MAGNETIC SEISMOLOGY OF THE LOWER SOLAR ATMOSPHERE

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ABSTRACT

The dynamic interaction between solar photospheric motions and the magnetic structures in the lower solar atmosphere is a rather challenging problem. Coherent motions, e.g. acoustic or $p/f$-modes have important implications for the local dynamics of the solar atmosphere ranging from the low chromosphere deeply even into the corona. Coherent atmospheric fields, however, may interact with global oscillations altering their properties measurable at photospheric levels. Combined ground-based and satellite high spatial and time resolution observations (SST in La Palma and SOHO, TRACE) supplemented with MHD modelling demonstrated $p$-mode leakage into the chromosphere, transition region and low corona causing, among others, spicule formation, moss and loop oscillations.

In my paper I review the latest results on how photospheric motions interact with the magnetic structures and waveguides of the transitional layer between the photosphere and the corona. A comprehensive study of this transitional layer, also called solar atmospheric boundary layer, allows us to perform lower atmospheric magneto-seismology. Theoretical and observational efforts on the coupling mechanism(s) of both coherent and random motions and magnetic fields to the low atmosphere are discussed. Key issues will be addressed, including what dynamic impact the photosphere has on the overlaying magnetic atmosphere; what is the role of magnetic wave guides in the photosphere - chromosphere - transition region dynamics; what are the possible scenarios and physical details of the boundary layer coupling mechanism(s); how $p/f$-modes resonantly interact to lower atmospheric MHD slow and Alfvén waves; are the global oscillations influenced by the solar magnetic carpet; how the coupling could be used for atmospheric diagnostics, lower atmospheric magnetic seismology and connectivity studies.

Key words: solar physics; helioseismology; magnetohydrodynamics; fundamental mode; global acoustic modes.

1. INTRODUCTION

There are now abundant evidences that the solar atmosphere is magnetically coupled from its lower boundary, the photosphere, through the chromosphere, transition region (TR) up to its open-ended upper region, the corona. Concurrently taken images in visible, UV, EUV and X-ray light representing an atmospheric scanning (i.e. mapping temperature) along height clearly indicate that magnetic field concentrations, e.g. an active region, will show up as strong brightenings. One can also observe, at the same time, how dynamic and inhomogeneous this coupling is on short and broad time and spatial scales.

Traditionally helioseismology is concerned about solar global acoustic oscillations that are trapped mainly in the interior of the solar resonator: from the deep interior to the upper convection zone, or photosphere. For more details on recent progress in helioseismology, see the review paper by Thompson (2006).

Recent high-resolution satellite imaging and spectroscopy provides us with unprecedented spatial and time resolution of a magnetically strongly structured solar atmosphere from dark filaments to X-ray bright points that show, among many other feature, periodic motions, i.e. waves and oscillations. The main differences between these periodic motions and global oscillations are, at current resolution limits, at least two-fold: (i) the observed oscillations are local and within the magnetic structure, i.e. locked to it, in contrast to solar acoustic oscillations that have a more global character; (ii) one usually sees one or in the best case two eigenmodes within a local waveguide, as opposed to the millions of modes observed by GONG, MDI of one waveguide. It is natural to raise the question, whether from such observations would one be able to derive diagnostic information, e.g. density along this magnetic structure, what is its structure, curvature, etc. The idea of coronal seismology was put forward by Roberts et al. (1984) in a seminal paper. By extending the concept to lower parts of the solar atmosphere, where plasma-$\beta$ is of the order of unity, complex interesting and very promising application can be made.

In a first instance, a typical lower atmospheric waveguides, could be modelled as an isolated magnetic flux tubes with practically no magnetic field in their environment. Another feature that distinguishes these waveguides from their coronal counterparts is that the density and pressure scale heights inside the waveguide are comparable to their radial dimensions, leading even in the simples approximation to a Klein-Gordon-type linear wave equation governing the perturbations (Roberts 2004).
In order to proceed with the seismologic approach for the lower solar atmosphere, one first has to be clear about the role of the presence of the boundary layer between the solar interior and the corona. The transition between the solar interior and the corona occurs in a rather narrow layer when compared to the approximately 700 Mm solar interior. This boundary layer, that includes the photosphere, chromosphere and TR is around 2-3 Mm thick and contains coherent and random magnetic and velocity fields, giving a very difficult task to describe even in simplest approximate terms wave perturbations. We identify four coupling field elements. Random flows (e.g. turbulent granular motion), coherent flows (meridional flows or the near-surface component of the differential rotation), random magnetic fields (e.g. the continuously emerging tiny magnetic fluxes or magnetic carpet) and coherent fields (large loops and their magnetic canopy region). Each of these field elements may have their own effect on wave and oscillation perturbations. Of course, some may be more determinant with more important contribution than others. The magnitude of these corrections to e.g. eigenfrequencies or line widths of an ideal and ‘clean’, purely acoustic solar global resonator has to be estimated one by one and. It turns out, unfortunately, these coupling field elements all contribute to line widths and frequency shifts of the acoustic p/f-mode oscillations on a rather equal basis. In helioseismologic terms corrections from this boundary layer are labelled as surface term contributions (Basu 2002) and in many helioseismic modellings the surface term is taken in some ad-hoc functional form.

One of the main aims of the present review is to summarise our current knowledge about progress made in identifying the importance of coupling field elements of the ‘surface term’ as far as wave and oscillation perturbations are concerned. In particular, we focus on two aspects: (i) How global solar oscillations may couple to the inhomogenenous boundary layer and how the presence of this lower atmospheric boundary layer shifts the frequencies of measurable global acoustic oscillations and contributes to their line widths. Through simple examples we demonstrate how random and coherent fields, both magnetic and velocity, can influence the wave coupling of internal oscillations with the solar atmosphere. (ii) What are the dynamic dynamical consequences of the coupling in the lower atmosphere (and upper) solar atmosphere, e.g. is leakage of the global oscillations accountable for spikes of running coronal waves. With this type of approach and methodology we introduce the atmospheric magneto-seismology which extends the term of coronal seismology, by applying similar basic principles to the lower atmospheric boundary layer.

