Predicting the Rossby Number in Convective Experiments

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Abstract

The Rossby number is a crucial parameter describing the degree of rotational constraint on the convective dynamics in stars and planets. However, it is not an input to computational models of convection but must be measured ex post facto. Here, we report the discovery of a new quantity, the predictive Rossby number, which is both tightly correlated with the Rossby number and specified in terms of common inputs to numerical models. The predictive Rossby number can be specified independent of Rayleigh number, allowing suites of numerical solutions to separate the degree of rotational constraint from the strength of the driving of convection. We examine the scaling of convective transport in terms of the Nusselt number and the degree of turbulence in terms of the Reynolds number of the flow, and we find scaling laws nearly identical to those in non-rotational convection at low Rossby number when the predictive Rossby number is held constant. Finally, we describe the boundary layers as a function of increasing turbulence at constant Rossby number.

Key words: convection – dynamo – hydrodynamics – Sun: rotation – turbulence

Supporting material: machine-readable table

1. Introduction

Rotation influences the dynamics of convective flows in stellar and planetary atmospheres. Many studies on the fundamental nature of rotating convection in both laboratory and numerical settings have provided great insight into the properties of convection in both the rapidly rotating regime and the transition to the rotationally unconstrained regime (King et al. 2009, 2012, 2013; Schmitz & Tilgner 2009; Zhong et al. 2009; Julien et al. 2012; Ecke & Niemela 2014; Stellmach et al. 2014; Cheng et al. 2015; Gastine et al. 2016) The scaling behavior of heat transport, the nature of convective flow structures, and the importance of boundary layer–bulk interactions in driving dynamics are well known. Yet, we do not know of any simple procedure for predicting the magnitude of vortical flow gradients purely from experimental control parameters, such as bulk rotation rate and thermal input.

In the astrophysical context, many studies of rotating convection have investigated questions inspired by the solar dynamo (Glatzmaier & Gilman 1982; Busse 2002; Brown et al. 2008, 2010, 2011; Augustson et al. 2012; Guerrero et al. 2013; Käpylä et al. 2014). Even when these simulations nominally rotate at the solar rate, they frequently produce distinctly different behaviors than the true Sun, such as anti-solar differential rotation profiles (Gastine et al. 2014; Brun et al. 2017). It seems that these differences occur because the simulations produce less rotationally constrained states than the Sun. The influence of rotation results from the local shear gradients, and these are not direct input parameters. Recent simulations predict significant rotational influence in the deep solar interior, which can drastically affect flows throughout the solar convection zone (Featherstone & Hindman 2016; Greer et al. 2016). In the planetary context, the balance between magnetic and rotational forces likely leads to the observed differences between ice giant and gas giant dynamos in our solar system (Soderlund et al. 2015). The work of Aurnou & King (2017) demonstrates the importance of studying a dynamical regime with the proper balance between Lorentz, Coriolis, and inertial forces when modeling astrophysical objects such as planetary dynamos.

In short, simulations must achieve the proper rotational balance if they are to explain the behavior of astrophysical objects. In Boussinesq studies, rotational constraint is often measured by comparing dynamical and thermal boundary layers or deviation in heat transport from the non-rotating state (Julien et al. 2012; King et al. 2012, 2013). Such measurements are not available for astrophysical objects where the degree of rotational influence is best assessed by the ratio between non-linear advection magnitude and the linear Coriolis accelerations. The Rossby number is the standard measure of this ratio,

\[
\text{Ro} \equiv \left| \frac{\nabla \times \mathbf{u}}{2|\Omega|} \right| \sim \left| \frac{(\nabla \times \mathbf{u}) \times \mathbf{u}}{2|\Omega| \times \mathbf{u}} \right|,
\]

where \( \Omega \) denotes the bulk rotation vector. Many proxies for the dynamical Rossby number exist that are based solely on input parameters, most notably the convective Rossby number. However, all proxies produce imperfect predictions for the true dynamically relevant quantity.

In this paper, we demonstrate an empirical method of predicting the output Rossby number of convection in a simple stratified system.

