TRANSPORT EFFECTS IN THE EVOLUTION OF THE GLOBAL SOLAR MAGNETIC FIELD

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Abstract. The axisymmetric component of the large-scale solar magnetic fields has a pronounced poleward branch at higher latitudes. In order to clarify the origin of this branch we construct an axisymmetric model of the passive transport of the mean poloidal magnetic field in the convective zone, including meridional circulation, anisotropic diffusivity, turbulent pumping and density pumping. For realistic values of the transport coefficients we find that diffusivity is prevalent, and the latitudinal distribution of the field at the surface simply reflects the conditions at the bottom of the convective zone. Pumping effects concentrate the field to the bottom of the convective zone; a significant part of this pumping occurs in a shallow subsurface layer, normally not resolved in dynamo models. The phase delay of the surface poloidal field relative to the bottom poloidal field is found to be small. These results support the double dynamo wave models, may be compatible with some form of a mixed transport scenario, and exclude the passive transport theory for the origin of the polar branch.

1. Introduction

The time–latitude distribution of the axisymmetric component of the large-scale solar magnetic fields is known to have a pronounced poleward branch at higher latitudes (Stenflo, 1988, 1991, 1994; Stenflo and Güdel, 1988; Ribes and Bonnafond, 1990; Mouradian and Soru-Escaut, 1991). This branch is also present in the butterfly diagram of a number of tracers of the magnetic field such as quiescent filaments, polar faculae, or the coronal green line (Callebaut and Makarov, 1992; Makarov and Sivaraman, 1989; Leroy and Noens, 1983).

Several conflicting explanations exist for this polar branch. Noting that the separation latitude of the two branches, 30–40°, approximately coincides with the latitude where the radial differential rotation changes sign according to helioseismology, one group of theories interprets it as the surface reflection of a high-latitude poleward propagating dynamo wave, coexisting with the low-latitude equatorward wave (Gilman, Morrow, and De Luca, 1989; Belvedere, Pidatella, and Proctor, 1990, 1991). The Parker–Yoshimura rule of sign (Belvedere, 1985) then naturally leads to the correct directions of propagation if α is negative, as expected in the lower overshooting layer.
where the dynamo should operate. (In fact this is only so for certain latitudinal profiles of $\alpha$, cf. Schmitt, 1993.) The most modern version of these models, also incorporating some transport effects, is due to Rüdiger and Brandenburg (1995). In what follows we will refer to these models as double wave models.

An alternative approach to the problem of the origin of the poleward drift is to interpret it in terms of magnetic transport processes in the solar photosphere and convective zone. The common ancestor of all such models was the now classic Babcock–Leighton interpretation of the solar cycle (Babcock, 1961, Leighton, 1964). The Babcock–Leighton model was what one may call a mixed transport model in as much as the poloidal fields were brought to the surface in a concentrated form in active regions, and thereafter they were passively transported to the poles by transport processes (diffusion and meridional circulation). In the more recent mixed transport models (Wang, Sheeley, and Nash, 1991; Choudhuri, Schüssler, and Dikpati, 1995) meridional circulation plays the main role in transporting the weak fields to the poles, and is also supposed to be responsible for the equatorward drift of toroidal fields in the dynamo layer. (In the original Babcock–Leighton theory the equatorward branch was still supposed to be due to a dynamo wave.)

A third type of models may be called passive transport models. These are a variation of the mixed transport models where the “active source” of the weak fields in the form of active regions is also substituted by a passive transport mechanism. Dikpati & Choudhuri (1994, 1995) suggested such a scenario wherein the poloidal fields are brought to the surface at low latitudes by meridional circulation. In these models the origin of the equatorward drift is not specified; instead, it is just given as a lower boundary condition. The resulting butterfly diagrams for the poloidal field were in good agreement with the observations provided that the value of the turbulent magnetic diffusivity used in the calculations was artificially reduced to about 10 km$^2$/s, i.e. nearly 2 orders of magnitude lower than expected.