Finally, let us make one more introductory remark: In spite of the discovery of the correlation between changes of global acoustic frequencies and the solar cycle (e.g. Woodard and Noyes 1985; Woodard and Libbrecht 1991), the focus of attention has generally not turned towards the importance of the atmospheric or boundary layer (magnetic) fields. By now there is an overwhelming evidence for the above mentioned correlation, and let us recall just a few: GONG (Global Oscillations Network Group) observations for a single multiplet (e.g. l=50, m=9) indicate a strong correlation between the line width, magnetic flux and sunspot numbers (Komm et al. 2000); BiSON (Birmingham Solar-Oscillations Network) observations of line widths variations of low-angular degree p-modes during the fall of Solar Cycle 22, averaged over 2.6 to 3.6 mH, showed 24±3% mean increase in the modal line width from activity minimum to maximum as a function of the 10.7 cm radio flux, which is an excellent proxy for low-atmospheric magnetic field (Chaplin et al. 2000); Last, but not least there is a recent example by Dziembowski and Goode (2005) who interpreted SOHO (Solar and Heliospheric Observatory) MDI (Michelson Doppler Imager) data on oscillation frequency changes between 1996 and 2004, focusing on differences between the activity minimum and maximum of Solar Cycle 23. They found that both the f-mode and p-mode frequencies are correlated with general measures of the Sun’s magnetic activity.

Here I will focus on frequency and line width changes that are attributed to random and coherent fields in the boundary layer. I will demonstrate the importance of the results achieved on the coupling of global perturbations to the lower solar atmosphere. We will conclude, how these global oscillations may penetrate into the solar atmosphere and me be accounted for a range of dynamical phenomena.

2. ACOUSTIC WAVE LEAKAGE

Solar global oscillations driven by pressure forces (p-modes) are acoustic waves trapped in the solar interior within a resonant cavity. The cavity is a region where a wave can propagate but is bounded from above and below by regions where it cannot propagate any longer. An outward propagating wave is reflected inward from the solar upper surface, or boundary layer, because of a sudden decrease in the plasma density, while the lower boundary of the cavity is formed by the increasing sound speed due to temperature rises. If the frequency of a global oscillation mode is above a certain value (i.e. the acoustic cut-off frequency), the wave leaks out from the cavity into the atmosphere above. Such leakage can influence the dynamics of the lower (and upper) atmosphere, as was found by De Pontieu et al. (2004). They showed that oscillatory spikes are driven by the leakage of global oscillations along inclined thin flux tubes. An interesting consequence of this leakage is when these photospheric motions propagate even further into the lower corona, causing loop oscillations (De Pontieu et al. 2005). The solar f-mode, a gravity wave driven by buoyancy forces at the solar surface, resides right below the solar photosphere and can be regarded as a trapped mode in a cavity of zero depth, so that the upper and lower turning points coincide. This mode is similar to the well-known surface modes propagating along the free surface of a deep ocean.
Early theoretical studies (Ulrich 1970; Leibacher and Stein 1971) have predicted that the eigenfrequencies of solar global oscillations follow separated parabolic ridges as function of the horizontal wave number. In the simplest possible model, in a plane parallel geometry a semi-infinite plasma with a free surface may represent the solar interior. Assuming a compressible plasma with polytropic temperature increase as a function of depth from its surface, one can solve the eigenvalue problem of this configuration. The solutions are parabolic ridges for the \( p \) and \( f \) modes. Soon after the parabolic ridges of frequencies were predicted, Deubner (1975) observed these separated ridges of power. High-resolution measurements show deviations from the ridges. The above semi-infinite model can be improved by assuming, say, an isothermal atmosphere above the solar interior. The eigenfrequencies of such two-layer solar model are shown in Figure 1a and the eigensolutions of a few eigenmodes can be seen in Figure 1b (courtesy B. Pintér). In order to improve our modelling approach the next step is to consider a three-layer model where there is a narrow (approximately \( L=2 \) Mm thick) boundary layer between the solar interior and the hot isothermal solar atmosphere. Again, solving for the frequencies and eigenfunctions of this geometry, one recovers the internal \( p/f \)-modes; however, they have a more pronounced tail in the boundary layer and in the atmosphere (figure 2). In this three-layer model the boundary layer \( g \)-modes also appear.

3. ATOMIC ROTATIONAL FIELD

So far no magnetic fields have been considered. We recall the observations, that the global acoustic frequencies show small but significant and systematic correlations with the solar cycle. Although the acoustic modes are strongly evanescent in the atmosphere, changes of magnetic fields and mean temperature at the boundary layer (i.e. lower atmosphere) could play a rather important role in the determination of their frequencies. Of course the mass inertia of the solar atmosphere is negligible compared to its interior counterpart. However, in general, when one solves an eigenvalue problem, the eigenfrequencies and eigensolutions can be rather sensitive to the boundary conditions, not just to the conditions in the domain itself. Since magnetic or flow fields in a boundary layer (i.e. photosphere - chromosphere -TR) can change the mean elasticity of the boundary itself or alter the upper turning points these effects could contribute small corrections to the eigenfrequencies.

The pioneering papers by Roberts & Campbell (1986) and Campbell & Roberts (1989; CR89 hereafter), opened a new series of studies of the coupling of solar global oscillations to the lower solar atmosphere. The influence of a chromospheric (i.e. boundary layer) magnetic field on \( p \) and \( f \)-mode frequencies was evaluated theoretically for a simple and elegant model of the solar plasma, consisting of a polytrope in the solar interior, above which is an isothermal atmosphere. The atmosphere is permeated by a horizontal magnetic field. The eigenfrequencies of this model are shown on figure 3. Frequency changes and shifts in phase factors due to the presence of a magnetic atmosphere were calculated analytically in the long-wavelength limit, and numerically for arbitrary
A promising model was put forward by Evans & Roberts (1991, 1992). The effect of the magnetic canopy on the solar acoustic modes was extended to allow for variations in height of the magnetic canopy. Analytical solutions in the limit of long horizontal wavelength were obtained; the solutions exhibit explicitly the dependence of frequency shifts on magnetic field strength, wavenumber, and canopy height. Frequency shifts are principally due to changes in canopy height. Full numerical solutions were also presented. It was found that changes in chromospheric magnetism can be manifested in $p$- and $f$-mode data sets gathered at different phases of the solar cycle. These predictions of solar-cycle variability in high-degree $p$-mode frequencies from a simple model of the magnetic canopy which permeates the solar atmosphere were compared with the observations of Libbrecht and Woodard (1990). Good agreement was found with

the observed frequency shifts for modes of frequency less than 4 mHz, through a mechanism in which an increasing magnetic field induces “stiffening” of the Sun’s chromosphere (see also Wright and Thompson 1992).