In Anders & Brown (2017) (hereafter AB17), we studied non-rotating compressible convection without magnetic fields in polytropic atmospheres. In this work, we extend AB17 to...
Figure 1. (a) Critical Rayleigh number, as a function of the Taylor number, plotted as a solid black line. The gray shaded region is subcritical, and rotation suppresses convection there. Paths of constant convective Rossby number (\(Rop\), orange solid lines) are shown. From thickest to thinnest, paths with \(Rop = [1.58, 0.96, 0.6]\) are plotted, and the value of \((Ta_{crit}, Ra_{crit})\) for each path is denoted by a circular marker (see Table 1). (b) Evolved Ro plotted vs. Ra along paths of \(Rop = [1.58, 0.96, 0.6]\) for (large, medium, small) orange triangles. For comparison, paths of constant \(S\) (blue squares, \(S = [3, 2]\) for (large, small) squares) and constant \(Rop\) (green triangles, \(Rop = [1, 0.3, 0.1]\) for (large, medium, small) triangles) are shown. (c) Evolved value of Ro shown as a function of \(Rop\) and \(Ro\). Each of the experiments in (b) is outlined by a black (circle, triangle, square) for points along constant (\(Rop\), \(Rop\), \(S\)) paths. The color inside of the marker represents the exact measured Ro of that experiment, while the colormap outside of the markers is a linear interpolation of the data set.

rotationally influenced, \(f\)-plane atmospheres (e.g., Brummell et al. 1996, 1998; Calkins et al. 2015). We determine how the input parameters we studied previously, which controlled the Mach and Reynolds numbers of the evolved flows, couple with the Taylor number (\(Ta\); Julien et al. 1996), which sets the magnitude of the rotational vector.

In Section 2, we describe our experiment and paths through parameter space. In Section 3, we present the results of our experiments and in Section 4 we offer concluding remarks.

2. Experiment

We study fully compressible, stratified convection under precisely the same atmospheric model as in AB17, but here we have included rotation. We study polytropic atmospheres with \(n_p = 3\) density scale heights and a superadiabatic excess of \(\epsilon = 10^{-4}\) such that flows are at low Mach number. We study a domain in which the gravity, \(g = -g_z\), and rotational vector, \(\Omega = \Omega_z\), are antiparallel (as in, e.g., Brummell et al. 1996; Julien et al. 1996).

We evolve the velocity (\(u\)), temperature (\(T\)), and log density (\(\ln(\rho)\)) according to the fully compressible Navier–Stokes equations in the same form presented in AB17, with the addition of the Coriolis term, \(2\Omega \times u\), to the left-hand side of the momentum equation. We impose impermeable, stress-free, fixed-temperature boundary conditions at the top and bottom of the domain.

We specify the kinematic viscosity (\(\nu\)), thermal diffusivity (\(\chi\)), and strength of rotation (\(\Omega\)) at the top of the domain by choosing the Rayleigh number (\(Ra\)), Prandtl number (\(Pr\)), and \(Ta\),

\[
Ra = \frac{gL_z^3 \Delta S/c_p}{\nu \chi}, \quad Pr = \frac{\nu}{\chi}, \quad Ta = \left(\frac{2\Omega L_z}{\nu}\right)^2,
\]

where \(L_z\) is the depth of the domain as defined in AB17, \(\Delta S \propto \epsilon n_p\) is the specific entropy difference between the top and bottom of the atmosphere, and the specific heat at constant pressure is \(c_p = \gamma/(\gamma - 1)\) with \(\gamma = 5/3\). Throughout this work we set \(Pr = 1\). The Taylor number relates to the oft-quoted Ekman number by the equality \(Ek = Ta^{-1/2}\).

Due to stratification, \(Ra\) and \(Ta\) both grow with depth as \((Ra, Ta) \propto \rho^2\) (see AB17). We non-dimensionalize our atmospheres at the top of the domain, and so all values of \(Ra\) and \(Ta\) quoted in this work are the minimal values of \(Ra\) and \(Ta\) in the domain at \(z = L_z\). For direct comparison to Boussinesq studies, past work has found that the value of \(Ra\) at the atmospheric midplane \((z = L_z/2)\) varies minimally with increasing stratification (Unno et al. 1960). For the atmospheres presented in this work, midplane \(Ra\) and \(Ta\) values are larger than reported top-of-atmosphere values by a factor of \(~70\), and values at the bottom of the atmosphere are larger by \(~400\).

