Coronal Heating & Solar Wind Acceleration by Drift Waves

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Abstract. An alternative approach to the coronal heating problem, based on the theory of drift waves, has been proposed. The drift mode is the only mode that is able to survive the drastically different (collisional/collisionless) extremes in the different layers of the solar atmosphere. As a matter of fact, this mode is over stable, i.e. able to grow, in both of these extreme situations, and has been called the universally growing mode in the literature. In collisional plasma of the lower layers of the solar atmosphere, the drift mode grows due to the electron collisions and this can be described within the two-fluid model. In the collisionless coronal plasma, however, the mode grows due to a pure kinetic effect, viz. the electron resonance effect in the presence of a density gradient.

It has been shown, with qualitative and quantitative arguments, that the drift waves have the potential to satisfy all coronal heating requirements. The basic ingredient required for the heating is the presence of density gradients in the direction perpendicular to the magnetic flux surfaces. The drift wave theory is well-established and has been explicitly verified experimentally in laboratory (fusion) plasmas, similar (hot, low-beta, highly conductive) to those in the solar atmosphere. In these circumstances, two mechanisms of the energy exchange and heating take place simultaneously: Landau damping in the direction parallel to the magnetic field, and stochastic heating in the perpendicular direction. The latter, in fact, is more effective on ions than on electrons, acts predominantly in the perpendicular direction, and heats heavy ions more efficiently than lighter ions. Moreover, for plasmas at a temperature of 1 MK and beyond, the parallel wave field resulting from the drift waves exceeds the Dreicer field so that the bulk plasma species (primarily electrons) can be accelerated/decelerated by the wave in the parallel direction. In addition, this acceleration is more effective on the particles that are already more energetic, resulting in a distribution function considerably different from a Maxwellian, similar to the observed kappa-distribution in the outer solar atmosphere and in the solar wind.

1. Introduction
The solar coronal heating problem, a 75 year old puzzle regarding the observational fact that the temperature in the corona of the Sun is more than two orders of magnitude higher than the surface temperature of the Sun, is still one of the major challenges in solar and astrophysics. As a matter of fact, none of the many proposed models is able to explain all the observational characteristics. This is partly related to the fact almost all the models proposed so far, use magnetohydrodynamics (MHD) as mathematical model. MHD is the simplest model, one-fluid and macroscopic (continuum model) and thus not valid on the length scales at which the actual dissipation takes place. As a matter of fact, in order to be conclusive, the model should take into account the physical conditions in the lower layers of the solar atmosphere. The density
and the temperature are not the only physical quantities that change drastically in the solar atmosphere. In fact, the chromosphere plays an important role as an intermediate layer between the photosphere and the corona, where the plasma beta (the ratio of the thermal pressure to the magnetic pressure) changes from above one to much below one ($\leq 1\%$), i.e. where the plasma changes from being thermally dominated in the low atmospheric layers to highly magnetically dominated in the low corona (further out in the corona the plasma $\beta$ rises again to order one at 1 AU). Moreover, the collisionality of the plasma varies from “strongly collisional” in the photosphere and in the lower part of the chromosphere, over “mildly collisional” in the upper part of the chromosphere, to the “collision-less” plasma regime in most of the corona and beyond. As a result, the ionization degree too varies drastically in the solar atmosphere. It can be as low as $10^{-4}$ at the temperature minimum (at the top of the photosphere), and is increasing to $\pm 50\%$ at a height of about 2000 km and to nearly 100% in the hot corona. Moreover, Carlsson et al. [1, 2] showed that the chromosphere is in a tricky non-equilibrium ionization and non-LTE (non-local thermodynamic equilibrium) regime.

Hence, the relatively narrow chromosphere is the layer where the plasma jumps from collisional to collisionless, from high beta to very low beta, from neutral to almost fully ionized, from opaque to optically thin (in UV and mm-waves), and increases drastically (by almost 3 orders of magnitude) in temperature. Improved observations by current missions, like Hinode (JAXA/NASA) and IRIS (NASA/ESA), make these complexities more apparent and provide even more evidence for the fact that the chromosphere plays a crucial role in many observed phenomena in the solar atmosphere. This is because the chromosphere is the photosphere/corona interface (cf. above) and the fact that the magnetic field, dominant in the corona, is anchored in the chromosphere and transition region and electric current systems (stressing the coronal field) penetrate the chromosphere.

