepancy between ion and electron temperatures in the corona, as well as iii) the observed large temperature anisotropy in the inner corona \((T_\perp > T_\parallel)\), and iv) the observed preferential heating of the heavier ions [2] are beyond the (single!) fluid model.

Here, we make a starting step toward the formulation of a new paradigm, based on the kinetic theory of the drift waves driven by density gradients that are omnipresent in the solar corona. It implies that the direct energy supply for the heating comes from the corona itself (from the density gradients), though still maintained and replenished by some mechanisms below the surface. These include a continuous restructuring of the magnetic field, implying consequent similar changes of the plasma density (due to the frozen-in conditions), and also the observed inflow of the plasma along the magnetic
loops [6]. To some extent this looks similar to the currently accepted scenarios mentioned above, where the magnetic field plays an essential role and is assumed as given. However, the suggested new approach not only enables to describe these drift waves (which are missing in the MHD picture), but also their dissipation is easy to explain in the self-consistent kinetic model that works on the (very small) length scales at which the actual dissipation takes place. Actually, two mechanisms of energy exchange and heating will be shown to take place simultaneously, one due to the Landau effect in the direction parallel to the magnetic field, and another one, stochastic heating, in the perpendicular direction. Moreover, this stochastic drift wave heating mechanism seems to satisfy all seven above-mentioned model requirements. This will be proved below using only established basic theory, verified experimentally in laboratory plasmas.

The drift mode is the only mode that is not only able to survive the drastically different extremes in various parts of the solar atmosphere, but, in fact, even manages to benefit (i.e., to grow) from each of them. The driving mechanism, however, is always the same, viz. the density gradients perpendicular to the ambient magnetic field vector. Numerous observations confirm the omnipresence of such irregularities of the plasma density across magnetic flux surfaces (see fig. 1). Extremely fine density filaments and threads have been observed in the solar atmosphere for a long time now, even from ground-based observations [7] such as those during the eclipse in 1991, showing a slow radial enlargement of the structures. Contour maps [8] reveal the existence of numerous structures of various sizes. Filamentary structures with length scales of the order of 1 km have been discussed [9]. Recent Hinode observations [10] confirm that the solar atmosphere is a highly structured and very inhomogeneous system, with radially spreading density filaments of various sizes pervading the whole domain. A very recent three-dimensional analysis of coronal loops reveals short-scale density irregularities within each loop separately [11]. Such density irregularities are, as a rule, associated with the magnetic field, thus creating a perfect environment for drift waves. The characteristic dimensions of the observed density irregularities are limited by the resolution of the current instruments that is presently a fraction of an arcsec. However, even extremely short, meter-size length scales cannot be excluded, especially in the corona [12]. Hence, in dealing with drift waves, we may operate with density inhomogeneity length scales that have any value from one meter up to thousands of kilometers (in the case of coronal plumes). Nevertheless, the role of the drift wave in the problem of solar coronal heating has been overlooked so far in the literature, probably due to the fact that it simply does not exist in the widely used MHD model.

Regarding the wave heating analysis, in the case of the practically collision-less corona, an efficient mechanism for the transfer of energy from the wave to the plasma is needed. Within the framework of the drift wave kinetic theory, this process develops as follows: the interaction between the wave and the electrons is destabilizing and the mode grows due to a Cherenkov-type interaction. At the same time, however, its energy is absorbed by the ions due to Landau damping. This may be seen from ref. [13], where the drift wave properties within the limits, \( k_z v_{\text{T}e} \ll \omega \ll k_z v_{\text{T}i} \), \( \omega \ll \Omega_i \), and \( |k_y/k_z| (T_e/T_i)^{1/2} \rho_i/L_n \gg 1 \), are described by the frequency

\[
\omega_r = -\frac{\omega_i \Lambda_0(b_i)}{1 - \Lambda_0(b_i) + T_i/T_e + k_y^2 \lambda_{d1}^2} \tag{1}
\]

and the growth rate

\[
\omega_i \approx \left( \frac{\pi}{2} \right)^{1/2} \frac{\omega^2_{\parallel}}{\omega_{\parallel} \Lambda_0(b_i)} \left[ \frac{T_i}{T_e} \frac{\omega_i - \omega_{\parallel}}{|v_{\parallel}| v_{T_i}} \exp[-\omega^2_{\parallel} (k_y^2 v_{\parallel}^2/r_{\text{T}i})] + \frac{\omega_r - \omega_{\parallel}}{k_y v_{T_i}} \right] \tag{2}
\]

