Global Solar Magnetic Field Organization in the Outer Corona: Influence on the Solar Wind Speed and Mass Flux Over the Cycle

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Abstract

The dynamics of the solar wind depends intrinsically on the structure of the global solar magnetic field, which undergoes fundamental changes over the 11-year solar cycle. For instance, the wind terminal velocity is thought to be anti-correlated with the expansion factor, a measure of how the magnetic field varies with height in the solar corona, usually computed at a fixed height ($\sim 2.5 R_\odot$, the source surface radius that approximates the distance at which all magnetic field lines become open). However, the magnetic field expansion affects the solar wind in a more detailed way, its influence on the solar wind properties remaining significant well beyond the source surface.

We demonstrate this using 3D global magnetohydrodynamic (MHD) simulations of the solar corona, constrained by surface magnetograms over half a solar cycle (1989–2001). A self-consistent expansion beyond the solar wind critical point (even up to $10 R_\odot$) makes our model comply with observed characteristics of the solar wind, namely, that the radial magnetic field intensity becomes latitude independent at some distance from the Sun, and that the mass flux is mostly independent of the terminal wind speed. We also show that near activity minimum, the expansion in the higher corona has more influence on the wind speed than the expansion below $2.5 R_\odot$.

Key words: magnetohydrodynamics (MHD) – methods: numerical – solar wind – Sun: corona

1. Introduction

Relating in situ measurements of the solar wind and remote observations of the Sun and the solar corona has been a continuous effort for the past 60 years, starting with the pioneering work of Parker (1958). Within this standard wind model, the observed $\sim 1$ MK coronal temperature is able to dynamically expand a hydrodynamical solar atmosphere to supersonic speed around 400 km s$^{-1}$ at 1 au. The solar magnetic field has further been found to play a fundamental influence on the structure of the corona and the solar speed. The solar wind speed, density, magnetic field, and currents see changes following the solar sunspot cycle of 11 years (Neugebauer 1975). Through this cycle, the solar magnetic field observed at the photosphere changes its geometry and its amplitude with the minimum of activity, showing a dipolar configuration and the maximum of activity a quadrupolar or multipolar configuration (DeRosa et al. 2012).

With its out of the ecliptic orbit, the Ulysses spacecraft has revealed, over a full cycle, the 3D structure of the solar wind and its relation to the magnetic field geometry (McComas et al. 1998; McComas 2003). The two components of the solar wind, the slow (around 400 km s$^{-1}$ at 1 au) and the fast (around 800 km s$^{-1}$ at 1 au) are believed to originate nearby and away from closed coronal loops, respectively. In situ measurements have revealed that the mass flux in the solar wind at 1 au is between $2 \times 10^{8}$ cm$^{-2}$ s$^{-1}$ in the fast streams and $(3-4) \times 10^{8}$ cm$^{-2}$ in slow streams (see e.g., Schwenn 1983; Goldstein et al. 1996, for observations made with Helios and Ulysses, respectively). Modeling the separation of closed loops and open regions in the corona is thus crucial to predicting the solar wind characteristics at Earth orbit. Ulysses observations have also revealed that the radial magnetic field is constant with latitude (Smith & Balogh 1995; Smith 2011). This empirically validates that magnetic field measurements made in the ecliptic plane are a good proxy for the total flux emerging from the Sun.

The well-known potential field source surface model (PFSS; see Altschuler & Newkirk 1969; Schatten et al. 1969; Schrijver & DeRosa 2003) gives, in an extremely computationally efficient way, a structure for the coronal magnetic field from observed magnetograms. Assuming a current free corona, up to the source surface radius where the field becomes radial, the PFSS model yields the location and the area of coronal holes as well as the open regions polarity (and thus the location of the heliospheric current sheet). The source surface radius is usually set around $2.5 R_\odot$, even though the optimal value is likely to vary over the solar cycle, to match, for instance, the variation of the solar wind flux (Lee et al. 2011; Arden et al. 2014).

The PFSS model yields the value of the expansion factor $f_{\exp}$, i.e., how a given flux tube expands deviates from a pure radial field, up to the source surface. This parameter has been observed to be anti-correlated with the terminal wind speed (Wang & Sheeley 1990). The PFSS model has thus been used, combined with other models able to propagate the solution to 1 au, to predict the solar wind speed and IMF polarity from magnetograms, notably within the Wang–Sheeley–Arge (WSA) model (Arge & Pizzo 2000; Arge et al. 2003). The relation obtained between the expansion factor at the source surface and the terminal wind speed in the WSA model needs, however, to be calibrated and corrected with observations. The PFSS model indeed fails to reproduce the interaction between the magnetic field and flow dynamics, yielding, for instance, a latitude-dependent profile of the radial magnetic field.

