

Breaking Taylor-Proudman Balance by Magnetic Fields in Stellar Convection Zones

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Abstract

We carry out high-resolution calculations for the stellar convection zone. The main purpose of this Letter is to investigate the effect of a small-scale dynamo on the differential rotation. The solar differential rotation deviates from the Taylor–Proudman state in which the angular velocity does not change along the rotational axis. To break the Taylor–Proudman state deep in the convection zone, it is thought that a latitudinal entropy gradient is required. In this Letter, we find that the small-scale dynamo has three roles in the deviation of the stellar differential rotation from the Taylor–Proudman state. 1) The shear of the angular velocity is suppressed. This leads to a situation where the latitudinal entropy gradient efficiently breaks the Taylor–Proudman state. 2) The perturbation of the entropy increases with the suppression of the turbulent velocity between upflows and downflows. 3) The convection velocity is reduced. This increases the effect of the rotation on the convection. The second and third factors increase the latitudinal entropy gradient and break the Taylor–Proudman state. We find that an efficient small-scale dynamo has a significant impact on the stellar differential rotation.

Key words: dynamo - stars: interiors - Sun: interior

Supporting material: animation

1. Introduction

The Sun is rotating differentially. The structure of the differential rotation is one of the most important factors in the generation of the magnetic field, i.e., the dynamo, because the differential rotation stretches the large-scale poloidal field to the toroidal field (Parker 1955). The detailed profile of the solar differential rotation has been revealed by global helioseismology (e.g., Schou et al. 1998). The profile shows the three features that deviate from the Taylor-Proudman state, where the angular velocity has a cylindrical profile, i.e., $\partial \langle \Omega \rangle / \partial z = 0$, where Ω and z are the angular velocity and the direction of the rotation axis, respectively. The angle brackets, $\langle \rangle$, show the longitudinal average. The tachocline, the near surface shear layer, and the conical profile of the angular velocity in the bulk of the convection zone are the deviations. The Taylor-Proudman theorem is derived from the longitudinal component of the vorticity equation as (for details see Rempel 2005; Hotta & Yokoyama 2011)

$$\frac{\partial \langle \omega_{\phi} \rangle}{\partial t} = (\nabla \times \langle \boldsymbol{v} \times \boldsymbol{\omega} \rangle)_{\phi} + \lambda \frac{\partial \langle \Omega \rangle^2}{\partial z} - \frac{g}{c_{\rm p} r} \frac{\partial \langle s \rangle}{\partial \theta} + \left(\nabla \times \left\langle \frac{1}{4\pi\rho} (\nabla \times \boldsymbol{B}) \times \boldsymbol{B} \right\rangle \right)_{\phi}, \tag{1}$$

where r, θ , and ϕ show the radius, the colatitude, and the longitude in the spherical geometry, respectively. $\lambda = r \sin \theta$ is the distance from the rotational axis. ρ , v, ω , Ω , s, g, c_p , and Bshow the density, the fluid velocity in the rotating frame, the vorticity, the angular velocity, the specific entropy, the gravitational acceleration, the heat capacity at constant pressure, and the magnetic field, respectively. When all of the other terms in Equation (1) are ignored, the Taylor– Proudman theorem, $\partial \langle \Omega \rangle / \partial z = 0$, is derived. The deviation of the Taylor–Proudman state should be explained with the transport term, $\nabla \times \langle v \times \omega \rangle$, the baroclinic term (latitudinal entropy gradient), $-g/(c_p r)\partial \langle s \rangle / \partial \theta$, or the Lorentz force term. Hotta et al. (2015b) suggest that the near surface shear layer is maintained by the transport term as a consequence of a sheared meridional flow. The tachocline and the conical profile of the differential rotation are thought to be maintained by the baroclinic term (Rempel 2005; Miesch et al. 2006; Brun et al. 2011). Rempel (2005) adopted the mean field model, i.e., the turbulence was not solved, and suggested that the interaction of the meridional flow and the subadiabatic overshoot region can maintain the latitudinal entropy gradient, and that the Taylor-Proudman state is broken. In the threedimensional (3D) model, Miesch et al. (2006) put the latitudinal entropy gradient at the bottom boundary and maintained the conical profile of the differential rotation. Brun et al. (2011) include the radiation zone and maintain the large entropy gradient in the overshoot region self-consistently, and as a result, the tachocline is also generated.