In the present paper we generalize the models of Dikpati & Choudhuri (1994, 1995) to include all turbulent transport effects such as anisotropic turbulent diffusion, turbulent pumping, and density pumping, as well as the meridional circulation. We construct models with realistic values of the transport coefficients in order to see whether the horizontal component of pumping can support the meridional circulation in producing a poleward migration despite the strong diffusive link between the surface and the bottom of the convective zone. (Such a possibility was suggested by Krivodubskij and Kichatinov, 1991 and Kichatinov, 1993.) Another motivation for this study is to check to what extent is the radial pumping able to concentrate the fields to the bottom of the convective zone, thereby restricting dynamo action to that layer, and to reduce the surface field strength to the
observed values, as suggested by Schüssler (1984), Petrovay and Szakály (1993), and Petrovay (1994a).

2. The model

We calculate the mean poloidal magnetic field in the solar convective zone as a function of space and time, assuming axial symmetry and antisymmetry to the equatorial plane. No dynamo is supposed to operate in the convective zone, i.e. the $\alpha$-effect is neglected. The dynamo layer below is not included in our computational volume; instead, the processes operating there are assumed to influence our model via the lower boundary conditions.

The equations of passive magnetic field transport with these assumptions are (Petrovay, 1994b):

$$\partial_t A_\phi = a_{\theta\theta} \partial_\theta^2 A_\phi + a_{\theta r} \partial_\theta \partial_r A_\phi + a_{rr} \partial_r^2 A_\phi + a_{\theta} \partial_\theta A_\phi + a_{r} \partial_r A_\phi + a_{0} A_\phi$$

$$\langle B_\theta \rangle = -\frac{A_\phi}{r} - \partial_r A_\phi$$

$$\langle B_r \rangle = \frac{\partial_\theta A_\phi}{r} + \cot \theta A_\phi$$

(1)

where

$$a_{\theta\theta} = b_{\theta\theta} = \beta_{\theta\theta}/r^2$$

$$a_{\theta r} = b_{\theta r} = 2\beta_{\theta r}/r$$

$$a_{rr} = b_{rr} = \beta_{rr}$$

$$a_{\theta} = (\gamma_\theta + \tilde{\gamma}_\theta - U_\theta)/r + (\beta_{\theta\theta} \cot \theta - \beta_{\theta r})/r^2$$

$$a_{r} = (\gamma_r + \tilde{\gamma}_r - U_r) + (\beta_{rr} + \beta_{\theta\theta} + \beta_{\theta r} \cot \theta)/r$$

$$a_{0} = \left[ (\gamma_r + \tilde{\gamma}_r - U_r) + \cot \theta(\gamma_\theta + \tilde{\gamma}_\theta - U_\theta) \right]/r$$

$$-\left[ \beta_{rr} + \beta_{\theta\theta} \cot^2 \theta + \beta_{\theta r} (1 + \cot \theta) \right]/r^2.$$  

(2)

$\mathbf{B}$ is the mean magnetic flux density, $A_\phi$ is the azimuthal component of the mean vector potential, $\mathbf{U}$ is the velocity of meridional circulation, $\beta$ is the anisotropic turbulent diffusivity, and $\gamma$ and $\tilde{\gamma}$ are the normal and anomalous components of the pumping (incorporating both the density gradient and the gradient of turbulence intensity, Moffatt, 1983). Note that in another widespread notation $\gamma$ and $\tilde{\gamma}$ are expressed by components of the anti-symmetric and symmetric parts of the $\alpha$-tensor. For $\mathbf{U}$ we use the same expression as given in Dikpati and Choudhuri (1994). For the diffusivity and pumping standard mean-field theory expressions are used; the formulae are collected in the Appendix. For the evaluation of those expressions we use the UKX convective zone model (Unno, Kondo, and Xiong, 1985).
differential rotation is assumed to be independent of depth in the computational regime.