When \(Ta\) is large, the wavenumber of convective onset increases according to \(k_{crit} \propto Ta^{1/6}\) (Chandrasekhar 1961; Calkins et al. 2015). We study horizontally periodic, 3D Cartesian domains with extents of \(x, y = [0, 4(2\pi/k_{crit})]\) and \(z = [0, L_z]\). At large values of \(Ta\), these domains are tall and slender, as in Stellmach et al. (2014). We evolve our simulations using the Dedalus\(^5\) pseudospectral framework, and our numerical methods are identical to those presented in AB17. The supplemental materials of this paper include a tar file which contains the code used to perform the simulations in this work; this tar archive has been deposited in Zenodo (Anders et al. 2019).

The critical value of \(Ra\) at which rapidly rotating convection onsets also depends on \(Ta\) (see the black line in Figure 1(a)), roughly according to \(Ra_{crit} \sim Ta^{1/3}\) (Chandrasekhar 1961; Calkins et al. 2015). Even taking account of linear theory, the dependence of the evolved non-linear fluid flows on the input parameters makes predicting the rotational constraint very

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5 http://dedalus-project.org/


### Table 1

| Parameter Space | $(Ra_{crit}, Ta_{crit})$ | $(Ra_{max}, Ta_{max})$ |
|-----------------|-------------------------|------------------------|
| 0.60            | $(10^{8.88}, 10^{1.19})$| $(10^{9.97}, 10^{3.75})$|
| 0.96            | $(10^{4.44}, 10^{3.30})$| $(10^{8.58}, 10^{1.49})$|
| 1.58            | $(10^{1.39}, 10^{3.33})$| $(10^{7.14}, 10^{6.99})$|

**Note.** Values of the critical Ra and Ta for each $Ro_p$ track are reported, as well as the maximal values of Ra and Ta studied on each track. All values reported are for the top of the atmosphere. A fuller set of simulations is reported in Table 2 along with midplane Ra and Ta values.

challenging. We will explore three paths through Ra–Ta space:

$$Ra = \left\{ \begin{array}{ll}
S Ra_{crit}(Ta), & (I) \\
(Ro)^2 Pr Ta, & (II) \\
(Ro_p)^2 Pr^{1/2} Ta^{3/4}, & (III)
\end{array} \right.$$  

Paths on constraint I are at constant supercriticality, $S \equiv Ra/Ra_{crit}(Ta)$ (blue dashed–dotted line in Figure 1(a)). Paths on constraint II (green dashed line in Figure 1(a)) are at a constant value of the classic *convecti*on Rossby number,

$$Ro_c = \sqrt{\frac{Ra}{Pr Ta}} = \frac{1}{2\Omega} \sqrt{\frac{g}{\rho_0} \frac{\Delta S}{L_z}},$$

which has provided (e.g., Brummell et al. 1996; Julien et al. 1996) a common proxy for the degree of rotational constraint. This parameter measures the importance of buoyancy relative to rotation without involving dissipation. Paths on constraint III (e.g., orange solid lines in Figure 1(a)) set constant a ratio which we call the “prediction Rossby number,”

$$Ro_p = \sqrt{\frac{Ra}{Pr^{1/2} Ta^{3/4}}} = \frac{1}{(2\Omega)^{3/4}} \sqrt{\frac{g}{\rho_0} \frac{\Delta S}{\chi^{1/2}}}.$$  

Unlike paths through parameter space which hold $Ro_c$ constant, paths with constant $Ro_p$ feel changes in diffusivities but not the depth of the domain. To our knowledge, these paths have not been reported in the literature, although the importance of $Ra/Ta^{3/4} = Ra Ek^{3/2}$ has been independently found by King et al. (2012) using a boundary layer analysis. We compare our results to their theory in Section 4.

In this work, we primarily study three values of $Ro_p$. These values are shown in Figure 1(a) and Table 1. Table 1 lists the values of $(Ra_{crit}, Ta_{crit})$ for each value of $Ro_p$, and also the maximum value of $(Ra, Ta)$ studied in this work for each path. We additionally walked two pathways at constant supercriticality (constraint I, $S = [2, 3]$) and three pathways at constant convective Rossby number (constraint II, $Ro_c = [1, 0.3, 0.1]$). Full details on all cases are provided in the Appendix and the supplemental materials.