Consequently, full magnetohydrodynamics (MHD) models are necessary to obtain a more coherent picture. Usmanov (1993) first used observed magnetograms within a full MHD model to analyze the properties of the solar corona and wind. (Mikić et al. 1999) have shown that MHD models constrained by photospheric magnetograms yield a good description of the
structure, shape, and size of coronal loops as they appear during eclipses. The radial field is also observed to be independent of latitude in these models as a consequence of the latitudinal balance between the thermal and dynamic pressure of the wind and the magnetic forces, realized at $\beta \geq 1$, in the higher corona (see Usmanov et al. 2000 for a detailed discussion). This effect has an influence on the expansion of the field lines beyond the canonical value of the source surface, thus it cannot be captured by the PFSS. Pinto et al. (2016) studied, in an axisymmetric setup, the expansion of the field lines over the solar cycle and showed that the link between the expansion in the whole domain and the wind terminal speed clearly depends on the cycle phase.

The purpose of this paper is to study the 3D structure of the solar corona and the heliospheric magnetic fields at a few solar radii from the Sun, over the solar cycle, and to yield a new understanding of the structure of flux tube that is obtained with MHD models. Although many papers have already been dedicated to this topic, they focus mainly on the structure of closed loops or open regions close to the Sun. For instance, Riley et al. (2006) and Réville et al. (2015b) have shown that below the source surface, the agreement between MHD simulations and the PFSS were acceptable, both providing similar values of the open magnetic flux (if the value of the source surface is well chosen; see Réville et al. 2015b). Deviations beyond the source surface are, however, significant and have consequences on the solar wind speed predictions as shown recently by Cohen (2015). The geometry of the field lines evolve continuously from the surface to and within the interplanetary medium. It is misleading to think that they become strictly radial near the Sun, whether it is beyond a source surface radius around $2.5 R_\odot$ or even beyond the sonic or the Alfvén point, located further on. We demonstrate this performing MHD simulations using as an input WSO magnetograms, one per year between 1989 and 2001. We observe a reorganization of the field that is taking place between the surface and about 10 solar radii, i.e., far beyond the typical value of the source surface used for the Sun.

In Section 2, we detail our numerical model and then describe the 13 simulations we made using Wilcox magnetograms to cover half a 22-year solar cycle for the period 1989–2001. The global structure of the corona is described in Section 3, and the average location of what we define as an optimal source surface radius, the sonic point, and the Alfvén point are analyzed over the cycle. Section 4 shows that the latitude independent $B_c$ established in our simulations is responsible for an enhanced field line expansion beyond the optimal source surface radius. In Section 5, we look at the influence of the expansion on the wind speed, and find that the well known anti-correlation between the total expansion factor and wind speed is reversed between the source surface and the higher corona during solar minimum. We also show the evolution of the mass flux with height and underline the importance of the expansion both below and beyond the source surface to match observations. We discuss these results and set the limitations of our model in Section 6.

2. Numerical Setup

2.1. MHD Model

The numerical simulations reported here were done using the PLUTO code (Mignone et al. 2007). We follow the exact same procedure as in Réville et al. (2016a) and solve the time-dependent ideal (non-resistive) MHD equations:

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \quad (1)$$

$$\frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot (\rho \mathbf{v} \mathbf{v} - \mathbf{BB} + \mathbf{Ip}) = -\rho \nabla \Phi + \mathbf{a}, \quad (2)$$

$$\frac{\partial}{\partial t}(E + \rho \Phi) + \nabla \cdot ((E + \rho \Phi) \mathbf{v} - \mathbf{B} (\mathbf{v} \cdot \mathbf{B})) = \rho \mathbf{v} \cdot \mathbf{a}, \quad (3)$$

$$\frac{\partial \mathbf{B}}{\partial t} + \nabla \cdot (\mathbf{vB} - \mathbf{Bv}) = 0, \quad (4)$$

where $E \equiv \rho e + \rho v^2/2 + B^2/2$ is the total energy, $\mathbf{B}$ is the magnetic field, $\rho$ is the mass density, $\mathbf{v}$ is the velocity field, $p = p_b + B^2/2$ is the total (thermal plus magnetic) pressure, and $I$ is the identity matrix.

The potential $\Phi$ accounts for the gravitational attraction of the star and $\Phi$ is a source term that contains the Coriolis and the latitudinal balance between the thermal and dynamic forces. It has shown recently by Cohen (2015) that the value of the source surface is well chosen; see Réville et al. 2016a, for a more detailed description of the expansion on the wind speed, and outcome of the simulations performed in this paper put forward in this paper.

