Dynamo Action of Jupiter’s Zonal Winds

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Abstract

The new data delivered by NASA’s Juno spacecraft significantly increase our understanding of Jupiter’s internal dynamics. The gravity data constrain the depth of the zonal flows observed at cloud level and suggest that they slow down considerably at a depth of about 0.96 rJ, where rJ is the mean radius at the one bar level. Juno’s magnetometer reveals the planet’s internal magnetic field. We combine the new zonal flow and magnetic field models with an updated electrical conductivity profile to assess the zonal wind induced dynamo action, concentrating on the outer part of Jupiter’s molecular hydrogen region where the conductivity increases very rapidly with depth. Dynamo action remains quasi-stationary and can thus reasonably be estimated where the magnetic Reynolds number remains smaller than one, which is roughly the region above 0.96 rJ. We calculate that the locally induced radial magnetic field reaches rms values of about 10^{-6} T in this region and may just be detectable by the Juno mission. Very localized dynamo action and a distinct pattern that reflects the zonal wind system increases the chance to disentangle this locally induced field from the background field. The estimates of the locally induced currents also allow calculating the zonal flow related Ohmic heating and associated entropy production. The respective quantities remain below new revised predictions for the total dissipative heating and total entropy production in Jupiter for any of the explored model combinations. Thus neither Ohmic heating nor entropy production offer additional constraints on the depth of the zonal winds.

1 Introduction

Two of the main objectives of NASA’s Juno mission are to measure Jupiter’s magnetic field with unprecedented resolution and to determine the depth of the fierce zonal winds observed in the planet’s cloud layer. The first Juno-based internal magnetic field model JRM09 (Connerney et al. 2018) already provides the internal magnetic field up to spherical harmonic degree 10 and shows several interesting features that seem unique to Jupiter’s dynamo (Moore et al. 2018). Better resolved models are expected as the mission continues.

Based on Juno gravity measurements (Iess et al. 2018), Kaspi et al. (2018) deduce that the speed of the equatorially antisymmetric zonal flow contributions must be significantly reduced at a depth of about 3000 km below the one bar level, which corresponds to a radius if 0.96 rJ. Kong et al. (2018) come to roughly similar conclusions with a different inversion procedure, but they also point out that the solution is not unique. While the gravity data only allow constraining the equatorially antisymmetric winds, the results likely also extend to the symmetric contributions. New interior models (Guillot et al. 2018)
Debras & Chabrier (2019) and also the width of the dominant equatorial jet (Gastine et al., 2014; Heimpel et al., 2016) both support the idea that the fast zonal winds are roughly confined to the outer 4% in radius.

The fast planetary rotation enforces geostrophic flow structures with minimal variation along the direction of the rotation axis. Geostrophic zonal winds are thus expected to reach right through the planet’s gaseous envelope, and it remains unclear which mechanism limits their extend in Jupiter. The demixing of hydrogen and helium and the subsequent precipitation of helium deeper into the planet offers one possible explanation (Militzer et al., 2016). This process would have established a helium gradient that suppresses convection. In Jupiter, this stable helium-rain layer may start somewhere between 0.93 and 0.90 $r_J$ and perhaps extends down to 0.80 $r_J$ (Debras & Chabrier, 2019). Note, however, that ab initio simulations by Schottler & Redmer (2018) predict that the hydrogen/helium demixing may not even have started. Recent analysis of gravity measurements by the Cassini spacecraft suggest that Saturn’s zonal winds may only reach down to about 0.85 $r_S$ (Iess et al., 2019; Galanti et al., 2019). Since the stably stratified layer is thought to start significantly deeper, at about 0.62 $r_S$ according to Schottler & Redmer (2018), it cannot be the reason for this limited depth extent of Saturn’s zonal winds.

A second possibility to brake the zonal winds at depth are Lorentz forces. Lorentz forces are tied to dynamo action and thus to the electrical conductivity profile. Ab initio simulations for Jupiter suggest that ionization effects lead to a super-exponentially increase of the electrical conductivity in the outermost molecular gas envelope. We will refer to this layer as Jupiter’s Steeply Decaying Conductivity Region (SDCR) in the following. At about 0.9 $r_J$, hydrogen, the planet’s main constituent, becomes metallic, and the conductivity increases much more smoothly with depth (French et al., 2012) (see panel a) of fig. 1). Though dynamo action and the potential braking of the zonal winds due to Lorentz forces are classically attributed to the metallic region, they may already become significant where the electrical conductivity reaches sizable levels in the SDCR.

Different dynamo-related arguments have been evoked to estimate the depth of the zonal winds without, however, directly addressing the role of the Lorentz forces. Liu et al. (2008) estimate that the Ohmic heating caused by zonal-wind related induction would exceed the total heat emitted from Jupiter’s interior, should the winds reach deeper than 0.96 $r_J$ with undiminished speed. Ridley & Holme (2016) argue that the secular variation of the magnetic field over 30 years of pre-Juno observations is rather small and thus likely incompatible with an advection by undiminished zonal winds. They conclude that the winds cannot reach to depths where the magnetic Reynolds number exceeds one and more significant induction can be expected. This puts the maximum depth somewhere between 0.96 $r_J$ and 0.97 $r_J$, as we will discuss below. A recent analysis by Moore et al. (2019) suggests that the observations over a 45 year time span including Juno data would by compatible with zonal wind velocities of 2.4 m/s at 0.95 $r_J$, two orders of magnitude smaller than observed in the cloud layer.

Another interesting question is how much the dynamo action in the SDCR contributes to Jupiter’s total magnetic field. Using a simplified mean-field approach, Cao & Stevenson (2017) predict that the radial component of the Locally Induced Field (LIF) may reach 1% of the background field and could thus be detectable by the Juno magnetometer. Wicht et al. (2019) analyze the dynamo action in the SDCR of fully self-consistent numerical simulations that yield Jupiter-like magnetic fields. Because of the dominance of Ohmic diffusion, the dynamo dynamics becomes quasi-stationary in the SDCR of their simulations. A consequence is that the locally induced electric currents and field can be estimated with decent precision when flow, electrical conductivity profile, and the surface magnetic field are known. Refined information on all three ingredients has recently become available for Jupiter, allowing for a fresh look on the problem.

Here we use three different zonal flow models, two electrical conductivity models, and the new Juno-based magnetic field model JRM09 to predict the electric currents and magnetic fields produced in Jupiter’s SDCR. In addition, we also derive new estimates for the total dissipative heating and related entropy production and explore whether either value is exceeded by the zonal-flow related Ohmic dissipation.

The article starts off with outlining the methods and introducing the used data in sect. 2. Sect. 3 discusses dissipative heating and entropy production in Jupiter. Estimates for dynamo action, Ohmic heating, and entropy production are then presented in sect. 4. Sect. 5 closes
the article with a discussion and conclusion.