Before we embark on studies of flow fields on global oscillations there is one more case of a magnetic boundary layer in a static plasma to be recalled. If the characteristic lengths of perturbations are small compared to the gravitational scale-heights one can approximate the boundary layer by a simple surface (e.g. there is a jump in density or temperature). In this case propagations could be both parallel to the atmospheric magnetic field lines (see e.g. Miles & Roberts 1992, Miles et al. 1992) or the horizontal wave number of the perturbations may have a finite angle to the field lines (e.g. Uberoi & Narayanan 1986, Gonzales & Gratton 1991, Jain & Roberts 1994a). These studies are relevant to $f$-modes.

4. RANDOM MAGNETIC CARPET

Similarly to granulation, the magnetic field in the solar atmospheric boundary layer is also very dynamic. High resolution magnetograms reveal that outside active regions the solar surface is covered with a mixed polarity network, which is termed magnetic carpet (Title & Schrijver 1998). The structure of this small-scale field changes rapidly on very short spatial and time scales and flux continuously emerges and disappears nearly homogeneously over the surface. Interestingly, the smallest magnetic structures show apparently no correlation with the solar cycle (Hagenaar et al. 2003), and therefore they are believed to originate in a separate, small-scale, dynamo process possibly close to the surface. Erdélyi et al. (2005) investigated the influence of this disorganised, small-scale atmospheric field on the $f$-mode frequencies. In a first approximation the magnetic carpet was modelled as a time-independent, stochastic field. Since, de-
pending on their spherical degree, some $f$-modes may have a life-time comparable to the characteristic replacement time (of the order of tens of hours) of the magnetic carpet, this limits the validity of their study. The magnetic field was taken to be independent of time because firstly they wanted to assess the effect of random magnetic field alone and in the case of a time-dependent field one would have to deal with generated flows in the initial state. They found that a time-independent random magnetic field can significantly increase the $f$-mode frequencies, in contrast with random steady flows which tend to have an opposite effect (Murawski & Roberts 1993a,b). Observations of $p$-modes show a more complex picture, and it is found that for some spherical degrees there is actually a frequency decrease as a function of the sunspot number - a good proxy for magnetic flux. In reality both magnetic and velocity fields are very dynamic at the solar surface, and therefore studying their interaction could be of crucial importance in interpreting the observed frequency shifts.

Bi et al. (2003) studied the influence of magnetic perturbations inside the Sun on the low-$l$ solar $p$-mode oscillations. They described the various possibilities of frequency shifts for a time-dependent source of MHD turbulence. For the magnetic perturbation contribution, they obtained the frequency shifts of modes with different degree as a function of the spectrum of fluctuating magnetic field. The frequency shift was found to increase with the strength of magnetic fields in the solar interior, and its temporal behaviour closely follows the phase of the synthetic solar activity cycle. This analysis indicates that the magnetic activities cause shifts of up to 0.3 $\mu$Hz. It is shown that the mode frequency, which is sensitive to the effect of magnetic fields, can also be used as a diagnostic tool for the presence of turbulent magnetic fields in the convection zone.

5. FLOW FIELDS

Flows at the boundary layer in the lower atmosphere may be random or coherent. By inverting the observational data of solar global oscillations one could potentially reconstruct the global flow structures. Large-scale subsurface flows were found by this technique (e.g. Braun and Fan 1998). They measured the mean frequencies of acoustic waves propagating towards and away from the poles of the Sun from observations made with MDI on board SOHO and the ground-based GONG. Significant frequency shifts between poleward- and equatorward-travelling waves measured over solar latitudes 20 - 60 degrees, which is consistent with the Doppler effect of a poleward meridional flow on the order of 10 m s$^{-1}$. From the variation of the frequency shifts of $p$-modes (with degree $l$ between 72 and 882) as a function of the lower turning point depth, they inferred the speed of the meridional flow, averaged over these latitudes, over a range in depth extending over the top half of the solar convection zone. Interestingly, there was no evidence for a significant equatorward return flow within this depth range.

Howe et al. (2000) have completed an analysis of the first 35 GONG Months (1 GONG Month = 36 days), covering the last solar minimum and the rising phase of Solar Cycle 23. The mode parameters have been estimated from 33 time series, each of 3-GONG Month duration, with centres spaced by 1 GONG Month. They reported on the temporal evolution of the rotational splitting coefficients up to 15th order. The coefficients do not correlate well with any surface magnetic flux measure yet considered, but Howe et al. found small though significant trends in their temporal evolution. Inverting the coefficients for two-dimensional rotation information and looking at deviations from the mean produces a picture of a systematic zonal flow migrating towards lower latitudes during the rising phase of the cycle. This flow is probably associated with the torsional oscillation. Similar trends were seen in the 1986-1990 BBSO data. These large-scale flow are important for at least two reasons. If there are slowly varying large-scale flows in the boundary layer (or around) they change the physical properties of the coupling mechanism between the interior and the atmosphere, resulting in changes in the eigenfunctions and eigenfrequencies. However, if there is damping of global oscillations (i.e. if line widths could be measured with sufficient accuracy), due to some dissipative mechanism present in the magnetic plasma, one could measure changes in the damping (i.e. in line widths) as a function of time in a slowly varying steady state. Inverting such measurements would give a clue as to the sub-surface flow structure.

To the best of our knowledge, Erdélyi et al. (1999) were the first who studied the effect of a sub-surface motion on magnetoacoustic-gravity (MAG) surface waves, representing the $f$-mode in a model of the solar interior - solar atmosphere interface. The main characteristics of their isothermal atmosphere was a magnetic (though constant) plasma-$\beta$, while in the sub-surface interior region there was a uniform and homogeneous equilibrium flow. They found that the flow causes a shift of the forward and backward propagating MAG modes, which in certain cases bifurcate. Erdélyi & Taroyan (1999, 2001) and Taroyan (2003) generalised the model by allowing the temperature to increase linearly with depth in the sub-surface zone. They derived the dispersion relation and analytical formulae for the frequencies of $p$- and $f$-modes in the limit of small wave numbers. Numerical solutions were presented for other cases.