The domain is a 3D Cartesian box of $60 R_\odot$ cubed, with 448 cells in each direction. The grid is refined in the domain $[-1.5 R_\odot, 1.5 R_\odot]$, as specified in the domain (see Réville et al. 2016a, for a more detailed
discussion on the boundary conditions). In this region, the surface magnetic field is imposed from observations and reconstructed with a PFSS at $t = 0$.

### 2.2. WSO Magnetograms

Our study is based on WSO magnetograms (Scherrer et al. 1977). These synoptic magnetograms identify the magnetic field in the line of sight with the radial component of the surface field, and we take from the work of DeRosa et al. (2012) the spherical harmonics decomposition of $B_r$ for each Carrington rotation:

$$B_r(1 R_S, \theta, \phi) = \sum_{\ell,m} a_{\ell,m} Y_{\ell,m}(\theta, \phi),$$

where the $Y_{\ell,m}$ are the scalar spherical harmonics functions. The WSO has continuously measured the solar magnetic field since 1976 (CR 1642). Figure 1 shows the typical properties of the solar magnetic field for the past 40 years. The magnetic energy at the surface, $\sum a_{\ell,m}^2$, oscillates, with a period of $\approx 11$ years, half the full magnetic cycle that includes the reversal. It reaches three maximums around 1980, 1991, and 2003, with a lower amplitude for the last cycle, which was also observed counting sunspots at the surface.

In Figure 1, we also show the parameter

$$f_{\text{disp}} = \frac{a_{1,0}^2 + a_{1,1}^2}{\sum_{\ell,m} a_{\ell,m}^2},$$

which quantifies the energy of the dipole component (axisymmetric and non-axisymmetric) over the total energy of the surface radial field in its spherical harmonics decomposition. Figure 1 illustrates the anti-correlation between the total magnetic field energy and the energy contained in the dipolar components over the whole period. Equivalently, we say that the Sun is almost dipolar at minimum of activity and multipolar components over the whole period. Equivalently, we say that magnetic field geometry is mostly axisymmetric dipolar component is, in effect, the dominant mode during those epochs.

### 3. Global Structure of the Corona

The results of the simulations are shown in Figure 2 for six simulations out of our 13 cases. The steady state is characterized by a balance between the outward thermal pressure, which creates the solar wind, and the magnetic forces that can contain the plasma in co-rotation with the Sun within closed loops. Using observed magnetograms all over the cycle creates a time-varying and complex distribution of loops and open regions.

Around minimum of activity, the structure is mostly dipolar, and the Alfvén surface—where the solar wind reaches the Alfvén speed $v_A = B/\sqrt{\mu_0 \rho}$—shown in gray shades, has its usual bilobal structure. Each lobe corresponds to the northern and southern hemispheres coronal holes. The axisymmetric dipolar component is, in effect, the dominant mode during those epochs. Near maximum of activity (i.e., for CR 1824 and CR 1954), we observe a more complex structure with smaller lobes in the Alfvén surface that corresponds to coronal holes and polarity connected open regions. The Alfvén point always touches the cusp of the helmet streamers, and is thus the closest to the star near the current sheets. In Figure 2, the polarity of the field lines is shown as reddish for positive and as yellowish for negative. When the magnetic field geometry is mostly dipolar, we observe a clear separation between positive and negative polarities—separation embodied by the heliospheric current sheet—located in the equatorial plane near the activity minimum. This clear separation disappears around maximum where current sheets can be located at all latitudes.

In Figure 3, we show the evolution of the total mass-loss rate $M = \int \rho v \cdot dS$. The mass-loss rate varies by some 20%,
between $M_2 \times 10^{-14} \, M_\odot \, \text{yr}^{-1}$ and $M_3 \times 10^{-14} \, M_\odot \, \text{yr}^{-1}$, within the observational constraints. The mass-loss minimum is located around 1990. The maximum is reached in year 1993, between solar activity maximum and solar activity minimum, then slowly decreases until 1999 and raises again in 2000. The correlation between the mass-loss variations and the structure of the solar magnetic field is subtle, but $M$ is globally anti-correlated with the magnetic energy (see Figure 1). Its variations match the observations of the Ulysses spacecraft as shown in McComas (2003), McComas et al. (2008). The in situ data starts with a minimum value of $M$ around $2.3 \times 10^{-14} \, M_\odot \, \text{yr}^{-1}$ in 1991 and reaches a maximum of $3.0 \times 10^{-14} \, M_\odot \, \text{yr}^{-1}$ in 1992–1993. The mass loss then slowly decreases until the new solar activity maximum phase in 2000, when it starts increasing again. We computed these values from the dynamical pressure, $\rho v^2$, and wind speed given in McComas et al. (2008), assuming a spherical symmetry.