2 Methods and Data

2.1 Estimating Dynamo Action

The ratio of inductive to diffusive effects in the induction equation,

$$\frac{\partial B}{\partial t} = \nabla \times \left( U \times B \right) - \nabla \times \lambda \nabla \times B ,$$

(1)

can be quantified by the magnetic Reynolds number

$$\text{Rm} = \frac{\langle U \rangle D}{\lambda} ,$$

(2)

where $\lambda = 1/(\mu \sigma r)$ is the magnetic diffusivity, with $\mu$ the magnetic permeability and $\sigma$ the electrical conductivity. Angular brackets generally denote rms values at a given radius throughout the paper; thus $\langle U \rangle$ stands for

$$\langle U \rangle = \left( \frac{1}{4\pi} \int_{0}^{2\pi} d\phi \int_{0}^{\infty} d\theta \sin \theta \ U \right)^{1/2} ,$$

(3)

$\theta$ being the colatitude and $\phi$ the longitude.

The typical length scale $D$ is hard to estimate, and the planetary radius is often used for simplicity. Where $\sigma$ decreases steeply in the SDCR, however, the length scale is determined by the conductivity or magnetic diffusivity scale height

$$D_{\lambda} = \frac{\lambda}{\partial \lambda / \partial r} = -\frac{\sigma}{\partial R / \partial r} ,$$

(4)

and the modified magnetic Reynolds number

$$\text{Rm}^{(1)} = \frac{\langle U \rangle D_{\lambda}}{\lambda}$$

(5)

should be used. Since $D_{\lambda}$ is small and $\lambda$ decreases steeply with radius, most of the SDCR is characterized by a small magnetic Reynolds number $\text{Rm}^{(1)} < 1$, and the magnetic field dynamics becomes quasi-stationary (Liu et al. 2008), obeying the simplified induction equation

$$\nabla \times \frac{j}{\sigma} \approx \nabla \times \left( U \times \hat{B} \right) .$$

(6)

Here, $j$ is the current density and $\hat{B}$ the strong background field produced by the dynamo acting deeper in the planet. The locally induced field $\hat{B}$ is given by Ampère’s law:

$$j = \nabla \times \hat{B} / \mu .$$

(7)

The steep $\sigma$ profile dominates the radial dependence of $j$ and $\hat{B}$ in the SDCR. The current density is thus dominated by the horizontal components, where radial gradients in $B$ contribute (Liu et al. 2008, Wicht et al. 2018):

$$j \approx j_{H} \approx \hat{r} \times \frac{\partial}{\partial r} \hat{B} .$$

(8)

Index $H$ denotes the horizontal components; the radial current can be neglected in comparison.

Along the same lines, the horizontal components of eqn. (6) can be approximated by

$$\frac{1}{r} \frac{\partial}{\partial r} \int_{r_{i}}^{r_{f}} \frac{\partial}{\partial \sigma} \left[ \nabla \times \left( U \times \hat{B} \right) \right] d r' ,$$

(9)

where $\hat{r}$ is the radial unit vector. Integration in radius yields the integral current density estimate introduced by Liu et al. (2008), which we identify with an upper index (I):

$$j_{H}^{(I)} = \frac{\sigma}{r} \left[ \frac{r}{\sigma} j_{H} \right] r_{i} + \hat{r} \times \frac{\sigma}{r} \int_{r_{i}}^{r_{f}} d r' \left[ \nabla \times \left( U \times \hat{B} \right) \right]_{H} ,$$

(10)

The square brackets with a lower index $r_{j}$ indicate that the expression should be evaluated at the outer boundary.

For a predominantly zonal flow, we can use the approximation

$$\nabla \times \left( U \times \hat{B} \right) \approx -\frac{U_{\phi}}{r \sin \theta} \left( \frac{\partial}{\partial \phi} \hat{B}_{\phi} \right) \hat{\theta} + \frac{1}{r} \left( \frac{\partial}{\partial r} \left( r \hat{U}_{\phi} \hat{B}_{r} \right) + \frac{\partial}{\partial \theta} \left( \hat{U}_{\phi} \hat{B}_{\theta} \right) \right) \hat{\phi} ,$$

(11)

where $U_{\phi}$ is the zonal flow component and $\hat{\theta}$ and $\hat{\phi}$ are unit vectors in latitudinal and azimuthal direction, respectively.

The integral estimates for the two horizontal current components are then given by

$$j_{H}^{(I)} = \frac{\sigma}{r} \left[ \frac{r}{\sigma} j_{0} \right] r_{i} - \frac{\sigma}{r} \left[ r \hat{U}_{\phi} \hat{B}_{r} \right]_{r_{f}} - \frac{\sigma}{r} \int_{r_{i}}^{r_{f}} d r' \frac{\partial}{\partial \theta} \left( \hat{U}_{\phi} \hat{B}_{\theta} \right)$$

(12)
and

\[ j^{(f)}_\phi = \sigma \left[ \frac{r}{\sigma} j_\phi \right]_r - \sigma \int_r^\infty \frac{U_\phi}{\sin \theta} \frac{\partial}{\partial \phi} \tilde{B}_\theta \, dr' . \]  

(13)

Since the latitudinal length scale of the zonal winds is smaller than the azimuthal length scale of the magnetic field, we expect that the latitudinal component dominates.

The integral estimate requires the knowledge of the surface currents. While the surface currents are certainly very small, the scaled version \( \sigma(r)/\sigma(r_j) j \) may remain significant. Liu et al. (2008) argue that neglecting the surface contribution at least provides a lower bound for the rms current density.

Wicht et al. (2019) confirm that the dynamics indeed becomes quasi-stationary in the SDCR of Jupiter-like dynamo simulations where \( Rm^{(1)} < 1 \) and show that \( j_\phi \) is indeed the dominant current component in the SDCR of their Jupiter-like dynamo simulations. They also report that the simplified Ohm’s law for a fast moving conductor,

\[ j^{(O)} = \sigma \left( U \times \hat{B} \right) , \]  

(14)

provides a significantly better estimate than \( j^{(f)} \). We identify the respective current estimate with an upper index \( (O) \). The general Ohm law,

\[ j = \sigma \left( U \times \hat{B} + E \right) , \]  

(15)

also contains currents driven by the electric field, which reduces to \( E = -\nabla \Phi \) in the quasi-stationary case, where \( \Phi \) is the electric potential. In the SDCR, this contribution likely proves secondary because the potential differences remain small compared to the induction by fast zonal winds (Wicht et al. 2019).

As the electrical conductivity decreases in the SDCR, the magnetic field approaches a potential field with its characteristic radial dependence. We use this dependence to approximate the background field with

\[ \tilde{B}_r(r, \theta, \phi) \approx \left( \frac{r_j}{r} \right)^{\ell+2} B_r(r_j, \theta, \phi) , \]  

(16)

where the index \( \ell \) denotes the magnetic field contribution at spherical harmonic degree \( \ell \). This provides a decent approximation as long as the LIF remains a small contribution of the total field (Wicht et al. 2019).