Vranjes et al. [3, 4] studied the important issue of the energy flux of Alfvén waves bearing in mind the facts that the solar photosphere is very weakly ionized and the dynamics of the plasma particles in this region is heavily influenced by the plasma-neutral collisions. The magnetization of the plasma constituents have been quantitatively examined, showing that the ions and electrons in the photosphere are both un-magnetized, their collision frequency with neutrals is much larger than the gyro-frequency. This implies that eventual Alfvén-type electromagnetic perturbations must involve the neutrals as well. This has important consequences, viz. i) in the presence of perturbations, the whole fluid (plasma + neutrals) moves; ii) the Alfvén velocity includes the total (plasma + neutrals) density and is thus considerably smaller compared to the collision-less case; iii) the perturbed velocity of a unit volume, which now includes both plasma and neutrals, becomes much smaller compared to the ideal (collision-less) case; and iv) the corresponding wave energy flux for the given parameters becomes much smaller compared to the ideal case.

Vranjes et al. [5] then investigated the properties of gas acoustic and ion acoustic modes in a collisional, weakly ionized plasma in the presence of un-magnetized ions and magnetized electrons. In such a plasma, an ion acoustic mode, driven by an electron flow along the magnetic field lines, can propagate almost at any angle with respect to the ambient field lines as long as the electrons are capable of participating in the perturbations by moving only along the field lines. Several effects, including the electron-ion collisions, the perturbations of the neutral gas, and the electromagnetic perturbations, are studied. The neutral sound mode couples to the current driven ion acoustic mode, and these two modes can interchange their identities in certain parameter regimes. The electromagnetic effects, which in the present model imply a bending of the magnetic field lines, result in a further destabilization of an already unstable ion acoustic wave.
2. Drift waves in the solar atmosphere?!

In 2009, Vranjes and Poedts [6] proposed a new paradigm for the coronal heating problem by applying a multi-component fluid and kinetic description of the drift wave instability to the solar atmospheric plasma. This novel drift wave heating theory appears to be able of satisfying numerous heating requirements of the solar atmosphere. Indeed, a self-consistent solar coronal heating model must not only provide an energy source for the extremely high temperature in the corona including a reliable and efficient mechanism for the transfer of energy from the source to the plasma particles, and this at a required heating rate (e.g. \( \approx 10^{-4} \) J/(m\(^3\)/s) in active regions) [7]. It should also work everywhere in the corona (with different heating requirements in different regions). Moreover, it must also explain the observed discrepancy between ion and electron temperatures (typically \( T_i > T_e \)) and the origin of the large temperature anisotropy (\( T_\perp > T_\parallel \)) with respect to the direction of the magnetic field, particularly for ions [8, 9]. In addition, it should explain the observed stronger heating of heavier ions [10]. Clearly, the latter three requirements can never be fulfilled by a single-fluid isotropic pressure model like MHD [11].

![Figure 1. Visualization of a drift wave in cylindric geometry, propagating in the poloidal and axial directions with poloidal wave number \( m = 2 \), in plasma with a \( r \)-dependent equilibrium density. The actual wave fronts are twisted around the axis and have a \( r \)-dependent amplitude which reaches a maximum in the area of the largest density gradient. From [12].](image)

The theoretical model proposed by Vranjes and Poedts [6, 12, 13] shows that a) the energy for driving the mentioned drift modes and instabilities, and for the heating of the solar corona is already present in the corona, viz. stored in the ubiquitous density gradients and, b) this energy is naturally transmitted to the different plasma species by well known effects that are, however, beyond the commonly used models and theories. Moreover, it is based on well established, basic plasma theory which has already been verified and confirmed by means of a series of laboratory plasma experiments. It was shown that the basic ingredient necessary for the heating is the presence of gradients of the density (eventually also of the temperature and magnetic field) in the direction perpendicular to the magnetic flux surfaces. Such density gradients are a source of free energy for the excitation of drift waves. The dispersion relation in two-component kinetic theory, under the conditions \( k_z v_{Ti} \ll \omega \ll k_z v_{Te}, \omega \ll \Omega_i \), and

\[
\frac{|k_y| \rho_i}{k_z} \left( \frac{T_e}{T_i} \right)^{1/2} \gg 1,
\]