On the other hand, it has been reported that the convection zone itself has the ability to make the latitudinal entropy gradient (Miesch et al. 2000). The cool downflow is bent to the equator by the Coriolis force, and the hot upflow is bent to the pole. As a result, the anisotropy of the thermal convection is able to form a negative latitudinal entropy gradient. We call this the CZ effect in this paper. Miesch et al. (2000) reported that this effect is not enough to explain the observed differential rotation, i.e., the generated entropy gradient is too small. In this Letter, we further explore the possibility of the CZ effect in 3D convection calculations.

One possibility for amplifying the CZ effect is the magnetic field. Hotta et al. (2015a, 2016) show that in high-resolution calculations, the magnetic field is on an equipartition level with the kinetic energy of the turbulent flow and that it significantly modifies the convective structure. In particular, the entropy structure is modified because the small-scale turbulent velocity between the up- and downflows is suppressed by the small-scale magnetic field. This has the possibility of changing the

CZ effect for the latitudinal entropy gradient. Hotta et al. (2016) have already shown that the high efficiency of the small-scale dynamo changes the distribution of the differential rotation (see their Figure 2 in the Supplementary Material). When the small-scale dynamo is very effective, the differential rotation tends to become a non-Taylor-Proudman state. A similar feature is seen in the high magnetic Prandtl number (low magnetic diffusivity) calculation in Käpylä et al. (2017). As Hotta et al. (2016) use the solar rotation rate, they adopt a very large thermal conduction on the entropy to avoid a high convection velocity and the resulting polar acceleration. Recently, a large number of studies have argued that the high-resolution calculations with the solar rotation rate tend to reproduce the high convection velocity and cause polar acceleration of the differential rotation, which is not consistent with the real Sun (e.g., Gastine et al. 2014; Fan & Fang 2014; Käpylä et al. 2014; Hotta et al. 2015b). To avoid this situation, Hotta et al. (2016) added a very large thermal conductivity to suppress the convection velocity. Meanwhile, the small-scale entropy feature is also suppressed in the calculation, and it becomes difficult to discuss the role of the small-scale dynamo in the creation of the latitudinal entropy gradient. In this Letter, we decided to change the rotation rate to $3\Omega_{\odot}$ and exclude the large thermal conduction on the entropy. Then we achieve an efficient small-scale dynamo in a relatively high-resolution calculation and discuss the importance of this small-scale dynamo on the latitudinal entropy gradient. We note that there will still be a discussion about the small-scale dynamo in a low magnetic Prandtl number, $P_{\rm m} = \nu/\eta$, where ν and η are the kinematic viscosity and the magnetic diffusivity, respectively (Käpylä et al. 2018). The low magnetic Prandtl number is achieved in the solar convection zone. This is beyond the scope of this study. Here we discuss the small-scale dynamo with $P_{\rm m} = 1$, which may be achieved only with numerical diffusivities.

2. Model

We solve the 3D magnetohydrodynamics (MHD) equations in the spherical geometry (r, θ, ϕ) . The equations solved in this study are listed as

$$\frac{\partial \rho_1}{\partial t} = -\frac{1}{\xi^2} \nabla \cdot (\rho \mathbf{v}),\tag{2}$$

$$\rho \frac{\partial \boldsymbol{v}}{\partial t} = -\rho(\boldsymbol{v} \cdot \nabla)\boldsymbol{v} - \rho_1 g \boldsymbol{e}_r + \nabla p_1 + 2\rho \boldsymbol{v} \,\boldsymbol{\Omega}_0 + \frac{1}{4\pi} (\nabla \times \boldsymbol{B}) \times \boldsymbol{B}, \qquad (3)$$

$$\frac{\partial \boldsymbol{B}}{\partial t} = \nabla \times (\boldsymbol{v} \times \boldsymbol{B}), \tag{4}$$