Given a surface field model and an electrical conductivity profile, Ohm’s law for a fast moving conductor and a predominantly zonal flow suggests

\[ j \approx j^{(O)}_\phi = \sigma \frac{U_\phi}{\sin \theta} \tilde{B}_r . \]  

(17)

When using this result to constrain the outer-boundary currents, the alternative integral estimates, eqn. (12) and eqn. (13), yield

\[ j^{(f)}_\mu = \sigma \frac{U_\phi}{\sin \theta} \tilde{B}_r - \frac{\sigma}{r} \int_r^\infty \frac{U_\phi}{\sin \theta} \frac{\partial}{\partial \theta} \left( U_\phi \tilde{B}_\theta \right) \]  

(18)

and

\[ j^{(f)}_\phi = -\frac{\sigma}{r} \int_r^\infty \frac{U_\phi}{\sin \theta} \frac{\partial}{\partial \theta} \tilde{B}_\theta , \]  

(19)

respectively. A comparison of the estimates shows that \( j^{(f)}_\mu \) and \( j^{(O)}_\phi \) will remain very similar at shallow depths. When the flow decays very deeply with depth, however, the integral contributions in eqn. (18) and eqn. (19) will dominate below some radius and cause larger deviations, as we will see below.

Calculating the LIF requires to uncurl Ampere’s law, which reduces to integrating eqn. (8) in the SDCR. When using \( j^{(O)}_\phi \), this yields

\[ \tilde{B}_H \approx \int_r^\infty \frac{\hat{r} \times \left( U \times \tilde{B} \right)}{\lambda} \, dr' . \]  

(20)

Since the electrical conductivity profile rules the radial dependence, the integral can be approximated by

\[ \tilde{B}_H \approx \frac{D_\lambda}{\lambda} \hat{r} \times \left( U \times \tilde{B} \right) . \]  

(21)

We have assumed here that the LIF vanishes at the outer boundary. For a dominantly azimuthal flow, the primary LIF component is also azimuthal:

\[ \tilde{B}_\phi \approx \frac{D_\lambda}{\lambda} \frac{U_\phi \tilde{B}_r}{\hat{r}} . \]  

(22)

This suggests that the rms value scales with \( Rm^{(1)} \),

\[ \langle \tilde{B}_H \rangle \approx Rm^{(1)} \langle \tilde{B} \rangle , \]  

(23)

assuming that the correlation between \( U_\phi \) and \( \tilde{B}_r \) is of little relevance.
The radial LIF can be estimated based on the radial component on the quasi-stationary induction equation (1):

\[
\lambda \nabla^2 \hat{B}_r \approx -\nabla \times (U \times \hat{B}) .
\] (24)

When approximating Ohmic dissipation by \( \lambda \hat{B}_r / D_\lambda^2 \), this yields:

\[
\hat{B}_r \approx -D_\lambda^2 \lambda \cdot \nabla \times (U \times \hat{B}) ,
\] (25)

which reduces to

\[
\hat{B}_r \approx -D_\lambda^2 \lambda \frac{U}{r \sin \theta} \frac{\partial}{\partial \phi} \hat{B}_r
\] (26)

for a predominantly zonal flow. This suggest that the rms radial LIF should roughly scale with the second modified magnetic Reynolds number

\[
\text{Rm}^{(2)} = \frac{\langle U \rangle D_\lambda^2}{\lambda D}
\] (27)

like

\[
\langle \hat{B}_r \rangle \approx \text{Rm}^{(2)} \frac{D}{D_\phi} \langle B \rangle.
\] (28)

Here \( D_\phi \) is the azimuthal length scale of the background field. Since \( D_\lambda \ll D_\phi \), the radial LIF is much smaller than its horizontal counterpart (Wicht et al. 2019).

### 2.2 Data

The electric current and LIF estimates discussed above require a conductivity profile, a zonal flow model, and a surface magnetic field model. For the heating and entropy estimates that we will derived in sect. 3, we also need density, temperature, and thermal conductivity profiles. We adopt the interior model calculated by Nettelmann et al. (2012) and French et al. (2012), which is the only one proving all the required information. Note, however, that recent Juno gravity data suggest that Jupiter’s interior may be more complex than anticipated in this model (Debras & Chabrier 2019).

Ab initio simulations of the electrical conductivity by French et al. (2012) provide 12 data points at different depths. Fig. 1 shows the values in the outer 20% of Jupiter’s radius and the parametrization \( \sigma_F(r) \) developed for our analysis. A linear branch,

\[
\sigma_F(r) = \sigma_r + (\sigma_m - \sigma_r) \frac{r - r_r}{r_m - r_r},
\] (29)

covers the smoother inner part \( r < r_m \). An exponential branch,

\[
\sigma_F(r) = \sigma_m \exp \left[ \frac{r - r_m}{r_m - r_r} + b \left( \frac{r - r_m}{r_m - r_r} \right)^2 \frac{\sigma_m - \sigma_r}{\sigma_m} \right],
\] (30)

describes the steeper decay for \( r_m < r < r_e \) with \( b = 7.2 \). Matching radius \( r_m = 0.89 r_J \) and reference radius \( r_r = \)

![Figure 1: (a) Electrical conductivity profiles in the outer 20% of Jupiter’s radius. The black line shows the parametrization \( \sigma_F(r) \) of the ab initio simulation data points (black circles) by French et al. (2012). The dotted red line shows the profile published in Zaghoo & Collins (2018), while the solid red line shows the extension \( \sigma_Z(r) \) used here. The profiles suggested by Liu et al. (2008) (green) and Nellis et al. (1999) (blue) are shown for comparison. (b) Magnetic diffusivity scale height \( D_\lambda/r_J \) for \( \sigma_F \) and \( \sigma_Z \).](image)
0.77 \ r_J \ are \ chosen \ where \ ab \ initio \ data \ points \ have \ been \ provided.

A double-exponential branch,

$$\sigma_F(r) = \sigma_e \exp\left( d \left( \exp\left( c \frac{r - r_e}{r_m - r_e} \right) - 1 \right) \right), \quad (31)$$

is required to capture the super-exponential decrease for \( r \geq r_e = 0.972 \ r_J \). The additional free parameter is \( c = 10 \), while \( \sigma_e = \sigma(r_e) \) and

$$d = \frac{1}{c} \left( 1 + 2 b \frac{r_e - r_m}{r_m - r_i} \right) \frac{\sigma_m - \sigma_r}{\sigma_m}. \quad (32)$$

The dotted red line in fig. 1 shows the conductivity model used to study dynamo action in Jupiter and Jupiter-like exoplanets by Zaghoo & Collins (2018). This is based on measurements which suggest a higher electrical conductivity in the metallic hydrogen phase than previous data. Unfortunately, Zaghoo & Collins (2018) do not discuss how the results were extrapolated to Jovian conditions. The solid red line in fig. 1 shows the respective parametrization \( \sigma_Z(r) \) used for our analysis, which retraces the published curve and connects to previously published parametrizations (green and blue) at lower densities (Nellis et al. 1996; Liu et al. 2008). Note, however, that these parametrizations are based on data (Weir et al. 1996) which may have been attributed to too low temperatures according to a recent analysis by Knudson et al. (2018).