Here, \( \Lambda_0(b_i) = I_0(b_i) \exp(-b_i) \), \( b_i = k_y^2 \rho_i^2 \), \( \rho_i = v_{T} \Omega_i \), \( \lambda_{d1} = v_{T} / |v_{\parallel}| \), \( \omega_{\parallel} = -\omega_{\parallel} T_i / T_e \), \( \omega_{\parallel} = k_y v_{\parallel} \), \( \delta_{\parallel} = -(v_{T}^2 / \Omega_i) e_z \times \nabla \times n_0 / n_0 \), and \( I_0 \) is the modified Bessel function of the first kind and of the order 0.

Equation (1) reveals the energy source already in the real part of the frequency \( \omega_r \propto \nabla \times n_0 \), while details of its growth due to the same source are described by eq. (2). It is seen that, for a wave frequency below \( \omega_{\parallel} \), the growth rate can be written as \( \omega_i = |\gamma_{el} - |\gamma_{ion}| \). The energy exchange between the plasma particles and the wave is provided by two mechanisms simultaneously, viz. one in the parallel and one in the perpendicular direction. The term \( |\gamma_{ion}| \) is responsible for the Landau dissipation of the wave energy and, consequently, for the parallel heating of the plasma. So, as long as the density gradient is present, there is a continuous precipitation of energy from the wave to the plasma. On the other hand, the term \( |\gamma_{el}| \)
results in the growth of the wave, and this implies another (stochastic) heating mechanism that also involves single particle interactions with the wave. Note that this process is described and even experimentally verified [14,15] for a hot, fully ionized plasma, viz. in a tokamak. For drift wave perturbations of the form \( \phi(x) \cos(k_y y + k_z z - \omega t) \), with \( |k_y| \gg |k_z| \), one finds the ion particle trajectory in the wave field from the following set of equations:

\[
\frac{d\chi}{d\tau} = \Upsilon, \quad \chi \equiv k_y x, \quad \Upsilon \equiv k_y y, \quad \tau = \Omega_i t, \quad (3)
\]

\[
\frac{d^2\Upsilon}{d\tau^2} = -\Upsilon + \left[ m_i k_y^2 \phi / (e B_0^2) \right] \sin(\Upsilon - \omega/\Omega_i). \quad (4)
\]

It has been shown [14] that stochastic heating takes place for a large enough wave amplitude, more precisely for \( a = k_y^2 \rho_i^2 \phi / (\kappa T_i) \gg 1 \). The maximum achieved bulk ion velocity, proportional to the wave amplitude, is given by

\[
v_{\text{max}} \simeq \left[ k_y^2 \rho_i^2 e \phi / (\kappa T_i) + 1.9|\Omega_i/k_y. \right] \quad (5)
\]

Ideally, this all requires \( |\gamma_{el}| \gg |\gamma_{ion}| \) so that the wave amplitude may grow and at some moment both heating mechanisms may take place simultaneously. In the solar corona this condition can easily and quickly be satisfied because of the almost unlimited range of the parallel wave number \( k_z \), so that the ion Landau damping can be made small, i.e., \(|\omega/k_z| \gg v_{Ti} \), and because of the large growth rate as shown in the example in fig. 2.

In the stochastic heating process due to the drift wave, the ions move in the perpendicular direction to large distances and feel the time-varying field of the wave due to the polarization drift \( \vec{v}_p = (\partial \vec{E} / \partial \vec{B}) / (\Omega_i B_0) \), and as a result their motion becomes stochastic. The polarization drift is in the direction of the wave vector, which emphasizes the crucial electrostatic nature of the wave in the given heating process. Also, this stochastic heating is highly anisotropic, and it takes place mainly in the direction normal to the magnetic field \( B_0 \) (both the \( x \)- and \( y \)-direction velocities are stochastic). The perpendicular heating in the experiment [14,15] was larger by about a factor 3 compared to the parallel one, and it is exceptionally fast (see below). At the same time, in view of the mass difference and the physical picture given above, this heating scenario predominantly acts on the ions.