$$\rho T \frac{\partial s_1}{\partial t} = -\rho T (\mathbf{v} \cdot \nabla) s + Q_{\text{rad}}, \tag{5}$$

$$p_1 = \left(\frac{\partial p}{\partial \rho}\right)_s \rho_1 + \left(\frac{\partial p}{\partial s}\right)_\rho s_1,\tag{6}$$

where p and T are the gas pressure and the temperature, respectively; the subscript 0 and 1 are the steady background and perturbation quantities, respectively; ξ is the factor for the reduced speed of sound technique (Hotta et al. 2012). We make

the speed of sound uniform as

$$\xi(r) = \xi_0 \frac{c_{\rm s}(r)}{c_{\rm s}(r_{\rm min})},$$
(7)

$$c_{\rm s} = \sqrt{\left(\frac{\partial p}{\partial \rho}\right)_{\rm s}},\tag{8}$$

where $r_{\rm min}$ and $c_{\rm s}$ are the location of the bottom boundary and the adiabatic speed of sound, respectively. We adopt $\xi_0 = 100$. The background density, ρ_0 , temperature, T_0 , and the pressure, p_0 , are calculated using Model S (Christensen-Dalsgaard et al. 1996; Hotta et al. 2014). The factors for the linearized equation of state, $(\partial p/\partial \rho)_s$ and $(\partial p/\partial s)_{\rho}$, are calculated with the OPAL repository (Rogers et al. 1996). $Q_{\rm rad}$ includes the radiative heating around the base of the convection zone and the artificial cooling around the top boundary as

$$Q_{\rm rad} = \frac{1}{r^2} \frac{d}{dr} \bigg[r^2 \kappa_{\rm r} \rho_0 c_{\rm p} \frac{dT_0}{dr} - r^2 F_{\rm s}(r) \bigg],\tag{9}$$

$$4\pi r^2 F_{\rm s} = L_{\odot} \exp\left[-\left(\frac{r-r_{\rm max}}{d_{\rm s}}\right)^2\right],\tag{10}$$

where $F_{\rm s}$ is the artificial energy flux for the cooling around the top boundary with $d_{\rm s} = 18.8$ Mm, which is two times the local pressure scale height at the top boundary. $\kappa_{\rm r}$ is the radiative diffusivity calculated with the OPAL repository; $L_{\odot} = 3.84 \times 10^{33}$ erg s⁻¹ is the solar luminosity.

The radial calculation domain extends from $r_{\rm min} = 0.71 R_{\odot}$ to $r_{\rm max} = 0.96 R_{\odot}$, where $R_{\odot} = 6.96 \times 10^{10} \, {\rm cm}$ is the solar radius. The whole sphere is covered with the Yin-Yang grid (Kageyama & Sato 2004). We prepare the number of grid points of $96(r) \times 384(\theta) \times 1152(\phi) \times 2$ (Yin-Yang). In the analyses, we convert the Yin-Yang grid to the almost equivalent ordinary spherical geometry with the number of grid points as $96(r) \times 768(\theta) \times 1536(\phi)$. Because the near surface layer is not included, we cannot discuss the near surface shear layer in this study. In this Letter, we show two cases: HD without the magnetic field and MHD with the magnetic field. We add small random perturbation on the entropy for both cases as an initial condition. For the MHD case, an axisymmetric longitudinal magnetic field, $|B_{\phi}| = 100$ G, is imposed. The magnetic field is antisymmetric about the equator. We calculate the HD and MHD cases for 7000 and 14,000 days, respectively.

We solve the equation with the fourth-order space-centered derivative and the four-step Runge–Kutta method (Vögler et al. 2005). An artificial viscosity suggested by Rempel (2014) is also adopted to stabilize the numerical calculation. We do not adopt any explicit diffusivity to maximize the resolution.

Averages for the analyses are done in 5000–7000 and 10,000–14,000 days for the HD and MHD cases, respectively.