Though model \( \sigma_Z(r) \) is somewhat arbitrary, it serves to illustrate the impact of conductivity uncertainties in our study. Close to \( r_J \) where conductivities remain insignificant, \( \sigma_F \) is many orders of magnitude larger than \( \sigma_Z \). The ratio \( \sigma_F/\sigma_Z \) decreases with depth, reaching \( 10^2 \) around 0.97 \( r_J \) and 10 around 0.96 \( r_J \). The two models finally cross at about 0.95 \( r_J \). At about 0.925 \( r_J \), the ratio reaches a minimum of 0.05 and then slowly increases with depth to 0.35 at 0.8 \( r_J \). Tab. 1 list values of both conductivity models for selected radii.

Panel b) of fig. 1 and selected values listed in tab. 1 demonstrate that the magnetic diffusivity scale heights \( D_\lambda \) differ much less than the conductivities themselves. Electric currents, locally induced fields, and Ohmic heating depend linearly on \( \sigma \) but on different powers of \( D_\lambda \). The differences between the results for the two conductivity models is thus predominantly determined by \( \sigma \) and can easily be scaled from one to the other.

Figure 2: Zonal flow models used in this study. Prograde flows are shown in red and yellow, while blue indicates retrograde directions.
Table 1: Rms flow velocities, electrical conductivities $\sigma$, magnetic diffusivities $\lambda$, diffusivity scale heights $D_1$, and magnetic Reynolds numbers $\text{Rm}^{(1)}$ at selected radii.

| $r/r_J$ | 0.98 | 0.97 | 0.96 | 0.95 | 0.94 | 0.92 | 0.90 |
|---------|------|------|------|------|------|------|------|
| $\langle U_G \rangle$ [m/s] | $3.0 \times 10^4$ | $3.2 \times 10^4$ | $3.5 \times 10^4$ | $3.8 \times 10^4$ | $4.1 \times 10^4$ | $4.5 \times 10^4$ | $4.9 \times 10^4$ |
| $\langle U_K \rangle$ [m/s] | $2.0 \times 10^4$ | $1.4 \times 10^4$ | $8.0 \times 10^3$ | $4.0 \times 10^3$ | $2.0 \times 10^3$ | $2.5 \times 10^3$ | $3.0 \times 10^3$ |
| $\langle U_Z \rangle$ [m/s] | $2.3 \times 10^4$ | $1.9 \times 10^4$ | $1.5 \times 10^4$ | $1.1 \times 10^4$ | $6.8 \times 10^3$ | $1.7 \times 10^3$ | $2.4 \times 10^3$ |

The three different zonal flow models explored here are illustrated in fig. 2. Tab. 1 lists rms values $\langle U_{\theta} \rangle$ at selected radii. All reproduce the observed zonal winds at $r = r_J$ (Porco et al. 2003; Vasavada & Showman 2005). We use running averages of the surface profiles with a window width of one degree and represent the result with 256 (nearly) evenly spaced latitudinal grid points for our calculations.

The three flow models differ at depth. The most simple one, $U_G$, assumes geostrophy in each hemisphere, i.e. the flow depends only on the distance $s = r \sin \theta$ to the rotation axis. Kaspi et al. (2018) describe the depth decay of the equatorially antisymmetric zonal flow with profiles constrained by the Juno gravity measurements. We apply their ‘latitude independent’ model version to the total zonal flow and refer to this model as $U_K$. The rms amplitude of $U_K$ has decreased by one order of magnitude at about 0.95 $r_J$ and by two orders of magnitude around 0.925 $r_J$.

We also consider the ‘deep’ model suggested by Kong et al. (2018), who assume an exponential depth decay and an additional linear dependence on the distance $z = \cos \theta$ to the equatorial plane. Like for $U_K$, our respective model $U_Z$ assumes that the depth and $z$ dependencies, which were originally derived for the equatorially antisymmetric contributions, apply to the whole flow. The rms velocity in $U_Z$ decays smoother with depth than in $U_K$, having decreased by one order of magnitude at about 0.935 $r_J$ and by two orders of magnitude at about 0.905 $r_J$.

Fig. 2 shows that $U_G$ and $U_K$ have discontinuities at the equatorial plane. These pose a problem when calculating the latitudinal zonal flow derivatives required for the integral estimate $f_I^{(1)}$ (see eqn. (18)). Formally, the derivative becomes infinite at the equator. Practically, however, the impact of the discontinuity depends on the model setup and on the methods used for calculating the derivatives. We tested the impact on rms current density estimates by comparing calculations covering all latitudes with counterparts where the derivatives were explicitly set to zero in a six-degree belt around the equator. Simple first order finite differences with 256 grid points at each radial level are generally used for calculating the derivative. For flow $U_Z$, which has been constructed to avoid the discontinuity (Kong et al. 2018), the belt contributes not more than one percent to $\langle f_I^{(1)} \rangle$ at any radius, which is less than the surface fraction it represents. For flow $U_K$, the contribution is even smaller due to the faster decay of the flow amplitude. However, for $U_G$ the belt contributes 20% to the rms current for radii below 0.94 $r_J$, which is a clear sign that the unphysical discontinuity causes problems. In order to
be on the safe side, we will only consider flow model $U_Z$ in connection with estimate $j_0^{(1)}$ below.

![Figure 3: Modified magnetic Reynolds numbers for the three zonal flow and the two conductivity models explored here. Vertical lines mark where the radii $r_1$ where $Rm^{(1)} = 1$. Green lines show profiles for a typical convective velocity of 10 cm/s suggested by scaling laws.](image)

The radius where $Rm^{(1)} = 1$, which we will refer to as $r_1$ in the following, roughly marks the point where the approximations discussed above break down (Wicht et al. 2019). Fig. 3h illustrates the $Rm^{(1)}$ profiles that result from combining $\sigma_F$ and $\sigma_Z$ with rms values for the three zonal flow models. Tab. 1 compares values at some selected radii. These modified magnetic Reynolds numbers exceed unity between $r_1 = 0.957 r_J$ for the combination $\sigma_Z$ and $U_K$ and $r_1 = 0.967 r_J$ for $\sigma_F$ and $U_G$. All $r_1$ values are listed in tab. 2.