(with \( L_n \) the inhomogeneity scale length (\( L_n = (n'_0/n_0)^{-1} \)) reads [12]:

\[
\frac{|k_y| \rho_i}{k_z} \left( \frac{T_e}{T_i} \right)^{1/2} \gg 1,
\]
\[ \omega_r = -\frac{\omega_{sl} \Lambda_0(b_i)}{1 - \Lambda_0(b_i) + T_i/T_e + k^2 \lambda^2_{di}}, \quad \omega_{si} = k_y \frac{v_{Ti}^2 n_i'}{\Omega_i n_{i0}}, \quad \nabla \times n_{i0} = -\vec{e}_x n_{i0}', \] (1)

for the real part of the frequency and

\[ \gamma \simeq -\left( \frac{\pi}{2} \right)^{1/2} \frac{\omega_r^2}{|\omega_{sl}| \Lambda_0(b_i)} \left[ \frac{T_i}{T_e} \frac{\omega_r - \omega_{se}}{|k_z| v_{Te}} \exp\left[ -\omega_r^2 / (k_z^2 v_{Te}^2) \right] + \frac{\omega_r - \omega_{si}}{|k_z| v_{Ti}} \exp\left[ -\omega_r^2 / (k_z^2 v_{Ti}^2) \right] \right], \] (2)

for the imaginary part. Here

\[ \Lambda_0(b_i) = I_0(b_i) \exp(-b_i), \quad b_i = k^2_{g} p_i^2, \quad \lambda_{di} = v_{Ti}/\omega_{pi}, \quad \omega_{se} = -k_y \frac{v_{Te}^2 n_i'}{\Omega_e n_{e0}}. \] (3)

Note that the presence of the energy source is seen already in the real part of the frequency \( \omega_r \propto \nabla \times n_0 \) (compare with Alfvén waves: \( \omega = k c_a \)). The details of the growth of the wave energy, which is due to the same source, is described by the imaginary part \( \gamma \).

Two mechanisms of the energy exchange and heating have been 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. The stochastic heating i) is due to the electrostatic nature of the waves, ii) is more effective on ions than on electrons, iii) acts predominantly in the perpendicular direction, iv) heats heavy ions more efficiently than lighter ions, and v) may easily provide a drift wave heating rate that is several (up to 4) orders of magnitude above the value that is presently believed to be sufficient for the coronal heating [12], i.e., \( \sim 6 \cdot 10^{-5} \text{ J/(m}^3\text{s)} \) for active regions and \( \sim 8 \cdot 10^{-6} \text{ J/(m}^3\text{s)} \) for coronal holes. This heating acts naturally through well known effects that are, however, beyond the current standard models and theories.

3. Simple numerical model

As mentioned before, the drift wave theory (both linear and nonlinear) has been tested and confirmed thoroughly via laboratory experiments in the context of controlled thermonuclear fusion research. We here report on the first results of a numerical quantitative study under parameter values that are more relevant for the solar atmosphere. We focussed on lower hybrid drift waves (LHDWs), a subclass of drift waves, that has also been studied and observed in the magnetosphere of the Earth. These waves have an angular frequency \( \omega \) that satisfies \( \omega_{ce} \leq \omega \ll \Omega_{ci} \) and \( k r_{L,e} \leq 1 \ll k r_{L,i} \), resp., with \( r_{L,s} \) the Larmor radius for species \( s \) [14]. In this parameter range, the electrons are strongly magnetized and can be described as a fluid, while the ions (protons and heavier ions) are treated as particles. The LHDWs can resonate with ions in the perpendicular direction (w.r.t. the magnetic field) and with electrons in the parallel direction. The latter electron-wave interaction destabilises the modes via the inverse Landau effect [6].