Given the topology changes of the solar surface magnetic field observed over the cycle phase, we expect the variation of the overall size of the Alfvén surface to be correlated with it. In Réville et al. (2015a), we showed that the location of the Alfvén was strongly influenced by the topology of the surface magnetic field and that it was a decreasing function of the complexity: the degree $\ell$ of the component in the spherical harmonics decomposition of the field (see also Finley & Matt 2017). Previous studies such as Pinto et al. (2011) were predicting a factor four between minimum and maximum of activity, due to the strong decreases of the dipole component at maximum. The strong modulation of $f_{\text{dip}}$ (more than a factor 10) would therefore suggest a significant modulation of the Alfvén radius. However, as shown in Table 1 and in Figure 3, this variation is, at most, 24% in average (the variation measured for all $(\theta, \phi)$ is larger). The clear anti-correlation

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**Figure 2.** 3D rendering of six of our MHD solutions of the solar wind using WSO maps. The fields lines are shown red or yellow depending on the polarity of the field. The Alfvén surface is shown in translucent gray. The surface radial magnetic field is shown using color contour maps. CR 1824 and CR 1954 represent clear cases of solar maximum with a non-axisymmetric dipole and a complex multipole structure, respectively.

**Figure 3.** Evolution of the mass-loss rate in the simulations (in red; left y-axis) and of the optimal source surface radius, the average spherical sonic point, and the average spherical Alfvén point over the cycle (right y-axis).
between $f_{\text{dip}}$ and the magnetic energy shown in Figure 1 is certainly the reason for this small variation and both the (increase) decrease of the dipole and the (decrease) increase of the magnetic energy must be looked at to understand the variation we obtain.

For instance, the maximum extension of the Alfvén surface is not exactly located at the minimum of activity, but rather between maximum and minimum for CR 1850 when the dipole component is regaining strength and the surface field has not yet fallen into its lowest energy levels. Similarly, the smallest average Alfvén surface are reached for CR 1954 and CR 1967, two years before the maximum of activity, when the dipole component and the magnetic energy are both weak. Also, the sharp drop observed for CR 1837 is due to the drop in amplitude around year 1991, shown in Figure 1 with a blue arrow.

Hence, taking into account the right variation of the field’s energy, or amplitude, is crucial to correctly capture the variation of the Alfvén surface over the cycle. This is likely the main difference between these results and the results of Pinto et al. (2011), where the surface magnetic field was taken from a kinematic dynamo code and was probably not recovering the right amplitude variation. They also find the mass loss to be maximum during maximum of activity. The integrated mass loss is strongly influenced by the spatial distribution of slow, dense wind, and the value at maximum of activity is probably enhanced by the axisymmetry of their model. This feature has been advanced in other works such as Wang (1998), which based its estimates on PFSS extrapolations and thus might not be fully consistent at estimating the mass flux (see Section 5). Nevertheless, the wind driving is likely to be an increasing function of the surface magnetic field energy (through the Alfvén waves energy flux for instance, see e.g., Suzuki 2006; Cranmer et al. 2007; Pinto & Rouillard 2017). This would certainly amplify the variation of the Alfvén surface over the cycle, with a clearer peak near activity minimum, and we can consider our results, obtained with fixed acceleration processes for the wind, as a minimum case.

In addition to the average Alfvén radius, Figure 3 shows the evolution of the average critical point, where the wind becomes supersonic, and of the optimal source surface radius $r_{\text{ss, opt}}$. The latter is defined from our previous work (Réville et al. 2015b), as the source surface radius that makes a PFSS model match the open flux of the simulation. We see that $r_{\text{ss, opt}}$ is always smaller than the Alfvén radius and its evolution follows the same variation than the average Alfvén radius. The peak in both quantities is for CR 1850. A slow decrease then follows from around $7 R_\odot$ to $5 R_\odot$ for the Alfvén radius and from $3 R_\odot$ to $2 R_\odot$ for $r_{\text{ss, opt}}$. This variation is in qualitative agreement with the variation suggested in Lee et al. (2011) and Arden et al. (2014).

It is interesting to see that the sonic point location is steadier over the cycle, located around $5.5 R_\odot$ (see Table 1). Its location is mostly controlled by the thermodynamics we impose for our simulation and which remain constant in all our sample. The sonic point is always located between the optimal source surface and the Alfvén point. The interaction between the magnetic field and the solar wind flow around the sonic point are crucial to understand to properties of the solar wind at 1 au (see e.g., Kopp & Holzer 1976; Wang & Sheeley 1991; Velli 2010), and we must investigate what is occurring in the higher corona, i.e., beyond the optimal source surface.