Green lines in fig. 3h show $Rm^{(1)}$ profiles for a typical convective velocity of 10 cm/s suggested by scaling laws (e.g. see Duarte et al. 2018). Numerical simulations show that the velocity increases with radius, an effect not taken into account here. The comparison of the different $Rm^{(1)}$ profiles suggests that zonal flow related dynamo action should dominate at least in the outer 9% in radius.

For Jupiter’s surface magnetic field we use the JRM09 model by Connerney et al. (2018), which provides information up to spherical harmonic degree $\ell = 10$. The more recent model by Moore et al. (2018) is only slightly different. In order to check the impact of smaller scale contributions, we also tested the numerical model G14 by Gastine et al. (2014), which reproduces Jupiter’s large scale field and provides harmonics up $\ell = 426$. Since it turned out that the impact of the smaller scales is very marginal, the results are not shown here.

3 Dissipative Heating and Entropy Production in Jupiter

Liu et al. (2008) constrain the depth of the zonal winds by assuming that the related total Ohmic heating should not exceed the heat flux out of the planet. Unfortunately, this assumption is not correct, as we will show in the following. In order to arrive at more meaningful constraints, we start with reviewing some fundamental considerations.

In a quasi-stationary state, where flow and magnetic field are maintained by buoyancy and induction against dissipative losses, the conservation of energy simply states that the heat flux $Q_o = Q(r_o)$ through the outer boundary is the sum of the flux $Q_i = Q(r_i)$ through the inner boundary and the total internal heating $H$:

$$Q_o = Q_i + H.$$  \hspace{1cm} (33)

Note that neither viscous nor Ohmic heating contribute to $H$. Since flow and magnetic field are maintained by the heat flux through the system, they cannot be counted as net heat sources (Hewitt et al. 1975; Braginsky & Roberts 1995).

When furthermore also neglecting the effects of helium segregation, core erosion, or planetary shrinking as potential energy sources, the only remaining contribution is the slow secular cooling of the planet. The volumetric heat source is then given by

$$h = \tilde{\rho} \tilde{T} \frac{\partial \tilde{S}}{\partial t},$$  \hspace{1cm} (34)

where the tilde indicates the hydrostatic, adiabatic, background state (Braginsky & Roberts 1995). Assuming that convection maintains an adiabat at all times, $\partial \tilde{S} / \partial t$ remains homogeneous throughout the convective region and obeys (Jones 2014):

$$\frac{\partial \tilde{S}}{\partial t} = (Q_o - Q_i) / \int_V dV \tilde{T}.$$  \hspace{1cm} (35)
Here, $\int_V dV$ denotes an integration over the whole convective volume. Note, however, that the thermal evolution could be more complex, should Jupiter indeed harbor stably stratified regions.

In order to get a handle on dissipative heating, one has to consider the local heat equation

$$\tilde{\rho} \tilde{T} \left( \frac{\partial s}{\partial t} + U \cdot \nabla s \right) = \nabla \cdot (k \nabla T) + h + \varphi , \quad (36)$$

where $\varphi$ denotes the volumetric dissipative heat source, and $k$ is the thermal conductivity. When assuming a steady state and adopting the anelastic approximation $\nabla \cdot (\tilde{\rho} U) = 0$, the integration over the shell between the inner boundary $r_i$ and radius $r$ yields

$$Q_D(r) + Q_A(r) = Q_i + \int_{r_i}^{r} dr' \int_F dF \left( h + \varphi + \tilde{\rho} s U_r \frac{\partial T}{\partial r} \right). \quad (37)$$

The left hand side is the total flux through the boundary at $r$, i.e. the sum of the diffusive contribution

$$Q_D(r) = -\int_F dF \int_{r_i}^{r} dr' \int_F dF k \frac{\partial T}{\partial r} \quad (38)$$

and the advective contribution

$$Q_A(r) = \int_{F(r)} dF \tilde{\rho} \tilde{T} s U_r . \quad (39)$$

The right hand side of eqn. (37) reflects the influx through the lower boundary $Q_i$ plus three volumetric contributions: the slow secular cooling, the dissipative heating, and the adiabatic cooling. Writing the adiabatic cooling in terms of $Q_A$ yields the relation

$$Q_D(r) + Q_A(r) = Q_i + \int_{r_i}^{r} dr' \int_F dF \left( h + \varphi \right) \quad (40)$$

$$- \int_{r_i}^{r} dr' \int_F dF Q_A(r') / D_T(r') ,$$

where $D_T = -\nabla / (\partial T / \partial r)$ is the thermal scale height.

Integrating eqn. (40) over the whole convective volume and using eqn. (33) reveals that the total dissipative heating $\Phi_T$ is balanced by the total adiabatic cooling:

$$\Phi_T = \int_V dV \varphi = \int_{r_i}^{r} dr' Q_A(r') / D_T(r') . \quad (41)$$

The total adiabatic cooling is actually identical to the buoyancy power $P$ that drives convection and thus the dynamo mechanism. Multiplying the buoyancy term in the Navier-Stokes equation with velocity and integrating over the convective volume to yield the total convective power input indeed gives the same expression (Braginsky & Roberts 1995). Eqn. (41) thus simply states that dissipation is balanced by the power input $P$ to the system, a fact used in many scaling laws to establish how the rms magnetic field strength or rms velocity scale with $P$ (Christensen & Aubert 2006; Christensen et al. 2009; Davidson 2013; Yadav et al. 2013).

Eqn. (41) requires to know $Q_A$ at each radius. Since $Q_A$ itself depends on the distribution of dissipative heat sources, however, an additional condition is required. Assuming that Ohmic heating and adiabatic cooling not only cancel globally but, at least roughly, also at each radius offers a simple solution used in most scaling laws (though never stated explicitly). With the exception of thin thermal boundary layers, the heat flux is then dominated by the advective contribution, so that

$$Q_A(r) \approx Q_i + \int_{r_i}^{r} dr \int_F dF h . \quad (42)$$

Adopting the interior model by Nettelmann et al. (2012) and French et al. (2012) and the observed flux $Q_o = 3.35 \times 10^{17}$ W from the planet’s interior (Guillot & Gautier 2015) allows calculating $h$ via eqn. (34). Because the inner core occupies only 10% in radius, $Q_i$ can be neglected. When, for example, assuming that $h$ also describes the cooling of the rocky core, $Q_i$ is two orders of magnitude smaller than $Q_o$.

Plugging eqn. (42) into eqn. (41) finally allows calculating the total dissipative heating:

$$\Phi_T = 1.20 \times 10^{18} \text{ W} . \quad (43)$$

The result reveals that dissipative heating can in fact exceed the heat flux out of Jupiter’s interior by a factor of 3.6. Gastine et al. (2014) came up with a power estimate that is about 50% smaller because they used a simplified formula provided by Christensen et al. (2009).