3.1. Set up

We considered a simple Harris current sheet configuration (see Fig. 2 and [15]) which assumes a local Maxwellian velocity distribution, shifted towards the currents direction with a velocity value equal to the diamagnetic drift velocity that arises from the pressure gradients in the system. Our goal is to check how efficient the free energy (in the density gradients) can be transferred to the plasma species, and ultimately, if the plasma species can absorb enough energy to sustain the energy requirements of solar corona. The profiles of the magnetic field and the density (illustrated in Fig. 2) are given by

\[ B_z(y) = B_0 \tanh \left( \frac{y}{L} \right) \vec{e}_z, \] (4)
Figure 2. Harris Current Sheet Configuration, illustrating the magnetic field $B_z(y)$ and particle density $n(y)$ profiles considered. The diamagnetic velocities of electrons and ions, generating the current, are denoted by $v_{de}$ and $v_{di}$, resp. [16].

and

$$n(y) = n_0 \sech^2 \left( \frac{y}{L} \right),$$

respectively, where $L$ denotes the thickness of the sheet, and $B_0$ and $n_0$ correspond to the values of the magnetic field and density at the centre of the sheet, respectively. The magnetic field profile yields a current density in the $x$-direction:

$$J_x(y) = \frac{1}{\mu_0} \nabla \times \vec{B} = \frac{B_0}{\mu_0 L} \sech^2 \left( \frac{y}{L} \right) \vec{e}_x. \tag{6}$$

We assume quasi-neutrality, i.e. a balance between magnetic and plasma pressure. Initially, $T_i = T_e$, and for two later cases with three species each, $T_i = T_e = T_{O+1}$ and $T_i = T_e = T_{He+2}$.

3.2. Numerics

The simulations are done with IPIC3D, the “implicit Particle-In-Cell 3D” code, which is designed for a multiprocessor architecture and implements implicit methods to solve the equations governing the physics of the particle-heating problem at hand, i.e. the Vlasov equations for the distribution functions of the considered particle species, in combination with the Maxwell equations for the electro-magnetic field. For details and stability conditions etc., we refer to Markidis et al. [17]. For the results discussed in the present paper, the iPIC3D code was applied to a Harris current sheet configuration in which the particle’s velocity distribution is Maxwellian. The diamagnetic velocity of each particle species satisfies the relation $v_{s,i}/v_{s,e} = -T_i/T_e$, so that electrons and ions drift to opposite directions and thus generate currents. At the centre of the
sheet, the magnetic field has minimal magnitude while the plasma density is maximal there, so that the motion at the central region is dominated by inertia effects. The density gradients are controlled by the parameter \( L \), the thickness of the plasma sheet. The actual value of the ion mass, specified by the parameter \( m_i/m_e \) in the code, is not the real one but this does not affect the dimensionless parameters \( \omega_{pe}/\Omega_{pe} \), \( L/r_{L,i} \) and \( v_{s,i}/v_{th,i} \), with \( \omega_{pe} \) the plasma frequency and \( \Omega_{pe} \) the electron gyro-frequency. We consider a simulation box of \( 10 \times 5 \times 1/40 \) in units of ion skin depth, \( d_i \). The high-resolution grid has \( 400 \times 200 \times 1 \) cells, i.e. \( \Delta x = \Delta y = \Delta z = 0.025 \ d_i \). Each cell has 9 computational particles and a typical simulation took 20000 cycles.

**Figure 3.** Top left: snapshot of the charge density plotted in the \( x - y \) domain. The central region of the current sheet is situated at \( y = 2.5 \ d_i \). Bottom left: the same quantity for ions that have an artificial mass ratio value \( m_i/m_e = 256 \). Right: distribution profiles of the charge density of electrons and ions plotted respecting the \( y \)-coordinate and evaluated at \( x = 5 \ d_i \). Distance is in units of ion skin depth \( d_i \) and time is normalized by the ion plasma period \( \omega_{pi}^{-1} \).
4. Simulation results

4.1. General case with two species only

As a first case, we choose $v_{th,e} = 0.05 c$, $T_i/T_e = 1$, and $m_i/m_e = 256$. In Fig. 3, snapshots of the electron and ion charge density are plotted in the plane perpendicular to the magnetic field, i.e. the $x−y$ plane, showing the perturbations that appear at the edges of the current sheet. In order to show how the density profiles change over time, we also plotted the same quantities over the $y$-direction along $x = 5 d_i$.

In Fig. 4, shows another snapshot of the same run at a later time, when the non-linear phase has developed. Notice the broader profile in nonlinear phase and the kinking of the current sheet during this phase. Later a kink mode develops due to the Lower Hybrid Drift Instability (not shown here).

![Figure 4](image-url)

**Figure 4.** The kinking of the current sheet during the non-linear phase, for the same parameter values and same plot options as Fig. 3.