### 4. Magnetic Field Organization in the Higher Corona

Figure 4 shows the radial magnetic field as a function of the latitude in the MHD model and the PFSS model for six Carrington rotations (same as in Figure 2) at different heights: $3 R_\odot$ and $10 R_\odot$. The colors show two different longitudes ($\phi = 0, \pi$). The profiles are similar for epochs far from the maximum of activity when the magnetic field is globally axisymmetric, while they vary significantly near maximum. In the MHD model, at a height of $3 R_\odot$—just above the optimal source surface—the radial magnetic field varies strongly with latitude. Changes of polarity can be identified when $B_r \rightarrow 0$, and we see $B_r$ decaying slowly on both sides of the current sheet, especially near minimum of activity. Higher in the corona, at $10 R_\odot$, the radial field forms a plateau at all latitudes, but near current sheets where it decays extremely sharply. The standard deviation of $B_r(\theta)$ is about 10% at $10 R_\odot$, against 33% at $3 R_\odot$ at activity minimum. This latitude independent radial magnetic field is in agreement with Ulysses observations far away from the Sun and should be recovered in every MHD solar wind model.

This feature is not in general obtained with the PFSS model. In Figure 4, radial magnetic field profiles of potential extrapolations made with the same surface magnetic field than the corresponding MHD models, and with the optimal source surface (see Table 1), are shown with dashed lines. The PFSS model has by definition a fixed structure in space beyond the source surface. In both panels, the structure of the radial magnetic field is consequently the same: strongly modulated over latitude and maximal in the core of coronal holes. The standard deviation of $B_r(\theta)$ is in this model and for activity minimum around 66%.

Smith & Balogh (1995) discussed this feature and noted that any latitudinal gradient of $B_r$ existing in the corona would tend to vanish through the action of the latitudinal term of the Lorentz force (created by the azimuthal current). It is worth noting that in general, the PFSS model does not yield a force free or even a current free magnetic field beyond the source surface.\(^2\) Given the case of a pure axisymmetric dipole, the latitudinal term in the Lorentz force beyond the source surface can be written as

$$\mathbf{J} \times \mathbf{B}_\theta = \mathbf{J}_q B_r - \mathbf{J}_r B_q,$$

where $\mathbf{J}_q$ and $\mathbf{J}_r$ are the azimuthal and radial components of the current density, respectively, $B_r$ and $B_q$ are the radial and azimuthal components of the magnetic field, and $\mathbf{J}$ is the current density. The Lorentz force is thus always equatorward. Put in other words, the gradient of $B_r$ (shown in Figure 4 for the PFSS model) would create a latitudinal field beyond the source surface.\(^2\) Given the case of a pure axisymmetric dipole, the latitudinal term in the Lorentz force beyond the source surface can be written as

$$\mathbf{J} \times \mathbf{B}_\theta = \mathbf{J}_q B_r - \mathbf{J}_r B_q = \left( \frac{1}{r} \frac{\partial}{\partial \theta} \right) B_r,$$

$$= \frac{1}{r^2} f (r_\odot)^2 \cos(\theta) \sin(\theta),$$

where $f (r_\odot) = 3 \mu / (2 r_\odot + 1/ r^2_\odot)$, and $\mu$ is the dipole moment. The $\theta$ component of the Lorentz force is thus always equatorward. Put in other words, the gradient of $B_r(\theta)$ shown in Figure 4 for the PFSS model would create a latitudinal magnetic pressure and disappear if left to evolve. To illustrate this process, we initialize a simulation of CR 1902 at the initial conditions of CR 1902 and follow the evolution of the latitudinal term of the momentum equation for point located at a height of $10 R_\odot$, until steady state is reached. Figure 5 shows the time evolution of the different terms.

When the simulation is initialized with $r_{\text{ss, opt}}$, the initial latitudinal Lorentz term is the dominant force and decreases to

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\(^2\) By definition, the field is current free below the source surface.
reach equilibrium with the pressure gradient and advection terms. When the simulation is initialized with a large \( r_{ss} \), the Lorentz force increases slightly but reaches a lower steady-state value than in the latter case. Note that the two steady-state are not exactly equivalent in Figure 5. The latitudinal profile of \( B_r \) is relatively close to the a PFSS solution (computed at \( r_{ss,\text{opt}} \)) shown in dashed lines. Further away, in the lower panels, \( B_r \) is constant with latitude—which is in better agreement with Ulysses observations—and deviates significantly from the PFSS model.

In Figure 6, we show the difference between the geometry of a given flux tube integrated from the surface in a PFSS solution and in the MHD simulations for two epochs (minimum and maximum of activity). For both models, field lines departing from the same footpoint at the base of the corona are identified by a single color. We see that field lines diverge in both model approximately at the optimal source surface (shown in dashed lines). Moreover, the difference is more important during minimum because of the more significant reorganization of the field lines that occurs beyond \( r_{ss,\text{opt}} \).