Considering the entropy rather than the heat balance avoids the need to come up with an additional condition (Hewitt et al. 1975; Gubbins et al. 1979; Braginsky & Roberts 1995; Gubbins et al. 2003). Dividing the heat
The conditions
\[ \Phi_O(r) \leq \Phi_T = 1.20 \times 10^{18} \text{ W} \] provides a possible constraint for the depth of the zonal winds in Jupiter.

The dissipative entropy production related to the Ohmic heating is given by
\[ \Theta_O(r) \approx \int_r^{r_o} \text{d}r' \int_F \frac{J^2}{\sigma T}. \]
This can be used for the alternative depth constraint
\[ \Theta_O(r) \leq \Theta_T = 2.0 \times 10^{15} \text{ W/K}. \]

4 Dynamo Action in Jupiter’s SDCR

4.1 Electric Currents and Locally Induced Field

We start with discussing the current estimates for the different zonal flow and conductivity model combinations. Fig. 4a) compares rms values of integral estimates \( J^O \) and Ohm’s law estimates \( J^O \) for conductivity model \( \sigma_F \) and flow \( U_Z \). The integral estimates of the latitudinal currents are at least 40 times larger than their azimuthal counterparts for all conductivity and flow model combinations. When using \( J^O \) as outer boundary condition for \( J^O \) (blue line in fig. 4a), both estimates remain very similar down to about 0.98 \( r_j \). At 0.97 \( r_j \), however, \( J^O \) is already about 50% larger than \( J^O \), and at 0.96 \( r_j \) the difference has increased to about 250%. When assuming a vanishing outer boundary current for \( J^O \), the differences are even larger: \( J^O \) (green line) is 3.5 times larger than \( J^O \) at 0.97 \( r_j \) and about 6 times larger at 0.96 \( r_j \).

Estimates of the rms horizontal and radial LIF are shown in fig. 4b), based on \( J^O \) and eqn. (20) for the horizontal and on eqn. (26) for the radial components. The radial LIF is between two and three orders of magnitude smaller than its horizontal counterpart. The rougher estimates, based on Rm(1) and Rm(2) respectively, provide values that are less than a factor two smaller and can thus safely be used for order of magnitude assessments. They correctly predict that the rms azimuthal LIF reaches the level of the background field at \( r_1 \).
and also that the ratio of radial to azimuthal LIF is about $Rm^{(2)}/Rm^{(1)} = D_1/r_J$. At $r_1$, the rms radial LIF is thus roughly three orders of magnitude smaller than the background field or the horizontal LIF. Tab. 2 lists the relative rms radial LIF (column 7) at $r_1$ (column 3) for all $\sigma$ and flow combinations when using $j^{(O)}$.

Wicht et al. (2019) demonstrate that the Ohm’s-law based estimates not only provide good rms but also decent local values for their Jupiter-like dynamo simulations. Fig. 5 shows the radial surface field of JRM09 in panel a) and the radial LIF for $\sigma_F$ and $U_Z$ at $r_1$ in panel b). A very distinct pattern of localized field patches can be found where the fast zonal jets around the equator interact with the strong blue patch in the JRM09 model.

The zonal flow pattern remains recognizable in the LIF, as is clearly demonstrated in fig. 6, which compares zonal flow profiles in panel a) with the azimuthal rms of the radial LIF in panel b). Due to the flow geometry, the currents and LIF show a depth-dependent phase shift relative to the surface jets. The equatorial jet, which is so prominent at the surface, contributes very little to dynamo action, since it does not reach down to depths where the electrical conductivity is more significant.

Fig. 7 compares spherical harmonic power spectra of the background radial field and the radial LIF. As already apparent from the map shown in fig. 5, the LIF is dominated by smaller scale contributions. The spherical harmonic degree spectrum results from the convolution of the complex latitudinal zonal flow structure with the background field. At $r_1$, the dipole contribution in the LIF is about $10^{-4}$ times smaller than the respective background field contribution. For degree $\ell = 10$, the ratio has increased to $10^{-2}$. The spectrum peaks at $\ell = 12$ but has also significant contributions from even higher degrees.

The spherical harmonic order spectrum, shown in panel b) of fig. 7 is very different. The action of the axisymmetric zonal flow on $\tilde{B}_r$ excites no additional harmonic orders so that the spectrum remains confined to $m \leq 10$. The LIF spectrum is rather flat but has no axisymmetric contribution. At $m = 10$, the rms LIF amplitude reaches roughly 25% of the background field.

The results for the conductivity model $\sigma_F$ presented so far can roughly be scaled to model $\sigma_Z$ by multiplying with the conductivity ratio $\sigma_Z/\sigma_F$. Around 0.97$r_J$, the LIF is two orders of magnitude weaker, and the difference decreases with depth, reaching about one order of magni-

Figure 4: (a) Rms current density estimates for flow model $U_Z$ and conductivity model $\sigma_F$. (b) Estimates of rms radial and horizontal LIF. The profile of $\langle \tilde{B}_r \rangle$ has been included for comparison. Vertical lines mark the radii where $Rm^{(1)} = 1$ for the different flow models. Current density estimates $j^{(O)}$ and conductivity model $\sigma_F$ have been used.
are less sensitive to the conductivity models. When using the level of heating remains negligible in outer two percent in radius. 

Table 2: While the first two columns define the model, columns 3 to 6 list different radii: radius $r_1$ where $Rm^{(1)} = 1$, $r_{10}$ where $Rm^{(1)} = 10$, $r_{\phi}$ where $\Phi_O = \Phi_T$, and $r_{\theta}$ where $\Theta_O = \Theta_T$. Column 7 provides the ratio of rms radial LIF to rms radial background field at $r_1$. Columns 8 and 9 give ratios $\Phi_O/\Phi_T$ and $\Theta_O/\Theta_T$ at $r_1$, respectively. Current density estimates $\hat{j}^{(0)}$ have been used.

| conduct. flow | $r_1/r_J$ | $r_{10}/r_J$ | $r_{\phi}/r_J$ | $r_{\theta}/r_J$ | $\langle \hat{B}_r \rangle/\langle \hat{B}_r \rangle$ | $\Phi_O/\Phi_T$ | $\Theta_O/\Theta_T$ |
|--------------|----------|-------------|---------------|----------------|---------------------------------|----------------|----------------|
| $\sigma_F$   | $U_G$    | 0.967       | 0.960         | 0.961          | 2.8x10^{-3}            | 1.2x10^{-1}  | 2.3x10^{-2} |
|              | $U_Z$    | 0.965       | 0.956         | 0.948          | 2.2x10^{-3}            | 8.0x10^{-2}  | 1.5x10^{-2} |
|              | $U_K$    | 0.963       | 0.954         | 0.946          | 3.0x10^{-3}            | 4.8x10^{-2}  | 9.0x10^{-3} |
| $\sigma_Z$   | $U_G$    | 0.960       | 0.956         | 0.956          | 1.8x10^{-3}            | 1.5x10^{-1}  | 2.5x10^{-2} |
|              | $U_Z$    | 0.958       | 0.954         | 0.953          | 1.6x10^{-3}            | 8.9x10^{-2}  | 1.5x10^{-2} |
|              | $U_K$    | 0.957       | 0.952         | 0.948          | 1.7x10^{-3}            | 2.9x10^{-2}  | 4.9x10^{-3} |