For the solution of Eq. (1) we employ the following boundary conditions. At the bottom of the computational regime \( A_\phi = A_\phi(\theta, t) \) was explicitly given, as mentioned above. For the other boundaries we set

\[
\begin{align*}
A_\phi &= 0 \quad \text{at } \theta = 0 \\
\partial_\theta A_\phi &= 0 \quad \text{at } \theta = \pi/2 \\
\partial_r A_\phi &= -A_\phi/r \quad \text{at the surface}
\end{align*}
\]

The latter condition implies that the field is vertical at the surface (Yoshimura, 1975; Brandenburg, 1994). Note that these simple boundary conditions are not necessarily unrealistic: the field is intermittent, and if most of the flux in the photosphere is present in flux tubes exceeding \( 10^{18} \text{Mx} \) then Eq. (1) should apply owing to the strong buoyancy of these tubes. If, on the other hand, the flux is mainly present in thinner tubes, then a potential field boundary condition would be better (cf. Petrovay, 1994b for a discussion of this problem). We performed some test runs with different forms of the boundary condition to find that this choice does not have a great influence on the results.

Equation (1) was solved by the alternating direction implicit (ADI) method on a 64×64 grid. The grid we use is generally uniform in both \( \theta \) and \( r \). However, immediately below the surface the scale heights of the density and of the turbulent velocity are very small. A grid uniform in \( r \) is unable to resolve this fine structure and the corresponding potentially important pumping effects. For this reason, in some calculations a non-uniform grid is employed, resolving these layers. (In practice this is realized by introducing a new variable \( \tilde{r} = c_1 \tanh(c_2 r - c_3) \), rewriting eq. (1) in \( \tilde{r} \), and solving it on a grid uniform in \( \tilde{r} \).)

3. Results

3.1. Stationary dipole

It is instructive to consider the case when the lower boundary condition corresponds to the vector potential at fixed \( r = 500 \text{Mm} \) of a time-independent dipole field centered on the solar center. The cross section of the resulting field configuration after reaching a stationary state is shown in Fig. 1. Field lines are nearly radial, as a consequence of diffusivity. The effects of rotation and meridional circulation are negligible (Fig. 2). The overall field strength varies only by a factor of 2 between the surface and the bottom of the convective zone. The horizontal field component, however, increases by three
orders of magnitude from the surface to the bottom, as a consequence of the pumping (Fig. 3). This is the only case where the effect of rotation is not quite negligible; besides, it is apparent that a proper resolution of the shallow layers reduces the surface value by nearly one order of magnitude. The overall field structure is however not greatly distorted by a uniform grid, neither here, nor in any of the other cases studied.

3.2. OSCILLATING DIPOLE

The case when the lower boundary condition is a dipole with a dipole moment oscillating sinusoidally around zero was also considered, in order to see whether transport effects can turn this “standing dynamo wave” into a migrating wave. The result was negative: owing to the overwhelming diffusivity, the surface field pattern simply reflects the pattern at the bottom with a slight time delay. This finding is contrary to the result of Kichatinov (1993) who in a simplified model found that the effect of pumping, while weak, is sufficient to turn a high-latitude standing wave component into a migrating wave. While differences in the details of the models may also contribute to the different results, we believe that perhaps the most important difference is that in Kichatinov’s model the region of flux transport spatially coincides with the region of the dynamo (i.e. it is a convection zone dynamo).
3.3. Migrating field

The lower boundary condition here was

\[ A_\phi = F(\theta, t; A_1, \Gamma_1, \lambda_1, 0) + F(\theta, t; A_2, \Gamma_2, \lambda_2, \delta_2) \]  

(10)

where

\[ F(\theta, t; A_i, \Gamma_i, \lambda_i, \delta_i) = A_i \left\{ 1 + \exp \left( \Gamma_i (\pi/4 - \theta) \right) \right\}^{-1} \times \cos [\omega t + 2\pi/\lambda_i (\pi/2 - \theta) + \delta_i] \]  

(11)

with \( \omega = 9.1 \cdot 10^{-9}/s \) the cycle frequency.