### 3. Results

In our stratified domains, for $Ta \geq 10^5$, a best fit to results from a linear stability analysis provides $Ra_{crit}(Ta) = 1.459 Ta^{7/3}$ and $k_{crit}(Ta) = 0.414 Ta^{1/6}$ for direct onset of convection. In Figure 1(a), the value of $Ra_{crit}(Ta)$ is shown. Sample paths for each criterion in Equation (3) through this parameter space are also shown. In this work, we often find it instructive to use one critical Ra for an entire $Ro_p$ path. This $Ra_{crit}$ is determined by the intersection of the onset curve and $Ro_p$ path (indicated by the orange circles in Figure 1(a), and quoted in Table 1). In the high-Ta regime, we find that $Ra_{crit} = 18.5 Ro_p^{15.16}$.

In Figure 1(b), we display the evolution of Ro with increasing Ra along various paths through parameter space. We find that Ro increases on constant Ro paths, decreases on constant $S$ paths, and remains roughly constant along constant $Ro_p$ paths. In Figure 1(c), the value of Ro is shown simultaneously as a function of $Ro_p$ and $Ro_c$ for all experiments conducted in this study. We find a general power law of the form Ro $= CRo^c Ro_p^d$. In the rotationally dominated regime where Ro $< 0.2$ and Re $> 5$ (see Equation (6)), we find $\alpha = -0.02$, and Ro can be said to be a function of $Ro_p$ alone. Under this assumption, we report a scaling of $Ro = (0.148 \pm 0.003)Ro_p^{0.34 \pm 0.07}$. In the less rotationally dominated regime of $Ro > 0.2$ and $Re > 5$, we find $\{C, \alpha, \beta\} = \{0.2, -0.19, 1.5\}$.

In Figure 2, sample snapshots of the evolved entropy field in the $x$–$y$ plane near the top and at the middle of the domain are shown. In the left column, flows are at $Ro \sim 1$ and resemble the classic granular structure of non-rotating convection (see, e.g., Figure 2 in AB17), where strong narrow downflow lanes punctuate broad upwellings. The narrow downflows at the top organize themselves into intense coherent structures at the midplane, and at the midplane the downflows have much stronger entropy fluctuations than the broad and slower upflows.

Ro decreases from left to right into the rotationally constrained regime. As Ro decreases, the narrow downflow lanes begin to disappear and the flows at the midplane become more symmetric. In the rotationally constrained regime ($Ro \sim 0.03$), the convective structures are distinctly different. Here we observe dynamically persistent, warm upflow columns surrounded by bulk weak downflow regions. At the midplane, the upflow columns have substantially higher entropy perturbations than the surrounding weak downflows that sheathe them, and the locations of the columns are tightly correlated with their positions at the top of the domain. These quasi-two-dimensional dynamics are similar to those seen in rapidly rotating Rayleigh–Bénard convection (e.g., Stellmach et al. 2014). The select cases displayed in Figure 2 each have an evolved volume-averaged $Re_\perp \approx 32$ (defined below in Equation (6)).

We measure the Nusselt number (Nu), which quantifies heat transport in a convective solution, as defined in AB17. In Figure 3(a), we plot Nu as a function of $Ra/Ra_{crit}$ at fixed $Ro_p$. We find that $Nu \propto [Ra_{0.29 \pm 0.01}, Ra_{0.29 \pm 0.01}, Ra_{0.23}]$ for $Ro_p = \{0.6, 0.957, 1.58\}$. In the regime of $Ro \lesssim 0.1$, these scaling laws are indistinguishable from a classic $Ra^{7/2}$ power-law scaling, which is observed in non-rotating Rayleigh–Bénard and stratified convection (Ahlers et al. 2009, AB17). Our results seem consistent with the stress-free, rotating Rayleigh–Bénard convection results of Schmitz & Tilgner (2009), whose re-arranged Equation (7) returns a best fit of $Nu \propto Ra_{0.26}$ at fixed $Ro_p$. Their work primarily spans the transition regime between rotationally constrained and unconstrained convection, and so it is perhaps not surprising that their power law is a blend of our rotationally constrained $Ra^{7/2}$ power law and the fairly rotationally unconstrained $Ra_{0.24}$ at $Ro_p = 1.58$.