4.2 Ohmic Heating and Entropy Constraint

We now use the electric current estimates to calculate Ohmic heating and entropy production. Panel c) of fig. 5 shows the map of Ohmic heat flux density $q = \int_r^r d\tau \hat{j}^2/\sigma$ at radius $r_1$ when using $\hat{j}^{(0)}$, $\sigma_F$, and $U_Z$. The currents induced by interaction between the fierce zonal jets close to the equator and the strong blue patch in JRM09 not only yield a highly localized LIF but also intense local heating. While the action of various other zonal jets reaches a lower level, the related pattern remains roughly recognizable in the form of thin heating bands. The azimuthal mean of $q$, shown in Panel c) of fig. 6, clearly illustrates the correlation between heating and the zonal jets. Like for the LIF, there is a depth-dependent phase shift between the observed surface zonal wind profile and the Ohmic heating pattern.

Panel a) of fig. 8 compares the Ohmic heating profiles $\Phi_O(r)$ for the different zonal flow and electrical conductivity models. Because of the extremely low conductivity, heating remains negligible in outer two percent in radius. When using $\hat{j}^{(0)}$, the outermost radius where $\Phi_O$ reaches the level of $\Phi_T$ is $r_{\phi} = 0.950 r_J$ for flow $U_G$ and both conductivity models. When using $U_K$ and $\sigma_T$, the Ohmic heating always remains below $\Phi_T$. Results based on $\hat{j}^{(1)}$ (not shown) are less sensitive to the differences between the three flow models at depth and are generally similar to the results for $U_G$ and $\hat{j}^{(0)}$.

The different $r_{\phi}$ values where $\Phi_O = \Phi_T$ have been marked by vertical lines in fig. 8 and are listed in column 5 of tab. 2. All are located below the radii $r_1$ where $Rm^{(1)} = 1$ for the respective model combinations (column 3) and thus in a region where the approximations employed here break down. The maximum Ohmic heating reached at $r_1$ remains nearly one order of magnitude below $\Phi_T$ (see column 8 of tab. 2).

Similar inferences hold for the entropy production shown in panel b) of fig. 8. The entropy condition is less strict than the power-based heat condition, and the radii $r_\theta$ where the different models exceed the threshold $\Theta_T$ (column 6 of tab. 2) are somewhat deeper than respective $r_{\phi}$ values. The largest value of $r_\theta = 0.955$ is found for the combination $U_G$ and $\sigma_Z$. The combination of $U_K$ and $\sigma_F$, on the other hand, yields the deepest value of $r_\theta = 0.929$.

The exploration of numerical dynamo simulations by Wicht et al. (2019) suggest that $\hat{j}^{(0)}$ may provide an acceptable estimate for a limited region below $r_1$, at least down to where $Rm^{(1)} = 5$. Column 4 of tab. 2 demonstrates that even the radius $r_{10}$ where $Rm^{(1)} = 10$ lies deeper than $r_{\phi}$ for most flow and conductivity combinations. The only exceptions are the results for the geostrophic flow. This could indicated that strictly geostrophic flows would indeed violate the heating constraint.
Figure 5: Maps of (a) the radial surface field in the Jupiter field model JRM09, (b) the radial LIF at \( r_1 = 0.965 \, r_J \), and (c) the local Ohmic heating \( q = \int \, \frac{dr \, \vec{f}^2}{\sigma} \) at \( r = r_1 = 0.965 \, r_J \). Flow model \( U_Z \), conductivity \( \sigma_F \), and current density estimates \( j^{(O)} \) (for panel c) have been used. Outward (inward) directed field is shown in red (blue).

Figure 6: (a) \( U_Z \) profiles at different depths; the grey background color highlights where surface jets are prograde. (b) Azimuthal rms radial LIF and (c) heat flux \( q = \int \, \frac{dr \, \vec{f}^2}{\sigma} \) at three different depths. Conductivity model \( \sigma_F \), flow \( U_Z \), and estimate \( j^{(O)} \) have been used.
Figure 7: Power spectra of rms radial field contributions (Mauersberger-Lowes) for JRM09, the downward continued $B_r$ and the radial LIF $\hat{B}_r$ at $r_1 = 0.965 \, r_J$. (a) shows the spherical harmonic degree spectrum, while (b) shows the harmonic order spectrum. The LIF has been amplified with $10^3$. Flow $U_Z$ and conductivity $\sigma_F$ have been used.

Figure 8: Profiles of (a) Ohmic heating and (b) entropy production in the layer above radius $r$ for current estimate $j^O$. In (a) the solid horizontal line shows the total convective power of $1.2 \times 10^{18}$ W, while the dotted horizontal line shows the heat flow of $Q_o = 3.35 \times 10^{17}$ W out of Jupiter’s interior. In (b) the horizontal line indicates the total dissipative entropy production predicted by the entropy flux $\theta_T = Q_o/T_o = 2.0 \times 10^{15}$ W/K through the outer boundary. Vertical lines mark the radii $r_1$ where $\text{Rm}^{(1)} = 1$ (see fig. 3).
5 Discussion and Conclusion

The dominance of Ohmic dissipation in the outer few percent of Jupiter’s radius leads to simple quasi-stationary dynamo action. This can be exploited for estimating the electric currents and the Locally Induced Fields with surprisingly high quality (Wicht et al., 2019), once a conductivity profile, a surface magnetic field model, and flow model are given. Here we explored two conductivity profiles, used the new Juno-based JRM09 field model, and tested two zonal flow models suggested from inversions of Juno gravity measurements. A geostrophic zonal flow model was also considered as a third option.

The estimates roughly apply to the upper four percent in radius, or roughly 3000 km, where the modified magnetic Reynolds number Rm(1) is smaller than one. The radial LIF in this quasi-stationary dynamo region typically reaches rms values in the order of µT with peak values up to 15 µT. Could such a small contribution be measured by the Juno magnetometer? The instrument has been designed to provide a nominal vector accuracy of 1 in 10^4. Since the surface field reaches peak values of about 2 mT, the LIF could indeed be detectable.

One would still have to separate the LIF from contributions produced deeper in the planet. What should help with this task, is the distinct pattern imprinted by the zonal flows, which also leads to a distinct magnetic spectrum. The LIF spectrum peaks at degree ℓ = 12 and has significant contributions at even higher degrees. At ℓ = 10, the largest degree provided by JRM09, the LIF amounts to about 1% of the background field, which seems smaller than the estimated JRM09 precision (Connerney et al., 2013). Updated future models, based on a larger number of Juno orbits, will provide smaller scale details and increase the chances of identifying the LIF. Another possibility is a dedicated analysis of measurements around the ‘big blue spot’ in JRM09, which is convincingly shown that the zonal winds must be significantly weaker below about 0.96 r_J, they cannot uniquely constrain their structure or amplitude at this depth.