The butterfly diagrams at the surface essentially reflect those at the bottom of the convective zone (Figs. 4 and 5). A high-latitude poleward branch and a low-latitude equatorward branch can thus only be produced at the surface if such a pattern is given as input at the bottom. Figure 6 shows that the field lines do not show a complicated topology inside the convective zone, although they are not exactly vertical either as the time dependence does not allow sufficient time for diffusion to smooth them. As a consequence, the strong downwards pumping of the horizontal field component is to some extent also “felt” by the vertical component, and,
Figure 3. Strength of the horizontal field component as a function of depth at $\theta = 45^\circ$ for the stationary dipole model (solid). Dots: Same with a uniform grid (shallow layers not resolved). Dashed: Model with rotation and circulation switched off.

in contrast to the stationary dipole case, the overall field strength is also significantly concentrated to the bottom of the convective zone (Fig. 3).

4. Conclusion

In our axisymmetric model of the passive transport of the large-scale mean poloidal magnetic field in the solar convective zone we found that the latitudinal distribution of the field at the surface reflects the conditions at the bottom of the convective zone, i.e. in this regard the convective zone behaves as a “steamy window”. This is due to the fact that with realistic values of the transport coefficients diffusivity is prevalent over all other effects. Passive transport theories of the origin of the poleward branch of the solar butterfly diagram are thus not viable, while double wave models are supported by these results. Owing to the large diffusivity the phase delay of the surface poloidal field relative to the bottom poloidal field is minimal. Coupled with the approximately $\pi$ phase difference between the toroidal and poloidal field components for negative $\alpha$ value (Dikpati and Choudhuri, 1995), this could bring the double wave models to accordance with the observed phase relationship.
A mixed transport scenario of the kind proposed by Wang, Sheeley, and Nash (1991) is not ruled out (note that our Figs. 1 and 6 agree with the field structure suggested in that paper), but the large diffusivity implies that the photospheric patterns should pervade the whole convective zone, being continuously reprocessed through it. Note that in fact the active source used in those models may also extend to deeper layers of the zone, as emerging magnetic loops are thought to shred a significant amount of flux during their rise (Petrovay and Szákály, 1993; Moreno-Insertis, Caligari, and Schüssler, 1995; Petrovay and Moreno-Insertis, 1997).

The vertical distribution of the field (Figs. 3 and 7) is strongly concentrated to the bottom of the convective zone owing to pumping effects. This is in agreement with the earlier proposal of Schüssler (1984) and the one-dimensional results of Petrovay and Szákály (1993) and Petrovay (1994a). The effective pumping may contribute to restricting the dynamo to the overshooting layer and it can reduce the surface field by about an order of magnitude, thereby offering a solution to the overtly high surface poloidal field values found by Rüdiger and Brandenburg (1995) in their double wave model. Note that a significant part of this pumping effect occurs in a shallow subsurface layer, normally not resolved in dynamo models.
In summary, our results support the double wave models, may be compatible with some form of the mixed transport scenario, and exclude the passive transport theory.

In order to decide between the double wave and mixed transport models, two key issues should be solved in the future. First, note that in the mixed transport model the direction of migration is opposite for poloidal and toroidal fields but independent of latitude. Thus, if e.g. a clear signature of migrating high-latitude toroidal fields were found, this could solve the problem in either way, depending on the direction of migration. Harvey (1994) claims that high latitude ephemeral active regions show an equatorward migration, while Callebaut and Makarov (1992) claim that at least 50% of polar faculae (well known for their poleward drift) correspond to dipoles with a preferential east–west orientation, thus forming part of the toroidal field. It has even been claimed that the highest-latitude part of the sunspot butterfly diagram also shows a poleward drift (Becker, 1959). A clarification of this issue would clearly be important.

The role of causality relations between the two branches should also be clarified in both theories. Surges in low-latitude activity are followed by poleward surges of the high-latitude field, thus supporting the mixed transport scenario. On the other hand, the well-known fact that the level of
Figure 6. Field line configurations in a meridional plane at 4 equidistant cycle phases for the migrating double wave model.

Low-latitude activity in a sunspot cycle can be predicted from high-latitude fields at the end of the previous cycle suggests a causal relationship of the opposite sense. It is not impossible that such causal relations can be accommodated in both scenarios, but this remains to be seen.