The influence of short-scale motion (i.e. granulation) modelled as a random flow on $f$-mode frequencies was first evaluated by Murawski & Roberts (1993a,b; see also Murawski & Goossens 1993, Ghosh et al. 1995, Gruzinov 1998 and Medrek et al. 1999). Erdélyi et al. (2004) has re-visited this problem. The $f$-mode is essentially a surface wave; hence the mode frequencies are less likely to be influenced by the solar stratification. Most probably the discrepancies are the result of near surface mechanisms, such as interactions with surface or sub-surface magnetic fields and flows. Erdélyi et al. (2004) followed the general approach of Murawski & Roberts, which is a valuable one, but corrected certain
errors which appeared in that paper. The simple model used by Murawski & Robertes and Erdélyi et al. gives a deviation of the f-modes from the theoretically predicted parabolic ridges which agrees qualitatively with observations. They found that turbulent background flows can reduce the eigenfrequencies of global solar f-modes by several per cent, as found in observations at high spherical degree. Extensive numerical simulations of the outer parts of the Sun carried out by, e.g., Rosenthal et al. (1999) demonstrated and quantified the influence of turbulent convection on solar oscillation frequencies.

In what follows we discuss in detail the modelling efforts when both magnetic and flow field are present in or around the boundary layer. In §2 the governing equations are discussed, followed by the solutions for a steady convective zone in §3. In §4 and 5 two particular magnetic profiles are studied, representing weak and strong magnetic activities in the boundary layer and in the solar atmosphere. The dispersion relation is derived in each case. Analytical solutions are obtained for the corrections to the eigenfrequencies in the limit of long wavelength approximation. §6 is devoted to a discussion of resonant coupling of solar global oscillations to the atmospheric boundary layer. We conclude in brief in §6.

6. GOVERNING EQUATIONS

In order to model mathematically the coupling of solar global oscillation to the boundary layer and the solar atmosphere we consider initially a two-layer model. To make analytical progress we introduce a constant horizontal flow in the lower (internal) region \( z > 0 \), while the upper atmospheric regions of the model are embedded in a uniform magnetic field. The models can be grouped into two categories: (i) an homogeneous and uniform atmospheric magnetic field that may represent the solar atmosphere at high solar activities; we call this the case of a strong magnetic field; or (ii) an homogeneous though non-uniform magnetic field, where the magnetic field strength decays to zero far away from the solar surface. This latter model may be applicable to the solar atmosphere at low magnetic activity (i.e. at solar minimum); we call it the weak field approximation. Such solar models, without a flow, have been considered by CR89 and ER90. We outline the dispersion relations and show the solutions analytically in the limit of small wavenumber, for both cases. The obtained dispersion relations are also evaluated numerically for arbitrary wavelengths. The results are in good agreement with the analysis based on helioseismic measurements of the sub-surface meridional flows (see, for example, Braun & Fan 1998).

The linearized governing equations in compressible MHD are in standard notation

\[
\frac{\partial u_1}{\partial t} + \nabla \cdot (\rho_0 u_1 + \rho_1 u_0) = 0, \tag{1}
\]

\[
\frac{\partial \rho_1}{\partial t} + \nabla \cdot (\rho_0 u_1 + \rho_1 u_0) = \rho_1 u_0 \cdot \nabla u_0 = -\nabla p_1 + \frac{1}{\mu} (\nabla \times B_1) \times B_0 + \frac{1}{\mu} (\nabla \times B_0) \times B_1 + \rho_1 \mathbf{g}, \tag{2}
\]

\[
\frac{\partial \mathbf{B}_1}{\partial t} = \nabla \times (u_0 \times \mathbf{B}_1) + \nabla \times (u_1 \times \mathbf{B}_0), \tag{3}
\]

\[
\frac{1}{\rho_0} [\rho_1 u_0 \cdot \nabla p_0 - \gamma \rho_0 u_0 \cdot \nabla (\rho_0 u_0)] + \frac{\partial p_1}{\partial t} + u_0 \cdot \nabla p_1 + u_1 \cdot \nabla p_0 - c_s^2 \left( \frac{\partial \rho_1}{\partial t} + \frac{\nabla (\rho_0 u_0)}{\rho_0} + \frac{\nabla \rho_0}{\rho_0} + \frac{\nabla \rho_1}{\rho_1} \right) = 0, \tag{4}
\]

where all equilibrium quantities, denoted by the index 0, depend only on depth \( z \) and \( c_s^2(z) = (\gamma \rho_0 / \rho_0)^{1/2} \) is the local sound speed. The governing equations are supplemented by the solenoidal condition, \( \nabla \cdot \mathbf{B} = 0 \). Further, we express all perturbed vector and scalar quantities in the form

\[
f_1 = (f_x, 0, f_z) \exp (i(\omega t - kx)), \quad q_1 = q \exp (i(\omega t - kx)),
\]

respectively, where \( f_x, f_z, q \) are functions of \( z \), and consider Eqs. (1)–(4) in the convection zone and magnetic boundary layer separately.

7. DISPERSION RELATION

7.1. The solar interior

In the convection zone \( z > 0 \) we suppose the magnetic field to be absent and to have a uniform homogeneous steady flow along the x-axis, i.e., \( u_0 = (V, 0, 0) \). Then from Eqs. (1)–(4) we obtain

\[
(\omega_D^2 - g^2 k^2) u_z = -\omega_D^2 c_s^2 \frac{4 \Delta}{dz} - g(\gamma \omega_D^2 - k^2 c_s^2) \Delta, \tag{5}
\]

\[
\frac{du_z}{dz} = -\frac{g k^2}{\omega_D^2} u_z - \left( \frac{k^2 c_s^2}{\omega_D^2} - 1 \right) \Delta. \tag{6}
\]

Here \( \omega_D = \omega - kV \) is the Doppler shifted frequency, \( \Delta = -ik u_x + du_z / dz \), and \( \gamma_p \) is the adiabatic index in the region \( z > 0 \). Elimination of \( u_z \) yields a second-order ordinary differential equation

\[
\frac{d^2 \Delta}{dz^2} + \left( \frac{c_s^2}{c_s^2} \right) \frac{d \Delta}{dz} + \left[ \omega_D^2 - k^2 c_s^2 \right]
\]

\[
-\frac{g k^2}{\omega_D^2} \left( \frac{c_s^2}{c_s^2} - (\gamma_p - 1)g \right) \Delta = 0, \tag{7}
\]

where the prime denotes the derivative with respect to \( z \).