In application to coronal magnetic structures, the indication or “proof” that the heating really takes place would be: i) an ion temperature anisotropy \( T_{ix} \gg T_{iz} \), ii) a possibly higher ion temperature in comparison to electrons, and iii) a better heating of heavier ions. Observations show that i) may be taken rather as a rule than as an exception [2,3], i.e., the perpendicular stochastic heating is more dominant compared to the parallel heating. There are also numerous indications that confirm the features ii) and iii). As an example we refer to graphs from ref. [4], where \( T_e < T_H < T_{Hi} \) throughout the corona and the solar wind.

An easy way to demonstrate that the heating can take place at various inhomogeneity scale-lengths \( L_n \) (in other words in various magnetic field or density structures in the solar corona) is to keep the ratio \( \lambda_i/L_n \) fixed. For example, setting \( L_n = s \times 100 \text{m} \), and \( \lambda_i = s \times 40000 \text{m} \), where \( s \) takes values, e.g., between 1 and 1000, we calculate the frequency and the growth rate for the fixed value \( \lambda_i = 0.5 \text{m} \), and we find out that the ratio \( \omega_i/\omega_i \approx 1 \). Note that such a variation of \( s \) may also be used to describe the natural change of the radial density gradient with the increased altitude, in other words the heating occurs everywhere along a given flux tube.

Assuming an initial perturbation of the order \( \phi \tilde{\phi} / (\kappa T_i) \approx 0.01 \), (i.e., \( \phi = 0.86 \text{V} \)) for the parameters from fig. 2 (and for \( \omega_i \approx 2.5 \times 10^2 \text{Hz} \), \( \lambda_i \approx 0.5 \text{m} \)) the growth time \( \tau_g = \text{ln}100/\omega_i \) till it becomes of the order of unity (\( \phi = 86 \text{V} \)) is about 0.02 s. For the temperature increased by \( 10^6 \text{K} \) (see table 1) this implies a heating rate of ions of the order of \( 5 \times 10^7 \text{K/s} \), which is similar to the

| Table 1: Plasma heating for hydrogen ions for two perpendicular wavelengths and for two values of the wave amplitude \( \phi \). The values in brackets are for helium. The maximum stochastic velocity is given by eq. (5) and the effective temperature obtained by heating is \( T_{\text{eff}} \). |
|---|---|---|
| \( \phi = 60 \text{ (V)} \) | \( \phi = 80 \text{ (V)} \) |
| \( \lambda_i \) (m) | \( T_{e,\text{ff}} \) (K) | \( T_{i,\text{ff}} \) (K) |
| 0.5 | \( 1.56 \times 10^6 \) (1.91 \( \times 10^6 \) \( ) \) | \( 2.03 \times 10^6 \) (2.92 \( \times 10^6 \) \( ) \) |
| 1 | \( 2.72 \times 10^6 \) (1.56 \( \times 10^6 \) \( ) \) | \( 3.06 \times 10^6 \) (2.03 \( \times 10^6 \) \( ) \) |
heating rate obtained in the experiments [14,15]. Observe also that the magnitude of the electric field which we are dealing with is of the same order as in those experiments.

The stronger heating of heavier ions (see fig. 3) can be understood from eq. (5) and after expressing the effective temperature in terms of the ion mass \( T_{\text{eff}}(m_i) = m_i v_{\text{max}}^2 / (3 \kappa) \). From the derivative \( d T_{\text{eff}}(m_i) / dm_i > 0 \), it follows that the heating increases with the ion mass if \( k_i^a \rho_i^a (e \phi / \kappa T_i)^2 > 1.9 \). For \( \phi = 60 \text{ V} \) we have the normalized temperature \( T_i(\lambda_y, \mu) = 0.881 + 0.057 \mu / \lambda_y^2 + 1.78 \lambda_y^4 / \mu \).