Acknowledgements

This work was financed in part by the DGES project no. 95-0028 and by the OTKA under grant no. F012817.
Figure 7. Total field strength as a function of depth at $\theta = 45^\circ$ at different cycle phases for a migrating double wave model.

**Appendix**

**Expressions for the transport coefficients**

The UKX model assumed fixed anisotropy $\beta_{rr}/\beta_{\theta\theta} = 2$ for the diffusivity tensor. We assume that the effect of rotation on the diffusivity remains the same as for quasi-isotropic turbulence (Kichatinov, 1988):

$$
\begin{align*}
\beta_{\theta\theta} &= 3\beta(\phi_3 + \phi_2 \sin^2 \theta) & \beta_{\phi\phi} &= 3\beta \phi_3 \\
\beta_{rr} &= 3\beta(2\phi_3 + \phi_2 \cos^2 \theta) & \beta_{\theta r} &= -\frac{3}{2} \beta \phi_2 \sin(2\theta) \\
\beta_{\theta\phi} &= \beta_{\phi r} = 0
\end{align*}
(12)
$$

$$
\begin{align*}
\phi_2(\Omega_*) &= \frac{1}{4\Omega_*^2} \left( \frac{\Omega_*^2 + 3}{\Omega_*} \arctan \Omega_* - 3 \right) \\
\phi_3(\Omega_*) &= \frac{1}{4\Omega_*^2} \left( \frac{\Omega_*^2 - 1}{\Omega_*} \arctan \Omega_* + 1 \right)
\end{align*}
(13)
(14)

$$
\begin{align*}
\beta &= \frac{1}{2} x^2 \tau \\
\Omega_* &= \tau \Omega / \pi \\
\tau &= \text{Max} \{ H_p / x; 1000 \text{s} \}
\end{align*}
(15)
(16)
(17)
where $x$ is the r.m.s. radial component of the turbulent velocity in the UKX model, $H_P$ is the pressure scale height, and $\Omega = \Omega(\theta)$ is the surface plasma differential rotation law. Note that the expression of $\beta$ used is valid for non-cellular flows only (free random walk). The cellular nature of the supergranular flow on the solar surface may reduce the diffusivity there (Ruzmaikin and Molchanov, 1997), but even this reduced value is well in excess of 100 km$^2$/s and the reduction should only be confined to the photosphere, as a similar cellular flow is not expected in the deeper layers.

The general expressions of the pumping velocities for weak magnetic fields are

$$\gamma_\theta + \tilde{\gamma}_\theta = \left( K_{1}\frac{A_1^3-1}{A_1^3+1} + K_\rho \right) \beta_\theta d_\rho / \rho + K_v \frac{A_1^3}{A_1^3+1} \left( \partial_\theta \beta_\theta + 2 \beta_\theta r / r + K_v \partial_r \beta_\theta r \right) \tag{18}$$

$$\gamma_r + \tilde{\gamma}_r = \left( K_{1}\frac{A_2^3-1}{A_2^3+1} - K_\rho \right) \beta_r d_\rho / \rho + K_v \frac{A_2^3}{A_2^3+1} \partial_r \beta_r r + K_v \left( \partial_\theta \beta_\theta r + \beta_r r - \beta_\theta \theta \right) / r + K_{\text{im}} \frac{A_2^3}{A_2^3+1} x \tag{19}$$

$$3A_2^3 = \left[ 1 + \left( \phi_2 / \phi_3 \right) \sin^2 \theta \right]^{1/2} \tag{20}$$

$$1A_2^3 = 1 + \left[ 1 + \left( \phi_2 / \phi_3 \right) \sin^2 \theta \right]^{-1/2} \tag{21}$$

where $K_v = 0.6$, $K_\rho = 0.15$, $K_\beta = 0.55$ and $K_{\text{im}} = 0.1$ account for (small) higher-order and intermittency corrections (Petrovay, 1994b, Petrovay and Zsargó, 1998); these corrections are of minor importance. Note that the magnetic field strength does not exceed about 10 G at any point in our computational volume (if we assume that the surface fields are about 1 G as observed); so buoyancy effects were altogether neglected here.

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