In the special case of a linear temperature profile, i.e., when the sound speed is given by

\[
c_s^2(z) = c_0^2 + c_s^2 v(z), \tag{8}
\]

Eq. (7) has the solution

\[
\Delta(z) = \exp[-k(z+z_0)][C_1 M(-a, m+2, 2kz+2kz_0)+
\]

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where \( z_0 = c_s^2 / \gamma \), and \( M \) and \( U \) are the confluent hypergeometric functions (Abramowitz & Stegun 1965); \( C_1 \) and \( C_2 \) are arbitrary constants. The parameter \( a \) is given by
\[
a = \frac{m + 1}{\gamma p} \frac{\omega_D^2}{2 g k} + \left( m - \frac{m + 1}{\gamma p} \right) \frac{g k}{2 \omega_D^2} - \frac{m}{2} - 1,
\]
and
\[
m = \frac{\gamma g}{c_s^2} - 1 = \frac{\gamma g z_0}{c_s^2} - 1
\]
is the polytropic index.

We require the kinetic energy density \( (\rho_0 u^2) / 2 \) to be finite as \( z \to \infty \). This implies that \( C_1 = 0 \) in Eq. (9), resulting in
\[
\Delta(z) = C_2 \exp[-k(z+z_0)] U(-a, m+2, 2kz+2kz_0).
\]

There are now two possibilities to consider, as mentioned before: (i) an atmosphere with strong or weak magnetic field. In both cases one will arrive to a dispersion relation that has solutions representing the global acoustic modes modified by the presence of the magnetic field.

### 7.2. Strong atmospheric magnetic field

In the transitional layer \( z < 0 \) the steady flow is absent and a uniform magnetic field
\[
B_0 = (B_0, 0, 0)
\]
is present. The magnetohydrostatic equilibrium is governed by
\[
\frac{d}{dz} \left( \rho_0(z) + \frac{B_0^2(z)}{2 \mu} \right) = \rho_0(z) g.
\]

From Eqs. (1)–(4) we obtain the equations
\[
\frac{du_z}{dz} = -\frac{g k^2}{\omega^2} u_z - \left( \frac{k c_s^2}{\omega^2} - 1 \right) \Delta,
\]
and
\[
\frac{d}{dz} \left[ \frac{\rho_0(c_s^2 + v_a^2)(\omega^2 - k^2 c_s^2)}{\omega^2 - k^2 c_s^2} \right] u_z = \frac{\rho_0 g^2 k^2}{\omega^2 - k^2 c_s^2} - \rho_0 g \left( k \frac{c_s^2}{\omega^2} - \frac{k^2 c_s^2}{\omega^2 - k^2 c_s^2} \right) \frac{d}{dz} \left( \frac{\rho_0 c_s^2}{\omega^2 - k^2 c_s^2} \right) u_z,
\]
where \( v_A(z) = B_0 / (\mu_0 \gamma) \) is the local Alfvén speed, and \( c_T = c_s v_A / (c_s^2 + v_a^2)^{1/2} \) is the MHD cusp speed.

We assume the boundary layer to be isothermal. In order to obtain the dispersion relation observe that the vertical component of the Lagrangian displacement and the Lagrangian perturbation of total pressure are continuous across the interfaces \( z = 0 \). This can be expressed by the following equations:
\[
[x_z]_{z=0} = 0, \quad [P + \rho_0 g x_z]_{z=0} = 0,
\]
where \( P = p + B_0 \cdot B / \mu \) and the square brackets denote the jump of the enclosed quantity across the interface. After some algebra (see ER90) we obtain the condition for the existence of a nontrivial solution yielding the dispersion relation
\[
2 a k \omega_D^2 U(1-a, m+3, 2kz) \frac{g k^2}{2 \omega_D^2} - 2k \omega_D - g k^2 = \frac{\rho_0 p}{\rho_0 c g k^2 c_s^2 + (v_a^2 + \gamma A_c)(\omega^2 - k^2 c_s^2)} f \varphi
\]
where
\[
\varphi = 1 - \frac{p q A_1 A_3}{k r A_2} \frac{F(p + 1, q + 1, r + 1, -\frac{A_1}{A_2})}{F(p, q, r, -\frac{A_1}{A_2})}
\]
and \( \rho_0 p \) and \( \rho_0 c \) are the limits of the density \( \rho_0 \) when \( z \to +0 \) and \( z \to -0 \), respectively. In the absence of an equilibrium flow, one recovers the dispersion relation derived by ER90.

#### 7.2.1. Small wavenumber approximation

We can have an outstanding analytical progress, giving a great insight into the coupling problem, in the long wavelength approximation. Following ER90 and introducing \( \Omega^2 = \omega^2 / (g k) \), \( \Omega_D^2 = \omega_D^2 / (g k) \), \( K = k z_0 \) and \( \Lambda = (m+1) c_s^2 / (\gamma p c_s^2) \), we can rewrite the dispersion relation (17) in the form
\[
2 a K \Omega_D^2 U(1-a, m+3, 2K) \frac{g k^2}{2 \omega_D^2} - (m+1) \Omega_D^2 - K(1+ \Omega_D^2) = \frac{\Lambda (\Omega^2 - K - 1) \Omega_D^2}{\chi (\Lambda + \varphi(1 + 1/\beta)(\Omega^2 - K)(1 + 1/\beta))}.
\]

First suppose that \( a \) does not tend to \(-1, 0 \) or a positive integer and \( \Omega \) does not tend to zero as \( K \to 0 \). We can determine the behaviour of the \( n = 1, 2, \ldots, p \)-modes in the limit of small \( K \):
\[
\Omega = \Omega_n + \frac{V}{c_s} \sqrt{\frac{\gamma p}{2 (m+1)}} (2K)^{1/2} + \frac{\gamma c \Gamma(1+n+m)}{\Gamma(1+n+m) \Gamma(2+m)} \times \left[ \frac{m+1}{\gamma p} + \frac{(m+1 - m)}{\gamma p} \right]^{-1} \left( 1 - \frac{\Omega_n^4 - 2 \Omega_n^2 (2K)^m + 1}{\Gamma(n) (\gamma c + 2 \beta) \Omega_n (1 - \Omega_n^2)} \right)^{m+1}
\]
where \( \Omega_n \) is determined from
\[
\frac{m+1}{\gamma p} \Omega_n^2 + \left( m - \frac{m+1}{\gamma p} \right) \frac{1}{\Omega_n^2} - m + 2 = 2(n-1)
\]
for \( n = 1, 2, \ldots, p \)-modes.