Next, we calculated the total magnetic and thermal energy density for the whole system by averaging out these energy contributions over the total number of cells and plotting them with respect to time. The results are displayed in Fig. 6, where physical values have been used to calculate the energy density values in SI units. For this case, the initial thermal energy of the
plasma species matched the case of the Earth magnetosphere. Notice that the magnetic energy of the system drops. As time elapses, the electrons both give and receive energy ($J_e \cdot E > 0$ or $< 0$), while the protons only receive energy ($J_p \cdot E > 0$), as shown in Fig. 5.

![Figure 5. Temporal evolution of the electromagnetic power density equation’s terms, represented in code units. The top plot represents the energy exchange between electrons and electrostatic fields and the lower plot corresponds to the energy exchange between ions and electrostatic fields. The time is normalized by the ion plasma period $\omega_{pi}^{-1}$, while the power density of the energy exchange terms is $J_s \cdot E d_i/(n_0 c^2 m_i^2)$, with $s$ the species.](image)

In Fig. 7, two snapshots are plotted of the distribution of different temperature ratios, for different spatial directions and different species. This demonstrates the obtained temperature anisotropies. As a matter of fact, notice that $T_i > T_e$, i.e. ions are heated more than electrons. Also, $T_{e,\parallel} > T_{e,\perp}$, so for electrons the parallel direction (w.r.t. the magnetic field) is heated preferably. For ions, on the other hand, $T_{i,\parallel} \leq T_{i,\perp}$, i.e. the perpendicular direction (w.r.t. the magnetic field) is preferably heated, confirming theory and observations.
Figure 6. Energy density in SI units for the case of the Earth magnetosphere. The time is normalized by the ion plasma period.

Figure 7. Two snapshots of the temperature ratios, calculated for the different spatial directions and for different plasma species. Left: an initial time-step. Right: a phase during the well developed kink mode. Indices $x, y, z$ denote the spatial coordinates of the Cartesian system in the Harris sheet configuration. The distance is normalized by the ion skin depth $d_i$ and time by the ion plasma period.

4.2. Case $v_{th,e} = 0.05c$, $T_i/T_e = 5$, $n_0 = 0.1 \text{ cm}^{-3}$ (magnetosphere)

Next, we varied the initial electron thermal velocity and/or the initial temperature ratio between the ions and electrons in different steps in the parameter study.

As a first step, we increased the initial ion temperature by a factor of five, keeping the
electron’s thermal velocity the same, i.e. $v_{th,e} = 0.05c$. In Fig. 8, the conversion to physical units was obtained using a particle density $n_0 = 0.1 \text{ cm}^{-3}$, typical for the Earth’s magnetosphere. As a result, the total energy content is much larger than before. Notice also that the magnetic energy drops faster and more than in the previous case (with $T_i/T_e = 1$). Also, in this case the electrons gain energy during nonlinear phase, in contrast to the case with $T_i/T_e = 1$, and also the protons receive more energy. Notice the fast heating rate (on scale of milliseconds).

![Figure 8. Energy density (SI) for $v_{th,e} = 0.05c$, $T_i/T_e = 5$.](image)

4.3. Case $v_{th,e} = 0.035c$, $T_i/T_e = 1$, $n_0 = 10^9 \text{ cm}^{-3}$ (solar)

Let us now consider the case with the electron’s velocity $v_{th,e} = 0.035c$. Reducing the only the electron temperature and using $n_0 = 10^9 \text{ cm}^{-3}$ for the conversion to SI, we now have a case more suitable for the solar corona. The resulting energy density evolution is displayed in Fig. 9. Notice the much larger energy content in this case, due to the higher density considered. In this case, the magnetic energy drops faster/more, the electrons gain more energy during the nonlinear phase and the electron energy increases more/faster than in the previous case. Notice that also the protons receive more energy and faster than in the previous cases. The heating occurs extremely fast, in a matter of nanoseconds!

4.4. Cases with three particles

Next, we investigated what happens when there are three different species. First, we consider oxygen atoms as the third species. In the simulation, we specify that the charge density of ions consists of an equal share of protons and Oxygen atoms, while their sum is equal to that of electrons. In other words, the considered oxygen is $O^{+1}$. Clearly, $m_{O^{+1}}/m_i = 16$ and we specify that, initially, $T_i = T_e = T_{O^{+1}}$. The other parameter values are set as before. This choice of parameters corresponds to the physical situation in the Earths magnetosphere again. For the conversion to SI units, we assumed a particle density $n_0 = 0.1 \text{ cm}^{-3}$, like before in this case.