Flows are distinctly different parallel to and perpendicular from the rotation vector, which aligns with gravity and stratification. We measure two forms of the rms Reynolds
number,
\[
\text{Re}_\parallel = \frac{|u| L_z}{\nu}, \quad \text{Re}_\perp = \frac{|u|}{\nu} \frac{2\pi}{k_{\text{crit}}},
\]
where the length scale in Re$_\perp$ is the wavelength of convective onset, and is related to the horizontal extent of our domain (see Section 2). From our work in AB17, we expect the rms velocity to scale as $|u| \propto \sqrt{\Delta S}$. By definition, $\nu \propto \sqrt{Ra/(Pr \Delta S)}$, and $L_z$ is a constant set by the stratification while $k_{\text{crit}} \propto Ta^{1/6}$. Along paths of constant Ro, we thus expect Re$_\parallel \propto Ra^{1/2}$ and Re$_\perp \propto Ra^{4/18}$ when Pr is held constant.

In Figure 3(b), we plot Re$_\parallel$ and Re$_\perp$ as a function of Ra/Ra$_{\text{crit}}$ at fixed Ro$_p$. We find that Re$_\parallel \propto \{Ra^{0.44 \pm 0.01}, Ra^{0.45 \pm 0.01}, Ra^{0.44}\}$ and Re$_\perp \propto \{Ra^{0.22 \pm 0.01}, Ra^{0.23 \pm 0.01}, Ra^{0.21}\}$ for Ro$_p = \{0.6, 0.957, 1.58\}$. These scalings are similar to but slightly weaker than our predictions in all cases. However, the scaling of Re$_\parallel \propto Ra^{0.45}$ is once again a power law observed frequently in non-rotating convection (Ahlers et al. 2009, AB17). We also observe that Re$_\parallel$ collapses for each Ro$_p$ track, while Re$_\parallel$ experiences an offset to larger values as Ro$_p$ shrinks. The offset in Re$_\parallel$ is unsurprising, because more rotationally constrained flows result in smaller boundary layers relative to the vertical extent of our stratified domain. The horizontal extent of our domain scales with the strength of rotation, and so, regardless of Ro$_p$, flows perpendicular to the rotational and buoyant direction are comparably turbulent at the same Ra/Ra$_{\text{crit}}$. We find Re$_\parallel$ and Re$_\perp$ are, respectively, good proxies for the horizontal and perpendicular resolution required to resolve an experiment.

Figure 4 shows time- and horizontally averaged profiles of Ro and the standard deviation of the entropy, $\sigma_x$. Figures 4(a) and (b) show these profiles for Ro$_p = 1.58$ (Ro $\approx 0.4$), while Figures 4(c) and (d) show these profiles for Ro$_p = 0.96$ (Ro $\approx 0.1$). The transition in profile behavior from low Ra (yellow) to high Ra (purple) is denoted by the color of the profile. As Ro increases at a constant value of Ro$_p$, both the thermal ($\sigma_x$) and dynamical (Ro) boundary layers become thinner. We measure the thickness of the thermal boundary layer ($\delta_x$) at the top of the domain by finding the location of the first maxima of $\sigma_x$ away from the boundary. We measure the thickness of the Ro boundary layer ($\delta_{Ro}$) in the same manner. In Figure 4(e), we plot $\delta_{Ro}/\delta_x$, the ratio of the sizes of these two boundary layers. As anticipated, the dynamical boundary layer ($\delta_{Ro}$) becomes thinner with respect to the thermal boundary layer ($\delta_x$) as Ro and Ro$_p$ decrease. However, the precise scaling of this boundary layer ratio with Ro$_p$ and Ra is unclear, and we cannot immediately compare these ratios to similar measures from the Rayleigh–Bénard convection literature, such as Figure 5 of King et al. (2013). They measure the dynamical boundary layer thickness as the peak location of the horizontal velocities, but our horizontal velocities are subject to stress-free boundary conditions, and we find that the maxima of horizontal velocities occur precisely at the boundaries. In Figure 4(f), we plot $\delta_x$ in units of the density scale height at the top of the atmosphere, and we plot vertical lines when this crosses 1. We find no systematic change in behavior when $\delta_x$ is smaller than the local density scale height.