There remains the case \( a \to -1 \) corresponding to the \( f \)-mode. The frequency behaviour of the \( f \)-mode in the limit of small \( K \) is given by the formula
\[
\Omega = 1 + \sqrt{\frac{\gamma p}{2 (m+1)}} \frac{V(2K)^{1/2}}{c_s} + \frac{\gamma c (2K)^m}{2 \Gamma(m+2) (2 \beta + \gamma c)}.
\]
Figure 4. Cyclic frequency shift $\Delta_\nu = \nu(l, B, V) - \nu(l, B, 0)$ (\mu Hz) in the absence of a magnetic field ($B = 0$) as a function of the spherical harmonic degree $l$ and the flow $V$ (km/s) for the f-mode and the $n = 1$ p-mode.

From (20) and (22) we see that in the small wavenumber limit the flow has a more dominant influence on p/f-mode frequencies than the magnetic field.

7.2.2. Numerical evaluations and discussion

To solve the dispersion relation (17) numerically, we adopt a marginally stable stratification of the convection zone by taking $m = 1/\gamma_p - 1$ and suppose that $\gamma_p = \gamma_c = 5/3$. The temperature in the isothermal atmosphere is assumed to correspond to that at the temperature minimum, yielding $T_p = T_c = 4170$ K. From the continuity condition for the total equilibrium pressure across the interface we arrive at the following expression for the ratio $\beta = c_{sc}^2 / v_A^2 = \gamma_c (2\mu p_B - B^2) / (2B^2)$, where the pressure $p_B = 86.82 N/m^2$. The horizontal wavenumber $k$ is related to the spherical harmonic degree by the formula $k = \sqrt{(l + 1)/R_{Sun}}$, where $R_{Sun} = 696$ Mm. The results are presented in terms of spherical harmonic degree $l$ and cyclic frequency $\nu$ (related to the angular frequency by the relation $\omega = 2\pi \nu$).

In Figs. 4a-b the frequency difference $\Delta_\nu = \nu(l, B, V) - \nu(l, B, 0)$ (\mu Hz) expressing the flow effect in the absence of a magnetic field ($B = 0$), is plotted as a function of the spherical harmonic degree $l$ and the flow $V$ (km/s) for the (a) f-mode and (b) $n = 1$ p-mode, respectively. The flow $V$ changes in the interval $(-0.1, 0.1)$. The kinetic energy density of the f-mode in the field-free case is given by the expression $p_H / (2RT_c) \exp{(\gamma_c g / c_{sc}^2 - 2k)} z$. To satisfy the finite kinetic energy density condition as $z \to -\infty$, we require that $k < \gamma_c g / (2c_{sc}^2)$ (see also ER90). This gives the condition $l < 3550$. The p-modes leak into the atmosphere after reaching their cut-off frequency. In Figs. 4a-b the harmonic degree $l$ is continued until the cut-off frequency is reached (for p-modes) or until the finite kinetic energy condition is fulfilled (for the f-mode). For $n = 1$ p-mode and for the f-mode the magnitude of $\Delta_\nu$ increases with $l$ and $V$, starting to decrease at high degrees. The largest frequency shift ($\approx 30$ \mu Hz) occurs at about $l = 2000$ when $V = \pm 0.1$ km/s for the f-mode. The frequency shift $\Delta_\nu$ for the $n = 1$ p-mode reaches its maximum value ($\approx 20$ \mu Hz) at $l = 1100$ when $V = \pm 0.1$ km/s.

In Figs. 5a-b the frequency difference $\Delta_\nu = \nu(l, B, V) - \nu(l, B, 0)$ (\mu Hz) expressing the flow effect in the presence of a uniform magnetic field with $B = 30$ G, is plotted as a function of $l$ and $V$ (km/s) for (a) the f-mode and (b) $n = 1$ p-mode, respectively. Note that in the presence of a uniform magnetic field there is no cut-off frequency. Comparing Figs. 4 and 5 one can conclude that the flow influence on both p- and f-modes is slightly stronger in the presence of the given magnetic field. The same behaviour for p-modes can be observed in Figs. 6a-c, where this effect becomes more pronounced with increasing radial order $n$. Also observe from Figs. 5a-c that in the presence of the magnetic field the magnitude of the frequency shift continues to increase for high values of $l$, unlike the case with no magnetic field.

A similar equilibrium, though without flows was used by ER90 to study the effect of a chromospheric magnetic field on the p- and f-modes. Comparing the results of ER90 with our results, we observe that the flow influence on the p- and f-modes is more significant than the magnetic field influence in the small wavenumber limit. However this is not true for an arbitrary wavelength. The magnetic field influence might be stronger than the flow influence or vice-versa, depending on their characteristic values. Frequencies with higher $l$ are more sensitive to atmospheric effects and hence the magnitude of the frequency shift caused by the magnetic field grows more rapidly.

Measurements of p-mode frequencies (Braun & Fan 1998) show that there is a frequency shift between poleward- and equatorward-travelling waves measured over solar latitudes $20^\circ - 60^\circ$, which is consistent with the Doppler effect of a poleward meridional flow on the order of $10$ m/s. The measurements show frequency shifts on the order of $1 - 3$ \mu Hz for about $l \approx 477$. In order to compare our results in Figs. 6a-b we have plotted the frequency shift $\Delta_\nu = \nu(l, B, V) - \nu(l, B, -V)$ for $V = 10$ m/s as a function of the frequency $\nu(l, B, 0)$. We have taken two values of the harmonic degree: $l = 70$ and $l = 470$. Cases with $B = 0$ and 10 G are considered.
Figure 6. Cyclic frequency shift \( \Delta \nu / \nu = \nu(l, B, V) - \nu(l, B, -V) \) (\( \mu \text{Hz} \)) with \( V = 10 \text{ m/s} \) and \( B = 0 \) (dashed lines), \( B = 10 \text{ G} \) (solid lines) as a function of the frequency \( \nu(l, B, 0) \).

Figure 6b shows the frequency shift predicted by our simple model in good agreement with the measurements indicated above. However, these measurements are done only for low frequencies. Our model predicts a decrease in the magnitude of the frequency shift with increasing radial order \( n \) and cut-offs at high frequencies when no magnetic effects are taken into account.