In Figure 4, we show the structure of the radial magnetic field on spheres of radii \( 3 R_\odot \) (upper panel) and \( 10 R_\odot \) (lower panel) at different epochs of the cycle as a function of latitude. Two longitudes \( (\varphi = 0, \pi) \) are represented with the color gradation. In the upper panels, the structure of \( B_r \) is relatively close to the a PFSS solution (computed at \( r_{ss,\text{opt}} \)) shown in dashed lines. Further away, in the lower panels, \( B_r \) is constant with latitude—which is in better agreement with Ulysses observations—and deviates significantly from the PFSS model.

### Figure 4
Structure of the radial magnetic field on spheres of radii \( 3 R_\odot \) (upper panel) and \( 10 R_\odot \) (lower panel) at different epochs of the cycle as a function of latitude. Two longitudes \( (\varphi = 0, \pi) \) are represented with the color gradation. In the upper panels, the structure of \( B_r \) is relatively close to the a PFSS solution (computed at \( r_{ss,\text{opt}} \)) shown in dashed lines. Further away, in the lower panels, \( B_r \) is constant with latitude—which is in better agreement with Ulysses observations—and deviates significantly from the PFSS model.

### Figure 5
Time evolution of the latitudinal terms in the momentum equation for a representative point of the domain, located in a coronal hole for CR 1902 at \( r = 10 R_\odot, \theta = \pi/4, \phi = 0 \). Time is expressed in average Alfvén crossing time in the domain, \( \langle \tau_A \rangle \). The plain lines represent the simulation initialized with \( r_{ss,\text{opt}} = 2.3 R_\odot \). We see the initial latitudinal Lorentz force diminishing to reach equilibrium with the pressure gradient and the advection term. Dashed lines illustrate the simulation initialized with \( r_{ss} = 15 R_\odot \), which starts with no currents and thus no Lorentz force.

### Figure 6
Difference between the magnetic field lines of a MHD simulation and a PFSS model at the optimal source surface for minimum of activity (left panel, CR 1902) and maximum of activity (right panel, CR 1954). For both models, field lines departing from the same footprint at the base of the corona are identified by a single color. We see that field lines diverge in both model approximately at the optimal source surface (shown in dashed lines). Moreover, the difference is more important during minimum because of the more significant reorganization of the field lines that occurs beyond \( r_{ss,\text{opt}} \).
The continuous expansion that occurs at various height in the corona must have an influence on the wind speed. Figure 7 shows the spatial distribution of the radial wind speed at height of 10 $R_\odot$ for CR 1837 (maximum of activity) and CR 1902 (minimum of activity). Due to the absence of Alfvén wave induced mechanisms, the maximum wind speed at 10 $R_\odot$ is around 250 km s$^{-1}$, while the slowest winds reach two thirds of this value. At the edge of the simulation domain, the fast wind goes beyond 400 km s$^{-1}$ and is close to the terminal speed given by a 1D polytropic model. Near activity maximum, slow and fast wind streams are located at any latitude, and we can identify the so-called ballerina skirt that forms a S shape along all longitudes. For activity minimum, the global axisymmetry of the magnetic field structure is found in the speed distribution, with faster streams at high latitude, and slower streams near the equator.

Following the observations of Wang & Sheeley (1990), the terminal wind speed of the solar wind is generally thought to be anti-correlated to the expansion factor, which computes the deviation from a purely radial magnetic field and which can be defined as

$$f_{\exp}(s) = \frac{B_0}{B(s)} \left( \frac{r_0}{s} \right)^2,$$

where $r_0$ is the radius where the field lines integration starts (here just above the surface), and $s$ is the curvilinear abscissa along the field line. More specifically, this expression corresponds to the total expansion factor and will measure the accumulated expansion along the field line. As soon as the field becomes purely radial, the total expansion factor becomes constant.

In Figure 8, we show the structure of 20 field lines for four Carrington rotations and the total expansion factor for open field lines among them. In all of the cases shown in Figure 8, the expansion profile is still evolving at 10 solar radii, i.e., well beyond $r_{\text{ss, opt}}$, which is shown with the black vertical line and defined non-locally thanks to a procedure based on the open flux of the simulation (see Section 2). As shown in Table 1, the expansion is also acting beyond the sonic point (defined for each field lines and superimposed with the same color code on the expansion profile in Figure 8) and the Alfvén point.