Our analysis of Jupiter’s heat balance shows that Ohmic heating can significantly exceed the heat flux Q_o out of the planet’s interior. Using the interior model by Nettelmann et al. (2012) and French et al. (2012) suggests a total dissipative heating of Φ_T = 3.58 Q_o = 1.20×10^{18} W.

It would be interesting to repeat this assessment for the newer Jupiter models that include stably stratified regions (Debras & Chabrier, 2019). However, the most important constraint is the knowledge of Q_o, and the somewhat different distribution of internal heat sources implied by the newer models can only have a limited effect.

While the total Ohmic heating remains typically one order of magnitude below Φ_T, we find extreme lateral variations. Peak values in the Ohmic heating density reach 25 W/m^2 around the ‘blue spot’ in JRM09, which is nearly five times larger than the mean heat flux density from Jupiter’s interior. These peak values are reached at the bottom of the quasi-stationary region, i.e. at a depth of 3000 km. This is much deeper than any (current) remote instrument could reach for. For example MWR, the micro-wave instrument on Juno, hopes to detect temperature radiation from up to 1 kbar, which corresponds to a depth of about 600 km. However, the local heating may trigger convective plumes that rise to shallower depths and thus become detectable.

We also estimated the entropy flux out of Jupiter’s interior to 2.0×10^{15} W/K. The entropy produced by zonal wind related Ohmic heating in the quasi-stationary region does not exceed this value for any model combination. This means that neither Ohmic heating nor the entropy production offer any reliable constraint on the depth of the zonal winds.

Below the quasi-stationary region, electric fields become a significant contribution to Ohm’s law, tend to oppose induction effects, and lead to weaker electric currents than predicted by our approximations. (Wicht et al., 2019) demonstrate that the currents remain roughly constant below the depth where Rm(1) ≈ 5 in their numerical simulations. However, this may be different in Jupiter where the magnetic Reynolds numbers reach values orders of magnitude higher than in their computer models.

Fig. 5 demonstrates that Rm(1) increases to a value of at least 10^3 at 0.90 r_J. This is a consequence of the electrical conductivity profiles that easily overcompensate the depth-decrease in zonal flow velocities indicated by Juno gravity measurements. The zonal flows may thus actually play a larger role for dynamo action below than in the quasi-stationary region. While the gravity data convincingly show that the zonal winds must be significantly weaker below about 0.96 r_J, they cannot uniquely constrain their structure or amplitude at this depth.

It has been speculated that the fast observed zonal winds may remain confined to a thin weather layer, where differential solar heating and also moist convection could
significantly contribute to the dynamics (see for example Showman (2007) or Thomson & McIntyre (2016)). Kong et al. (2018) show that the gravity signal can then be explained by an independent zonal flow system that reaches down to about $0.7 r_J$ with typical amplitudes of about 1 m/s and has larger latitudinal scales than the surface winds. The strongest local dynamo action happens towards the bottom of the quasi-stationary region where models $U_K$ and $U_Z$ reach velocities of about $10 \text{ m/s}$. The currents and magnetic fields induced by this alternative flow model should thus be roughly an order of magnitude weaker than for $U_K$ or $U_Z$. Consequently, Ohmic heating and entropy production would be two orders of magnitude lower and play practically no role for the global power or entropy budgets.

Below $0.96 r_J$, full 3d numerical simulations would be required to model the zonal-wind related dynamo action. However, since they cannot be run at altogether realistic parameters and generally yield a much simpler zonal wind pattern, the results must be interpreted with care (Gastine et al. 2014; Jones 2014; Duarte et al. 2018; Dietrich & Jones 2018). These simulation suggest that even the weaker zonal winds at depth would significantly shear the large scale field produced by the deeper primary dynamo action. The resulting strong longitudinal (toroidal) flux bundles are converted into observable radial field by the small scale convective flows present in this region. The combined action of primary and secondary dynamo typically yields a radial surface field that is characterized by longitudinal banded structures and large scale patches with wavenumber one or two, resulting in a morphology is often reminiscent of the recent Juno-based field model JRM09 (Gastine et al. 2014; Duarte et al. 2018; Dietrich & Jones 2018).

**References**

Braginsky, S. & Roberts, P. 1995, Geophys. Astrophys. Fluid Dyn., 79, 1

Cao, H. & Stevenson, D. J. 2017, J. Geophys. Res., 122, 686

Christensen, U. & Aubert, J. 2006, Geophys. J. Int., 116, 97

Christensen, U. R., Holzwarth, V., & Reiners, A. 2009, Nature, 457, 167

Connerney, J. E. P., Kotsiaros, S., Oliversen, R. J., et al. 2018, Geophys. Res. Lett., 45, 2590

Currie, L. K. & Browning, M. K. 2017, Astrophys. J., 845, L17

Davidson, P. A. 2013, Geophys. J. Int., 195, 67

Debras, F. & Chabrier, G. 2019, Astrophys. J., 872, 100

Dietrich, W. & Jones, C. A. 2018, Icarus, 305, 15

Duarte, L. D. V., Wicht, J., & Gastine, T. 2018, Icarus, 299, 206

French, M., Becker, A., Lorenzen, W., et al. 2012, Astrophys. J. Supp., 202, 5

Galanti, E., Kaspi, Y., Miguel, Y., et al. 2019, Geophys. Res. Lett., 46, 616

Gastine, T., Wicht, J., Duarte, L. D. V., Heimpel, M., & Becker, A. 2014, Geophys. Res. Lett., 41, 5410

Gubbins, D., Alfè, D., Masters, G., Price, G., & Gillan, M. 2003, Geophys. J. Int., 155, 609

Gubbins, D., Masters, T. G., & Jacobs, J. A. 1979, Geophys. J., 59, 57

Guillot, T. & Gautier, D. 2015, in Treatise on Geophysics, Vol. 10, Planets and Moons, second edn. edn., ed. T. Spohn (Elsevier), 439–464

Guillot, T., Miguel, Y., Militzer, B., et al. 2018, Nature, 555, 227

Heimpel, M., Gastine, T., & Wicht, J. 2016, Nature Geosci., 9, 19

Hewitt, J. M., McKenzie, D. P., & Weiss, N. O. 1975, J. Fluid Mech., 68, 721

Iess, L., Folkner, W. M., Durante, D., et al. 2018, Nature, 555, 220

Iess, L., Militzer, B., Kaspi, Y., et al. 2019, Science

Jones, C. A. 2014, Icarus, 241, 148
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