7.3. Weak atmospheric magnetic field

In the lower atmosphere (\( z < 0 \)) the steady flow is absent and a uni-directional magnetic field \( B_0 = (B_0(z), 0, 0) \) is present. The magnetohydrostatic equation is given by Eq. (13), while perturbations are governed by Eq. (14) - (15). We again assume the transitional layer to be isothermal with the temperature equal to that at the top of the field-free medium, i.e., in \( z < 0 \) we take \( c_s(z) = c_0 \). Also we suppose that the Alfvén speed is constant. With these assumptions the plasma \( \beta = c_s^2/v_A^2 \) parameter is constant and the governing equation reduces to (e.g., CR89)

\[
\frac{d^2 u_z}{dz^2} + \frac{1}{H_B} \frac{du_z}{dz} + A_{uz} = 0,
\]

where

\[
A = \frac{(\Gamma - 1)k^2g^2 + (\omega^2 - k^2c_0^2)(\omega^2 - k^2c_0^2)}{(c_0^2 + v_A^2)(\omega^2 - k^2c_0^2)},
\]

\[
\Gamma = \frac{2\beta \gamma}{2\beta + \gamma}, \quad H_B^{-1} = \frac{\rho_0}{\rho_0} = \frac{\gamma g}{c_0^2},
\]

are the magnetically modified adiabatic exponent and pressure scale height, respectively (in the absence of a magnetic field, \( \Gamma = \gamma \) and \( H_B = H_0 \)). The solutions can be written in the the form

\[
\exp \left[ \frac{z}{2H_B} \left( -1 \pm \sqrt{1 - 4AH_B^2} \right)^{1/2} \right].
\]

We suppose that \( 4AH_B^2 < 1 \) and choose the plus sign. This corresponds to evanescent disturbances in the chromosphere. So we have \( u_z = \exp(\lambda z) \) where

\[
\lambda = \frac{1}{2H_B} \left[ -1 + \sqrt{1 - 4AH_B^2} \right], \quad z < 0.
\]

Figure 7. Cyclic frequency \( \nu = \omega/(2\pi) \) in mHz for \( f \)-mode (with \( k = 2.7 \text{ Mm}^{-1} \)) and \( n = 1 \) (with \( k = 1.3 \text{ Mm}^{-1} \)) p-mode as a function of flow in km/s and no magnetic field.

Figure 8. Frequency difference \( \Delta \nu \equiv \nu(V) - \nu(0) \) in \( \mu \text{Hz} \) with \( V = 1 \text{ km/s} \) and no magnetic field for the \( f \)-mode and the \( n = 1 \) p-mode as a function of the wavenumber \( k \) (\text{Mm}^{-1}).

7.4. Boundary conditions and the dispersion relation

Again, from continuity of the normal component of Lagrangian displacement (\( -iu_z/\omega \) in \( z < 0 \), and \( -iu_z/\omega_D \) in \( z > 0 \)) across \( z = 0 \), and from continuity of the equilibrium total pressure across \( z = 0 \) we have the dispersion relation for the case of weak field approximation

\[
2k\omega_D^2c_0^2 \frac{U(-a, m + 2, 2kz_0)}{U(-a, m + 2, 2kz_0)} + \gamma g\omega_D^2 - k\omega_D^2c_0^2 -\frac{gk^2c_0^2}{gk^2c_0^2 + (c_0^2 + v_A^2)(\omega^2 - k^2c_0^2)}.
\]

In the absence of flow \( \omega_D = \omega \) and (24) reduces to the dispersion relation found CR89.

7.4.1. Limit of long wavelength

It is convenient to treat the \( p \)- and \( f \)-modes separately.

\( p \)-modes:
To obtain a correction to the eigenfrequency, set
\[ \Omega_D^2 = \Omega_n^2 + pv_1(2kz_0)^s + pv_2 + \delta_0(2kz_0)^{s+\frac{3}{2}}, \]  
(25)
with \( pv_1, pv_2, \delta_0 \) and \( s \) to be determined. After some algebra, we obtain \( s = m + 3/2 \) and
\[ pv_1 = -\frac{2\sqrt{2} \Gamma(1 + m + n)(\Omega_n^4 + 1)V}{\Gamma(m + 1)\Gamma(m + 2)\Gamma(n)c_0 m \sqrt{m \Omega_n}(\Omega_n^4 - 1)}, \]  
(26)
\[ pv_2 = \frac{c_0^2 m^2 \Omega_n^2 (\Omega_n^4 - 1)^2}{\Gamma(m + 1)\Gamma(m + 2)\Gamma(n)c_0 m \sqrt{m \Omega_n}(\Omega_n^4 - 1)^2}, \]  
(27)
and
\[ \delta_0 = -\frac{\Gamma(1 + m + n)}{m(m + 1)\Gamma(m + 1)\Gamma(m + 2)\Gamma(n)} \times \left[ \frac{a_0 \Omega_n^8 + a_4 \Omega_n^4 + a_0}{\Omega_n^8 (\Omega_n^4 - 1) - \frac{m \Omega_n^2}{m + 2}} \right]. \]  
(28)

Thus,
\[ \Omega_D = \Omega_n + \frac{pv_1}{2\Omega_n} (2kz_0)^{m+\frac{3}{2}} + \frac{pv_2 + \delta_0}{2\Omega_n} (2kz_0)^{m+2}, \]  
(29)

or
\[ \Omega = \Omega_n + \frac{V}{c_0 \sqrt{2m}} (2kz_0)^{\frac{s}{2}} + \frac{pv_1}{2\Omega_n} (2kz_0)^{m+\frac{3}{2}} + \frac{pv_2 + \delta_0}{2\Omega_n} (2kz_0)^{m+2}, \]  
(30)
with \( pv_1, pv_2 \) and \( \delta_0 \) defined above.