For the same parameters as above, the maximum energy released per unit volume is \( \Sigma_{\text{max}} = n_0 m_i v_{\text{max}}^2 / 2 = 0.04 \text{ J} / \text{m}^3 \). The energy release rate \( \Gamma_{\text{max}} = \Sigma_{\text{max}} / \tau_g \approx 1.1 \times 10^{-3} \text{ J} / (\text{m}^3 \text{s}) \) amounts to four orders of magnitude above the necessary value. However, for \( L_n = 100 \text{ km} \) (i.e., setting \( s = 1000 \)) we obtain \( \omega_i = 0.13 \text{ Hz}, \omega_r = 0.254 \text{ Hz}, \tau_g = 32.6 \text{ s} \), and consequently \( \Gamma_{\text{max}} = 1.2 \times 10^{-3} \text{ J} / (\text{m}^3 \text{s}) \), that is about one order of magnitude above the heating rate accepted as necessary. Similar estimates may be done for still larger \( L_n \), yet the conditions under which the previous expressions are derived become violated and a numerical approach is required in this case. In reality the collisions and nonlinearity lead to the flattening of the density profile in the region occupied by the wave [16], which should result in the saturation of the growth. Also, energy diffusion in the perpendicular direction should act in the same way. In addition, the drift wave is as a rule coupled to the Alfvén wave [17], with the coupling proportional to \( k_y \rho_i \). All these effects will more effectively act on short scales, and the actual values for \( \Gamma \) are expected to be below \( \Gamma_{\text{max}} \). Therefore, the apparently too big release of energy at short scales, as formally obtained above, may in reality be considerably reduced. Clearly, more accurate estimates and more detailed description may be obtained only numerically.

To summarize, the proposed mechanism is based on a novel paradigm that allows a self-consistent solution model. The heating mechanism implies instabilities on time and spatial scales that are currently not directly observable by space probes. However, all the effects presented here are directly experimentally verified under laboratory conditions. Their indirect confirmation in the context of the solar corona seems to be also beyond doubt. This is because the consequences of the heating process, as enlisted earlier in the text (temperature anisotropy, better heating of heavier ions, and hotter ions than electrons), are indeed verified by satellite observations.

***

These results are obtained in the framework of the projects G.0304.07 (FWO-Vlaanderen), C 90205 (Prodex 9), and GOA/2009-009 (K.U. Leuven). TRACE (fig. 1) is a mission of the Stanford-Lockheed Institute for Space Research and part of the NASA Small Explorer program.

REFERENCES

[1] ASCHWANDEL M. J., Astrophys. J., 560 (2001) 1035.
[2] CUSERI I., MULLAN D. and POLETTO G., Space Sci. Rev., 87 (1999) 153.
[3] LI X., HABBAL S. R., KOHL J. L. and NOCI G., Astrophys. J., 501 (1998) L133.
[4] HANSTEEN V. N., LEER E. and HOLTZER T. E., Astrophys. J., 482 (1997) 498.
[5] KLIMCHUK J., Sol. Phys., 234 (2006) 41.
[6] SCHRIEVER C. J. et al., Sol. Phys., 187 (1999) 261.
[7] NOVEMBER L. J. and KOUTCHMY S., Astrophys. J., 466 (1997) 512.
[8] KAROVSKA M. and HABBAL S. R., Astrophys. J., 371 (1991) 371.
[9] WOO R., Nature, 379 (1996) 321.
[10] DE PONTIEU et al., Science, 318 (2007) 1574.
[11] ASCHWANDEL M. J., WÜLSER J. P., NITTA N. V. and LEMEN R. R., Astrophys. J., 679 (2008) 827.
[12] VRANJES J. and POEDTS S., Astron. Astrophys., 482 (2008) 653.
[13] ICHIMARU S., Basic Principles of Plasma Physics (The Benjamin/Cummings Publishing Company, Reading, Mass.) 1989, p. 133.
[14] SANDERS S. J., BELLAN P. M. and STERN R. A., Phys. Plasmas, 5 (1998) 716.
[15] MCDONALD M. J., BELLAN P. M. and STERN R. A., Phys. Rev. Lett., 59 (1987) 1436.
[16] LEE W. W. and OKUDA H., Phys. Rev. Lett., 36 (1976) 870.
[17] VRANJES J. and POEDTS S., Astron. Astrophys., 458 (2006) 635.