Figure 10 displays the energy evolution for this case. It is clear that the oxygen particles are heated more than the protons (cf. observations). Also, the protons heated more than the
electrons, but less than in the first case we considered (with only two species), as part of the energy now goes to the heavier ion species. The heating rate is extremely high again. Heating occurs on time scales of milliseconds.

Figure 9. $v_{th,e} = 0.035 \, c$, $T_i/T_e = 1$, energy density (SI)

Figure 10. The energy density in SI units for the case of the Earth magnetosphere but now with three species (electrons, protons and oxygen atoms (singly ionised, $O^{+1}$)), with parameter choices $v_{th,e} = 0.05 \, c$, $T_i = T_e = T_{O^{+1}}$, $m_{O^{+1}}/m_i = 16$. Time is normalized to the ion plasma period.
The last case we consider in the present paper, is also a case with three species, but this time we chose Helium ions as the third species. In this case, $v_{\text{th},e} = 0.05c$, and we choose to set $T_i = T_e = T_{He^{+2}}$, initially. Of course, $m_{He^{+2}}/m_i = 4$. For the conversion to SI units, we again assumed a particle density $n_0 = 0.1 \text{ cm}^{-3}$, like in the previous case, for comparison.

Figure 11 displays the energy evolution for this case. It is remarkable that in this case, the (heavier) helium particle population is heated less than the protons, unlike any other heavier ion, just like theory predicts [12]. Note also that the time scales for the heating are extremely short gain, order of milliseconds.

Figure 11. The energy density in SI units for the case of the Earth magnetosphere but now with three species (electrons, protons and helium atoms (doubly ionised, $He^{+2}$), with parameter choices $v_{\text{th},e} = 0.05c$, $T_i = T_e = T_{He^{+2}}$, $m_{He^{+2}}/m_i = 4$.

5. Conclusions
Vranjes and Poedts [6, 12] proposed a new explanation for the heating of the solar corona based on drift waves. As well-known from basic plasma theory, drift waves are always present in (inhomogeneous) plasmas, because these waves are intrinsically unstable (in the literature, therefore, this is called the “universal instability”). Any gradient in density (and/or temperature or magnetic field strength) across the magnetic flux surfaces, causes this ‘universal’ instability that drives these waves to occur. In other words, the drifts waves are ‘over stable’, i.e. with a frequency with both a real, oscillatory part, and an imaginary (in complex notation at least) growth rate. This distinguishes them clearly from the MHD waves that are often considered in the context of coronal heating, and that need to be driven at the right frequency in order to occur. Given the highly inhomogeneous characteristics of the plasma in the solar atmosphere, drift waves must be always self-generated there, and occur ubiquitously, even though they have not yet been observed (in the literature, any wave mode observed in the solar atmosphere is labeled as one or a combination of the MHD modes).

The alternative heating mechanism proposed by Vranjes and Poedts [6, 12], claiming to have the potential to satisfy all the observational requirements, was tested in a first attempt to quantify its efficiency. As a matter of fact, Vranjes and Poedts only applied linear theory and derived their preliminary conclusions from the large instability windows with high growth rates.
they obtained for coronal plasma parameter values. Here, a first attempt was made to take into account the non-linear saturation that inevitably occurs after the first (short) linear growth of the mode, to which the previous studies were limited. This was achieved by a relatively simple first numerical set-up and using an advanced particle-in-cell code (iPIC3D). The (first, simple) numerical simulations confirmed the linear theory results on the resulting temperature anisotropy, the fact that ions are heated more than electrons in this mechanism, the fact that heavier ions are heated stronger than lighter ions, except for helium ions, and the extremely fast heating time scales. These first results are, therefore, encouraging and need to be extended in order to be able to draw final conclusions.

As a matter of fact, a full quantitative confirmation of this heating mechanism needs much more advanced numerical simulations and parameter studies to include, e.g., the effect inhomogeneity along field lines, possible geometry effects (on the nonlinear saturation), etc. Clearly, also observational evidence would be welcome for drift waves in the solar atmosphere.

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