The same process occurs in all coronal holes even when the field geometry is more complex, near maximum of activity. However, in this configuration, the expansion around small closed loops is much more important, making the expansion factor at the source surface actually the most relevant parameter...
to determine the wind speed. Figure 9 deepens this analysis. We show the wind speed at \( r = R_\odot \) as a function of the total expansion factor computed for the 40 open field lines (among which are the one of Figure 8) at three different heights: the optimal source surface, the sonic point, and at 10 solar radii. The scattered points are colored according to the parameter \( f_{\text{dip}} \), which is a good proxy for the cycle phase (see Figure 1). Considering all points as a whole, we observe a global anti-correlation between the expansion factor and the wind speed, which is in agreement with observations (see Wang & Sheeley 1990; Wang 2016). At the source surface, the global anti-correlation is shown with a brownish line, and no correlation can be extracted for the point near minimum alone.

Nevertheless, higher in the corona, we can obtain two different regressions, for simulations whose \( f_{\text{dip}} \) parameter is over or below 50% of the maximum value \( (f_{\text{dip,max}} = 0.35) \) for the period considered. For near minimum cases, we get a positive correlation (shown by a black line) between the wind speed and the expansion factor starting at the sonic point. This correlation gets better as we go up in height at \( R_\odot \). Hence, for a simple, near-dipolar configuration of the solar magnetic field, where most of the open field lines come from the large
polar coronal holes, the expansion around and beyond the sonic point operates a significant reorganization. Of course, beyond the source surface, the total area is conserved and both super and under radial expansion must compensate. As such, super (under) radial expansion increases (decreases) the wind speed, and the correlation between the terminal speed and the total expansion factor is reversed (see also Pinto & Rouillard 2017). The additional acceleration provided only by the super radial expansion is maximum on field lines that cross local maxima of $B_r(\theta, r_{ss,opt})$. This process must influence the acceleration of the fast solar wind, which is thought to occur at large distances from the Sun, where Alfvén waves dissipate. The correlation (shown by the orange line) stays globally negative for points near maximum of activity everywhere in the corona.

In Figure 10, we show the mass flux as a function of the expansion factor at the source surface $f_{ss}$ at different heights in the corona. As shown by the blue dots, a large expansion below the source surface (associated with slow streams) increases the mass flux at the coronal base, through the velocity that is left free in the boundary. The mass flux at the surface spans over an order of magnitude, and the mass loading diminishes the energy per particle and the terminal speed. This explains the good correlation between the wind speed and the expansion factor at the source surface. However, near the source surface, the mass flux/expression correlation is reversed with faster wind (low $f_{ss}$) having a larger mass flux. The standard deviation from the mean is 22% at $r = r_{ss}$. The latter correlation is not observed in the solar wind where the mass flux is almost constant (although slightly larger in the slow solar wind, see Holzer et al. 2005). Expansion above the source surface acts to correct this relationship, decreasing the mass flux and its energy per particle. This process must influence the acceleration of the fast solar wind, which is thought to occur at large distances from the Sun, where Alfvén waves dissipate. The correlation (shown by the orange line) stays globally negative for points near maximum of activity everywhere in the corona.

As the energy input from the chromosphere. Indeed, rapid expansion below $2.5 R_\odot$ can lead to an enhanced local deposition of heat in the low corona, which helps lift material from the chromosphere as shown by Wang et al. (2009). Although the heating mechanisms used in this paper are not related to the magnetic field at the coronal base, we also get higher mass flux for higher surface field strength. This can be understood by the correlation between high expansion in the low corona and high magnetic flux concentration at the surface associated with small-scale structures. Our model thus agrees with more elaborated thermodynamical model of the corona. Also, the fact that we keep the coronal temperature fixed at the base of the corona helps the simulations comply with this observation (because the mass flux depends exponentially on the coronal temperature, see e.g., Hansteen & Velli 2012). Reaching a quantitative agreement with solar wind observables needs both an accurate description of the heating in the low corona and of the continuous expansion up to tens of solar radii (see e.g., Pinto & Rouillard 2017, where the expansion in higher corona has been added to a multi flux tube model based on a PFSS).

6. Discussions

In this study, we showed that the varying structure of the surface solar magnetic field over the 11-year cycle has a strong influence on the higher corona, and the organization of the interplanetary magnetic field. Using fixed thermodynamics, we showed that the characteristic surfaces of the solar corona varies over the cycle by 20%–30% at most, which is less than what others works have suggested (see Pinto et al. 2011). The optimal source surface radii, that contains all coronal loops, are as expected between two and three solar radii. The sonic point is located between 5 and 6 $R_\odot$, while the Alfvén point is located slightly further on. Taking into account both the organization of the field (the topology) and its energy (or surface flux/strength) is crucial to obtain these results. Both seem to balance over the cycle to limit variations of the Alfvén surface. Although more realistic heating prescriptions are likely to have a strong influence on the location of the critical surfaces, and their variation over the cycle, a similar result has been obtained with 3D MHD simulations for two epochs (minimum and maximum) with an Alfvén wave turbulence heating in the work of Garraffo et al. (2016). This balance may be a reliable output of the dynamo mechanisms at the source of the solar magnetic field.