**f-mode:**

Carrying out a similar analysis we can write the solutions to the dispersion relation for the \( f \)-mode in the form
\[ \Omega_D^2 = 1 + f v_1(kz_0)^s + (f v_2 + f_0)(kz_0)^{s+\frac{3}{2}}, \]  
(31)
with \( f v_1, f v_2, f_0 \) and \( s \) to be determined. After some algebra we obtain
\[ fv_1 = -\frac{2m+2V}{c_0 \sqrt{m \Gamma(m + 2)}}, \]  
\[ fv_2 = \frac{2m+1V^2}{c_0^2 m \Gamma(m + 2)}, \]  
\[ f_0 = \frac{(1 + 2\beta)2^m}{\beta(1 + 2\beta)\Gamma(m + 2)}, \]  
(32)
and
\[ \Omega = 1 + \frac{V}{c_0 \sqrt{m}} \sqrt{kz_0} - \frac{2m+1V(kz_0)^{m+\frac{3}{2}}}{c_0 m \sqrt{m \Omega_n}(\Omega_n^4 - 1)} + \frac{2^m}{\Gamma(m + 2)} \times \left[ \frac{V^2}{c_0^2 m} + \frac{(1 + 2\beta)}{2\beta(1 + 2\beta)} \right] (kz_0)^{m+2}. \]  
(33)

7.4.2. Numerical results for weak fields

For numerical investigations we have taken the sound speed in the isothermal atmosphere to correspond to that at the temperature minimum, yielding \( c_0 = 6.76 \text{ km/s} \) for adiabatic index \( \gamma = 5/3 \). The scale height \( H_0 \) is then \( 100 \text{ km} \) and the polytropic index \( m \) is \( 3/2 \). For conditions typical of the temperature minimum, we may take \( \beta = (180/B_0)^2 \), where \( B_0 \) is the field at the base of the magnetic atmosphere, measured in gauss (see also Campbell & Roberts 1989). Figures 7–10 are obtained by solving the full dispersion relation for the weak magnetic field case. Cut-offs in the figures are denoted by dotted lines.

8. RESONANT COUPLING OF SOLAR OSCILLATIONS TO THE ATMOSPHERE

So far we have mainly presented results of a very simple solar model, where we pointed out that both steady states and atmospheric magnetic fields could be important when evaluating helioseismic observational data. One key feature of the boundary layer was that the profile of the magnetic field was selected such that the Alfvén speed was constant in the boundary layer. In the following we relax this condition on the Alfvén speed.

A three-layer model with an intermediate zone, where the magnetic field, together with the Alfvén speed, varies continuously from zero was introduced by Tirry et al.
The importance of the continuum is that global modes may interact resonantly to local boundary layer Alfvén and/or slow oscillations at the height where their frequency matches the frequency of the global mode, hence the model can be used to investigate the effects of resonant coupling between global modes and local atmospheric MHD oscillations. Pintér & Goossens (1999) discussed the case of parallel propagation in this model. For propagation parallel to the magnetic field, the global oscillation modes couple to slow continuum modes only and this was found to occur for a rather large range of realistic parameters. In addition to the damping of global oscillation modes due to resonant absorption, it was also found that the interaction of global eigenmodes with slow continuum modes leads to an unanticipated behaviour in the global eigenmodes. The rather strange behaviour in the slow continuum involves the disappearance, appearance and splitting, and merging of global modes. Additionally, frequency shifts of global modes due to the magnetic field were examined. Pintér et al. (2005) extended the analysis to non-parallel propagation. They demonstrated that obliquely propagating global modes can couple also to local MHD Alfvén and slow continuum modes. They investigated the magnetic effects on global mode frequencies, especially the frequency shifts and damping rates caused by the resonant interaction with local slow and Alfvén waves. Due to the presence of a number of characteristic frequencies in the model, they found not only the $f$- and $p$-modes but also Lamb modes, with frequencies near the characteristic cut-off frequencies. Atmospheric gravity modes also appeared as solutions to the linear MHD equations. A theoretical study of the influence of the orientation of wave vector also raises the question how to measure the angle between the direction of the wave propagation and the atmospheric magnetic field lines. Although the number of observed oscillations is rapidly increasing, the resolution of detection has to be enhanced to obtain more features of the waves, such as the direction of their propagation with respect to the local magnetic field lines.

A straightforward application of the resonant coupling of acoustic oscillations in a steady state is the rotational splitting of helioseismic modes influenced by a magnetic atmosphere. Pintér et al. (2001b) studied the splitting of sectoral ($m = \pm l$) helioseismic eigenmodes. The Solar interior was in a steady state, with sub-photospheric plasma flow along the equator representing solar rotation. The Cartesian geometry employed restricted the study to sectoral modes with $l \geq 50$, which guarantees that the modes do not penetrate deeply into the solar interior and therefore experience an approximately uniform rotation. The mean increase of $\Delta v_{nlm}$ with $B$ for the $p$-modes they studied, for $l = 100$, is around 370 nHz, which is a 0.41% relative increase. For GONG and MDI data, the observational error of measuring $\Delta v_{nlm}$ due to rotational splitting is better than 0.25% (private communication with R. Howe). Hence, the effect obtained in the present model is on the verge of detectability, and ought to be detectable by combining a number of modes. On the other hand, there are other competing effects – such as those due to zonal flows – which are of about the same order. One possible way of helping to differentiate between the several competing shifts would be to evaluate the modelled rotational splitting for all $m$. This requires a move to spherical geometry.

Another application of the resonant coupling is the damping of helioseismic modes in a steady state (Pintér et al. 2001a). The frequencies and the line-widths of eigenmodes are affected by sub-surface flow and atmospheric magnetic fields. A key contribution to the effects comes from the universal mechanism of resonant absorption. When both atmospheric magnetic field and sub-surface flows are present, a complex picture of competition between these two effects was found. Their Table 1 shows the sensitivity of the line-width of the $f$- and $p_1$-modes to an equilibrium flow varying between $[-0.1c_s, 0.1c_s]$. The ratio $(\Gamma(V, l) - \Gamma(V=0, l)) V^{-1}$ is given for different values of $l$. The line-width of the $f$-mode increases or decreases linearly with $V$ in the interval $V \in [-0.1c_s, 0.1c_s]$, while that of the $p$-mode always increases, also linearly, in the given interval of $V$ and $l$. For larger $l$ the effects are more complicated, and the line-width of $p$-modes can also decrease with $V$. In addition, for larger values of $V$ the dependence becomes non-linear.

9. CONCLUSION

In the present paper I have tried to give an overview of studies on the coupling of global solar acoustic oscillations to the lower magnetised solar atmosphere. In analogy to critical layers in fluid dynamics, I introduced the
terminology of boundary layer for this narrow transitional layer (embracing the photosphere, chromosphere and transition region) between the solar interior and the solar corona. The effects of both magnetic fields and flows on the acoustic oscillations were investigated. The fields were split into their coherent and random parts. Each of these components has, in its own peculiar way, influences on the coupling of solar global oscillations to the atmosphere. At the moment it is hard to determine which effect is dominant or which one is less relevant. The Solar Dynamics Observatory (SDO) to be launched soon may shed light on this exciting and rapidly evolving question of coupling of solar acoustic oscillations to the solar atmosphere.

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