3. Results

The overall convection structures in the HD (panel (a)) and MHD (panel (b)) cases are shown in Figure 1. The perturbation of the entropy $s' = s - \langle s \rangle$ multiplied by the background density is shown. Compared with the result in Hotta et al. (2016), in which the thermal conductivity on the entropy is significantly large, we can observe the small-scale entropy structure in Figure 1. A rotationally aligned structure as a



Figure 1. 3D volume rendering of the perturbation of the specific entropy multiplied by the background density for the (a) HD and (b) MHD cases are shown. A corresponding movie for panel (b) is available online. The animation continues for seven seconds covering 40 days of the global convection. The turbulent nature of the thermal convection is clearly seen in the animation.

(An animation of this figure is available.)



Figure 2. Panels (a) and (c) show the streamline of the meridional flow for the HD and MHD cases, respectively. The solid and dashed lines show the clockwise and counterclockwise flows, respectively. The angular velocity, $\Omega/(2\pi)$, is shown for the HD and MHD cases in panels (b) and (d), respectively, where $\Omega = v_{\phi}/\lambda + \Omega_0$.

result of the Coriolis force is also observed (see also the corresponding movie online). Because of the high resolution, the small-scale dynamo is well excited. We confirm that the small-scale magnetic energy exceeds the kinetic energy at a small scale, $\ell > 200$, where ℓ is the spherical harmonic degree. The turbulent magnetic energy around the base of the convection zone ($<0.8 R_{\odot}$) exceeds the kinetic energy. These facts indicate that the flow is significantly affected by the small-scale magnetic field. This is seen in the comparison between panels (a) and (b) in Figure 1. During the average period (10,000–14,000 days in the MHD case), the large-scale magnetic field does not show any polarity reversal. We note that the polarity reversal is seen around 4500 and 8000 days.

Figures 2(a) and (c) show the streamline of the mass flux of the meridional flow $\rho_0 \langle \mathbf{v}_m \rangle$ for the HD and MHD cases, respectively, where \mathbf{v}_m includes the radial and latitudinal components of the

velocity. The meridional flow in the HD case (Figure 2(a)) shows a north-south aligned feature around the tangential cylinder. The overall flow pattern is counterclockwise (clockwise) in the northern (southern) hemisphere. In the MHD case (Figure 2(c)), we see a clockwise (counterclockwise) meridional flow cell around the base of the convection zone in the northern (southern) hemisphere. Figures 2(b) and (d) show the angular velocity for the HD and MHD cases, respectively. Because of the strong Lorentz force of the dynamo-generated magnetic field, the shear of the differential rotation, i.e., $\Delta\Omega$, decreased in the MHD case (see the difference in the color bar). In addition, the polar region is accelerated as well as the equator in the HD case. This feature is frequently seen in hydrodynamic calculations (Miesch et al. 2000; Gastine et al. 2012). The most important difference between the HD and MHD case concerns the Taylor-Proudman balance. In the HD case, the contour lines of the angular velocity

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Figure 3. Each term in Equation (11) in the MHD case is shown.



Figure 4. Latitudinal distributions of the entropy, \hat{s} , and the temperature, \hat{T} , are shown in the HD and MHD cases. Panels (a) and (b) show the result of the HD case, and (c) and (d) show the result of the MHD case. The entropy \hat{s} (temperature \hat{T}) is shown in panels (a) and (c) ((b) and (d)).

are almost parallel to the rotational axis (Figure 2(b)). This indicates that the HD case is in the Taylor–Proudman state. In contrast, the angular velocity in the MHD case shows significant deviation from the Taylor–Proudman state. This indicates that some terms in Equation (1) are balanced with the Coriolis force term.

To investigate the force balance on the meridional plane, we transpose Equation (1) as

$$COR = ADV + BAR + MAG,$$
(11)

$$\operatorname{COR} = -\lambda \frac{\partial \langle \Omega \rangle^2}{\partial z},\tag{12}$$

$$ADV = (\nabla \times \langle \boldsymbol{\nu} \times \boldsymbol{\omega} \rangle)_{\phi}, \tag{13}$$

$$BAR = -\frac{g}{c_{p}r}\frac{\partial\langle s\rangle}{\partial\theta},$$
 (14)

$$MAG = \left(\nabla \times \left\langle \frac{1}{4\pi\rho} (\nabla \times \boldsymbol{B}) \times \boldsymbol{B} \right\rangle \right)_{\phi}, \qquad (15)$$

where $\partial \langle \omega_{\phi} \rangle / \partial t$ is ignored, because the calculation is already in a statistically steady state. Figure 3 shows the balance in the MHD case. Minor contributions from the advection term ADV and the Lorentz force MAG around the top boundary are seen in Figures 3(b) and (d), but the Coriolis force (COR, Figure 3(a)) is mainly balanced with the baroclinic term, i.e., the entropy gradient (BAR, Figure 3(c)). Although the difference in the differential rotation is caused by including the magnetic field, the Lorentz force MAG is not the main contribution to the difference.