The constant radial field in latitude observed by Ulysses, is the result of an equilibrium between magnetic forces and fluid dynamics, which is occurring at different heights over the cycle. The PFSS model, which imposes an oversimplified structure of the corona beyond $2.5 R_\odot$, cannot capture this process. Moreover, because the magnetic pressure decays strongly close to the star near maximum of activity, this equilibrium is met closer than during minimum of activity. In terms of field expansion, this translates into an extended expansion in the higher corona, more significant near minimum. As a consequence, depending on the heights considered for the total expansion factor, the net effect on the wind velocity can be reversed between the low corona (at the source surface) and the higher corona (here around 10 $R_\odot$) for certain period of the cycle, characterized by a strong dipolar component.
This reorganization influences the solar wind properties by increasing the speed and decreasing the mass flux in the core of coronal holes, where $B_r(\theta)$ is maximum at the source surface. It is also necessary to get a mass flux relatively independent of the wind speed at large distances. In the theoretical model of Wang & Sheeley (1991), used to explain the observations of Wang & Sheeley (1990), a nonlinear relationship between the expansion factor at the source surface $f_s$ or the sonic point $f_c$ and the total expansion at Earth orbit $f_E$ is introduced. The relationship reads $f_i \approx f_c = 3 f_s^{1/2}$, which means that whenever $f_s$ is less than nine, the total expansion will increase between the sonic point and the observer, thus accelerating the wind and reducing the mass flux. For $f_s > 9$, the opposite will occur, yielding a rather constant mass flux for all wind speed. The transition value used in the relation of Wang & Sheeley (1991) matches the expansion limit we observe at minimum of activity, and confirms that the expansion parametrized in Wang & Sheeley (1991) is due to the pressure equilibrium met in the higher corona at solar minimum.

Hence, despite simple driving mechanisms, our model renders a general process that should be taken into account in every more sophisticated solar wind model. Moreover, our study might have consequences on the interpretation of the WSA semi-empirical relationship (see Arge & Pizzo 2000 and Arge et al. 2003). In this formulation, another parameter, $\theta_0$, is used to control the terminal wind speed. It represents the minimum angular distance to a coronal hole boundary and is thought to be strongly correlated to $f_s$ (McGregor et al. 2011), so that fast wind streams emerge from deep within coronal holes while slow wind originates near the streamers. We have shown that within large coronal holes (large $\theta_0$), the expansion beyond the source surface has a strong influence on the wind speed. It is likely that part of the efficiency of this parameter is due to what happens in the upper corona rather than what happens near the photosphere. A proper description of the expansion below and beyond the sonic point over an extended and realistic profile of a solar wind flux tube will certainly provide more accurate and physically motivated predictive laws of the wind terminal speed.

As demonstrated by many works (see e.g., Leer et al. 1982; Withbroe 1988), an extended energy deposition is needed to reproduce the fast solar wind. A more detailed thermodynamics is key to achieve a fully consistent picture, and recent works have introduced advanced Alfvén wave related pressure and heating in 3D MHD simulations (Lionello et al. 2009; Sokolov et al. 2013; van der Holst et al. 2014). However, we believe that additional heating and momentum deposition will not qualitatively change the results we get in this paper. As stated earlier, the model of Wang & Sheeley (1991)—which includes these processes—needs the high corona reorganization described in this work. More realistic models will likely get the sonic point closer to the Sun, near the optimal source surface radius in open regions (fast solar wind). Because the PFSS model is able to relativistically precisely recover the structure of the corona up to $2.5 R_\odot$, the observed homogenization of the radial magnetic field must occur beyond. Thus, beyond the sonic point where Alfvén waves are thought to be dissipated, the influence of expansion on the solar wind speed may be even more important than presented here.

Our work could also contribute to a more accurate description of the expansion in local models that focus on a given flux tube. These models are now the only ones able to capture a fully self-consistent description of the propagation and the dissipation of Alfvén waves (Suzuki & Inutsuka 2006; Cranmer et al. 2007; Matsumoto & Suzuki 2012). Yet, they generally use fast and very localized expansion prescriptions of what happens below the source surface (see the general formulation of Kopp & Holzer 1976). Improving these kinds of models and comparing them with observations will better help understand the acceleration profile of the solar wind in open regions. The upcoming spacecraft Parker Solar Probe (Fox et al. 2016) will continuously measure the evolution of the wind speed and the mass flux down to $9 R_\odot$. Solar Orbiter (Müller et al. 2013), traveling out of the ecliptic, will be the successor of Ulysses while approaching the Sun much closer. Both satellites will probe the nature and the location of the different mechanisms addressed in part in this paper, which are responsible for the solar wind observables at 1 au.

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