### 4. Discussion

We studied low-Mach-number, stratified, compressible convection under the influence of rotation. We examined three paths through Ra–Ta space, and showed that the newly defined predictive Rossby number, $Ro_p = Ra/(Pr^{1/2}Ta^{1/4})$, determines the value of the evolved Rossby number. Astonishingly, along these constant Ro$_p$ pathways, particularly when Ro $\lesssim 0.1$, we find Nu $\propto Ra^{2/7}$ and Re$_\parallel \propto Ro^{0.45}$. These scalings are indistinguishable from those of Re and...
Figure 3. Scaling laws for paths at \( R_{op} = 1.58 \) (\( Ro \approx 0.4 \)), \( R_{op} = 0.96 \) (\( Ro \approx 0.1 \)), and \( R_{op} = 0.6 \) (\( Ro \approx 0.03 \)). Numbers are plotted vs. \( Ra/Ra_{crit} \), where \( Ra_{crit} \) is given in Table 1. (a) \( Nu \), as defined in AB17. (b) \( Re_{\parallel} \) and \( Re_{\perp} \), as defined in Equation (6). All values of \( R_{op} \) trace out similar \( Nu \) and \( Re_{\perp} \) tracks, whereas \( Re_{\parallel} \) tracks shift upwards as \( Ro \) decreases.

Figure 4. Horizontally averaged profiles of the standard deviation of entropy \((\sigma_s, a)\) and Rossby number \((Ro, b)\) vs. height for \( R_{op} = 1.58 \) (\( Ro \approx 0.4 \)). Similar profiles are shown in (c) and (d) for \( R_{op} = 0.96 \) (\( Ro \approx 0.1 \)). The color of the profiles denotes the value of \( Ra/Ra_{crit} \), where \( Ra_{crit} \) is given in Table 1. (c) The ratio of the thicknesses of the dynamical boundary layers \((\delta_{Ro})\) and thermal boundary layers \((\delta_s)\) is shown vs. \( Ra/Ra_{crit} \) for fixed \( R_{op} \). (f) \( \delta_s \) is plotted vs. \( Ra/Ra_{crit} \) in units of the density scale height at the top of the atmosphere \((H_p)\). Vertical lines denote when \( \delta_s/H_p = 1 \) for each value of \( R_{op} \).
Nu with Ra in non-rotating Boussinesq convection (Ahlers et al. 2009). Julien et al. (2012) theorized that in the rapidly rotating asymptotic limit, \( \text{Nu} \propto (\text{Ra}^{5/2}/\text{Ta}) = (\text{Ra}/\text{Ra}_{\text{crit}}(\text{Ta}))^{5/2} \). Thus, at fixed Ta, a very sharp Ra^{5/2} scaling law is expected. At a fixed Ta = 10^{14}, Stellmach et al. (2014) found that the Ra^{5/2} scaling described the results of stress-free direct numerical simulations in Boussinesq cylinders very well. Gastine et al. (2016) studied Boussinesq convection in spherical shells with no-slip boundaries, and also found good agreement with the theory of Julien et al. (2012) for various Ra at Ta \( \gg 10^{10} \).

Here, when we run simulations at fixed \( \text{Ro}_p \), the value of Ta is coupled to that of Ra, and both increase simultaneously. Recasting the scaling of Julien et al. (2012) into this perspective, we find \( \text{Nu} \propto \text{Ra}^{3/2}/\text{Ta} = \text{Ro}_p^{8/3}\text{Ra}^{1/6} \propto (\text{Ra}/\text{Ra}_{\text{crit}})^{1/6} \), where in this final result we use the \( \text{Ra}_{\text{crit}} \) value of the whole \( \text{Ro}_p \) path, such as those specified in Table 1. This \( \text{Ra}^{1/6} \) scaling is much weaker than the \( \text{Ra}^{2/3} \) law we find here. We leave it to future work to explain this discrepancy between Boussinesq theory and our observed Nu versus Ra scaling.

In this work, we experimentally arrived at the \( \text{Ra}/\text{Ta}^{3/4} = \text{Ra} \text{Ek}^{3/2} \) scaling in \( \text{Ro}_p \), but this relationship was independently discovered by King et al. (2012). Arguing that the thermal boundary layers should scale as \( \delta_\text{T} \propto \text{Ra}^{-1/3} \) and rotational Ekman boundary layers should scale as \( \delta_\text{Ro} \propto \text{Ta}^{-1/4} = \text{Ek}^{1/2} \), they expect these boundary layers to be equal in size when \( \text{Ra}/\text{Ta}^{3/4} \sim 1 \). They demonstrate that, when \( 2 \lesssim \text{Ra}/\text{Ta}^{3/4} \lesssim 20 \), flows are in the transitional regime, and for \( \text{Ra}/\text{Ta}^{3/4} \lesssim 2 \), flows are rotationally constrained. We remind the reader that Boussinesq values of Ra and Ta are not the same as their values in our stratified domains here, as diffusivities change with depth (see Section 2). Taking into account this change with depth, our simulations fall in King et al. (2012)’s rotationally constrained \( (\text{Ro}_p = 0.6) \) and near-constrained transitional regime \( (\text{Ro}_p = (0.957, 1.58)) \). The measured values of Ro in Figure 1(b) and the observed dynamics in Figure 2 agree with this interpretation.