Figure 4 shows the latitudinal distributions of the entropy and the temperature in the HD and MHD cases. The hat for a value Q is expressed as

$$\hat{Q} = \langle Q \rangle - \bar{Q}, \tag{16}$$

where \bar{Q} is the horizontal average of the value Q. The value \hat{Q} is useful to show the latitudinal distribution while dismissing



Figure 5. Latitudinal distributions of the entropy $(\hat{s}, \text{ panel (a)})$ and the temperature $(\hat{T}, \text{ panel (b)})$ at the bottom boundary are shown. Panel (c) shows the rms entropy perturbation srms. Panel (d) shows the rms convection velocity. The dashed and solid lines show the radial and horizontal components. Black and red lines show the results of the HD and MHD cases, respectively.

the radial distribution. Figures 4(a) and (c) show \hat{s} in the HD and MHD cases, respectively. It is clear that the latitudinal entropy gradient is increased throughout the convection zone in the MHD case. The reasons for this are explained in the following paragraph. For reference, the latitudinal temperature distributions, \hat{T} , of the HD and MHD cases are shown in Figures 4(b) and (d), respectively. In the HD case, the temperature decreases to the pole from the middle latitude, whereas in the MHD case, the temperature increases monotonically to the pole. This is also shown in Figures 5(a) and (b)as a one-dimensional (1D) plot at the base of the convection zone.

There are two reasons that the latitudinal entropy gradient is increased in the MHD case. One is that the perturbation of the entropy is increased in this case. Figure 5(c) shows the rootmean-square (rms) entropy in the HD (black) and MHD (red) cases. It is clear that the rms entropy in the MHD case is increased compared with the HD case. This effect is already reported in the high-resolution calculation of the small-scale dynamo in a nonrotating Cartesian geometry (Hotta et al. 2015a). When the turbulent velocity between the upand downflows is suppressed by the small-scale magnetic field, the amplitude of the entropy perturbation is increased. This increases the anisotropy of the latitudinal energy flux, i.e., the CZ effect. The other reason is the suppression of the convection velocity. Figure 5(d) shows the rms velocity in the HD (black) and MHD (red) cases. The dashed and solid lines show the

radial and horizontal components, respectively. The convection velocity in the MHD case is clearly suppressed because of the Lorentz force. This indicates that the Coriolis force becomes more effective on the slow convection in the MHD case. This also increases the anisotropy of the latitudinal energy flux. This is confirmed by estimating the correlation of the entropy perturbation and the latitudinal velocity, $\langle s' v_{\theta} \rangle$, which is proportional to the latitudinal enthalpy flux. The average is done in the initial phase of the calculation (50-500 days), where the latitudinal entropy gradient is not well established and the contribution from the isotropic turbulent thermal conduction is small. We can see a poleward anisotropic energy flux in almost all of the locations in both of the cases. In the HD case, we see an equatorward energy flux in the small region close to the equator and the base of the convection zone. We summarize the reasons. The entropy in the upflow (downflow) becomes hotter (cooler) due to the magnetic field. There are several possible influences by the magnetic field to increase the entropy perturbation. While the magnetic field decreases the convection velocity, the total energy flux does not change and the entropy perturbation increases to compensate for the slow convection velocity. In addition, the time for the cooling around the top boundary becomes long due to slow convection velocity with the magnetic field. This also increases the entropy perturbation. Since the magnetic field becomes strong and the convective motion is significantly affected, the entropy perturbation and convection velocity are decoupled. This causes a deviation from the classical mixing length theory.

Then, the hot upflow (cool downflow) is bent poleward (equatorward) by the Coriolis force. Since the convection velocity is suppressed, this second step is also promoted by the magnetic field. As a result the latitudinal entropy gradient (hot pole and cool equator) is amplified by the magnetic field.

Regarding the deviation from the Taylor-Proudman state, the Lorentz force on the differential rotation also has a role. As explained, the Lorentz force is strong enough to reduce the shear of the differential rotation, $\Delta\Omega$. We transform Equation (11) by including only the terms COR and BAR as

$$-2\lambda \langle \Omega \rangle \frac{\partial \langle \Omega_{\rm l} \rangle}{\partial z} = -\frac{g}{c_{\rm p} r} \frac{\partial \langle s \rangle}{\partial \theta},\tag{17}$$

where $\Omega_1 = v_{\phi}/\lambda$. This relation shows that even with the same latitudinal entropy gradient, the Taylor-Proudman state is more easily broken with a small shear of the angular velocity, $\langle \Omega_1 \rangle$. Thus, the Lorentz force on the differential rotation also promotes the deviation of the Taylor-Proudman state.

4. Summary

In this Letter, we carry out a relatively high-resolution calculation of the stellar global convection in the spherical shell. We investigate the effects of the efficient small-scale dynamo on the shape of the differential rotation. Compared with our previous calculation (Hotta et al. 2016), we exclude the strong thermal conductivity on the entropy in order not to suppress the small-scale thermal features. Here we compare the calculations with and without the magnetic field. Because of the relatively high resolution, the efficient small-scale dynamo is achieved, even covering the whole convection zone.

One important finding is that the differential rotation deviates significantly from the Taylor-Proudman state when the magnetic field is included. The reasons for this deviation are summarized as follows:

- 1. The Lorentz force suppresses the differential rotation. Thus, the criterion, i.e., the required latitudinal entropy gradient to break the Taylor-Proudman balance, is relaxed.
- 2. The entropy perturbation is increased by suppressing turbulent velocity between the up- and downflows by the small-scale magnetic field. This increases the anisotropy of the latitudinal energy flux.
- 3. The convection velocity is suppressed by the magnetic field. Thus, the rotational effect is increased in the MHD calculation. This also increases the anisotropy of the latitudinal energy flux.

Reason 3 above possibly relates to reason 1. When the convection velocity is suppressed, the angular momentum transport by convection tends to be suppressed, leading to the reduction in the differential rotation. Reasons 2 and 3 increase the latitudinal entropy gradient that breaks the Taylor-Proudman balance. Reasons 2 and 3 are possibly achieved with a high effective thermal Prandtl number, in which the thermal conductivity is low and the kinetic viscosity is strong. There are several attempts at mimicking the small-scale dynamo and magnetic field by adopting a high effective Prandtl number (O'Mara et al. 2016; Bekki et al. 2017; Karak et al. 2018). This type of attempt is valuable, because the small-scale dynamo that

requires a high numerical cost can be reproduced with a reasonable cost. To mimic the small-scale dynamo correctly, detailed comparisons between the calculation in this study and calculations with high Prandtl numbers are required. In Karak et al. (2018), the generation of the negative latitudinal entropy gradient in their high Prandtl number calculations is discussed. They suggest that the entropy gradient is caused by the interaction of the meridional flow and the subadiabatic layer around the base of the convection zone. In this calculation, the subadiabatic layer around the base of the convection zone in the MHD case is confirmed, whereas the HD case does not have such a layer. The meridional flow, however, in this calculation, is clockwise (counterclockwise) around the base of the convection zone. This tends to generate the positive (negative) entropy gradient in the northern (southern) hemisphere. Thus, we conclude that the interaction between the meridional flow and the subadiabatic layer does not contribute to the deviation of the Taylor–Proudman state in this study.

In this study, we do not include the radiation zone (Brun et al. 2011), which also tends to amplify the latitudinal entropy gradient. In addition, the property of the small-scale dynamo is affected by the radiation zone (Hotta 2017). A combination of the small-scale dynamo and the radiation zone would determine the detailed shape of the stellar differential rotation. This will be addressed in the future.

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