We note briefly that the scaling \( \text{Ra} \propto \text{Ta}^{3/4} \) is very similar to another theorized boundary between fully rotationally constrained convection and partially constrained convection predicted in Boussinesq theory, of \( \text{Ra} \propto \text{Ta}^{4/5} \) (Julien et al. 2012; Gastine et al. 2016). This \( \text{Ta}^{4/5} \) scaling also arises through arguments of geostrophic balance in the boundary layers, and is a steeper scaling than the \( \text{Ta}^{3/4} \) scaling present in \( \text{Ro}_p \). This suggests that, at sufficiently low \( \text{Ro}_p \), a suite of simulations across many orders of magnitude of Ra will not only have the same volume-averaged value of Ro (as in Figure 1(b)), but will also maintain proper force balances within the boundary layers.

Our results suggest that by choosing the desired value of \( \text{Ro}_p \), experimenters can select the degree of rotational constraint present in their simulations. We find that \( \text{Ro} \propto \text{Ro}_p^{3.34 \pm 0.07} \), which is within 2\( \sigma \) of the estimate in King et al. (2013) who, although defining Ro very differently from our vorticity-based definition here, find \( \text{Ro} \propto \text{Ro}_p^{3.84 \pm 0.28} \). We note briefly that they claim that the value of Ro is strongly dependent upon the Prandtl number studied, and that low Ro can be achieved at high Pr without achieving a rotationally constrained flow. We studied only \( \text{Pr} = 1 \) here, and leave it to future work to determine if the scaling of \( \text{Ro}_p \propto \text{Pr}^{-1/4} \) is the correct scaling to predict the evolved Rossby number.

Despite the added complexity of stratification, and despite our using stress-free rather than no-slip boundaries, the boundary layer scaling arguments put forth in King et al. (2012) seem to hold up in our systems. This is reminiscent of what we found in AB17, in which convection in stratified domains, regardless of Mach number, produced boundary-layer-dominated scaling laws of Nu that were nearly identical to the scaling laws found in Boussinesq Rayleigh–Bénard convection.

We close by noting that once \( \text{Ro}_p \) is chosen such that a convective system has the same Rossby number as an astrophysical object of choice, it is straightforward to increase the turbulent nature of simulations by increasing Ra, just as in the non-rotating case. Although all the results reported here are for a Cartesian geometry with antiparallel gravity and rotation, preliminary 3D spherical simulations suggest that \( \text{Ro}_p \) also specifies Ro in more complex geometries (B. P. Brown et al. 2019, in preparation).

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Appendix

Table of Simulations

Information for select simulations in this work are shown in Table 2. The simulation at minimum (Ra, Ta) and maximum (Ra, Ta) for each of the \( \text{Ro}_p \), \( \text{Ro}_c \), and \( \delta \) paths in Figure 1(b) are shown. This information for the displayed simulations and all other simulations in this work is included in the online version of Table 2 and is deposited in Zenodo (Anders et al. 2019).
Table 2
Table of Simulation Information

| $R_a$ | $T_a$ | $R_{p, a}$ | $S$ | $L_x/L_z$ | $R_o$ | $R_{e, a}$ | $R_{e, z}$ | Nu |
|-------|-------|------------|-----|-----------|-------|------------|----------|
| 1.8 x $10^3$ | 4.1 x $10^3$ | 0.60 | 1.4 x $10^3$ | 3.0 x $10^6$ | 1.03 | 0.067 | 1.1 | 0.51 | (256, 32, 32) | 0.015 | 19.4 | 2.5 | 1.2 |
| 1.2 x $10^4$ | 5.2 x $10^3$ | 0.60 | 9.2 x $10^4$ | 3.8 x $10^4$ | 1.03 | 0.015 | 3.0 | 0.07 | (2048, 64, 64) | 0.026 | 1771 | 32.0 | 15.4 |
| 5.2 x $10^4$ | 4.6 x $10^3$ | 0.96 | 3.8 x $10^4$ | 3.4 x $10^3$ | 1.64 | 0.333 | 1.13 | 2.28 | (64, 64, 64) | 0.074 | 4.2 | 2.5 | 1.2 |
| 3.8 x $10^4$ | 3.1 x $10^4$ | 0.96 | 2.8 x $10^3$ | 6.3 x $10^3$ | 1.64 | 0.055 | 6.0 | 0.12 | (2048, 64, 64) | 0.129 | 3906 | 113 | 66.9 |
| 6.1 x $10^5$ | 5.0 x $10^3$ | 1.58 | 5.8 x $10^3$ | 7.4 x $10^3$ | 2.70 | 0.888 | 1.56 | 4.44 | (64, 64, 64) | 0.303 | 4.4 | 4.9 | 1.7 |
| 1.4 x $10^6$ | 9.7 x $10^3$ | 1.58 | 1.0 x $10^3$ | 7.2 x $10^10$ | 2.70 | 0.119 | 10.0 | 0.30 | (512, 128, 128) | 0.376 | 1257 | 94.9 | 40.1 |

| $R_o$ | $T_o$ | $R_{p, o}$ | $S$ | $L_x/L_z$ | $R_o$ | $R_{e, o}$ | $R_{e, z}$ | Nu |
|-------|-------|------------|-----|-----------|-------|------------|----------|
| 8.6 x $10^4$ | 8.6 x $10^4$ | 0.74 | 6.3 x $10^4$ | 6.3 x $10^6$ | 1.26 | 0.1 | 1.47 | 0.68 | (128, 128, 128) | 0.051 | 40.2 | 6.7 | 2.2 |
| 2.6 x $10^4$ | 2.6 x $10^4$ | 1.13 | 1.9 x $10^4$ | 1.9 x $10^6$ | 1.93 | 0.1 | 4.64 | 0.39 | (256, 512, 512) | 0.27 | 565 | 53.2 | 33.3 |
| 1.4 x $10^4$ | 1.6 x $10^4$ | 1.01 | 1.1 x $10^4$ | 1.2 x $10^6$ | 1.72 | 0.3 | 1.47 | 1.90 | (64, 128, 128) | 0.124 | 11.3 | 5.4 | 1.8 |
| 1.1 x $10^4$ | 1.2 x $10^4$ | 2.29 | 7.8 x $10^3$ | 8.6 x $10^6$ | 3.93 | 0.3 | 14.7 | 0.65 | (192, 384, 384) | 0.808 | 529 | 83.5 | 27.3 |
| 5.5 x $10^3$ | 5.5 x $10^4$ | 1.65 | 4.0 x $10^4$ | 4.0 x $10^6$ | 2.82 | 1.0 | 1.47 | 4.84 | (64, 128, 128) | 0.303 | 3.6 | 4.4 | 1.5 |
| 2.8 x $10^4$ | 2.8 x $10^4$ | 6.39 | 2.0 x $10^4$ | 2.0 x $10^6$ | 10.93 | 1.0 | 100 | 0.82 | (512, 512, 512) | 3.557 | 1099 | 220 | 46.6 |

Note. Input parameters and output parameters for select simulations are shown. For each of the eight paths in Figure 1(b), we show information for the lowest and highest (Ra, Ta) point on that path. The first six rows show information for constant $R_o$, paths, the next six for constant $R_a$, paths, and the last four for constant $S$ paths. We show the input Ra, Ta, and $R_o$, at the top of the atmosphere, as well as their stratification-weighted values at the midplane of the atmosphere, which provide a more direct comparison to Boussinesq values (Unno et al. 1960). We also provide the input $R_o$, at the top of the atmosphere, $S$, aspect ratio ($L_x/L_z$), and coefficient resolution ($\Delta x, \Delta y, \Delta z$). Each dimension of the physical grid is $3/2$ the size of the coefficient grid for adequate decalining of quadratic non-linear terms. Output values of $R_o$, $R_{e, o}$, $R_{e, z}$, and Nu are also provided. This table is published in its entirety in the machine-readable form and deposited in Zenodo (Anders et al. 2019).

(This table is available in its entirety in machine-readable form.)

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