Breakdown of large-scale circulation in turbulent rotating convection

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Abstract. - Turbulent rotating convection in a cylinder is investigated both numerically and experimentally at Rayleigh number $Ra = 10^9$ and Prandtl number $\sigma = 6.4$. In this Letter we discuss two topics: the breakdown under rotation of the domain-filling large-scale circulation (LSC) typical for confined convection, and the convective heat transfer through the fluid layer, expressed by the Nusselt number. The presence of the LSC is addressed for several rotation rates. For Rossby numbers $Ro \lesssim 1.2$ no LSC is found (the Rossby number indicates relative importance of buoyancy over rotation, hence small $Ro$ indicates strong rotation). For larger Rossby numbers a precession of the LSC in anticyclonic direction (counter to the background rotation) is observed. It is shown that the heat transfer has a maximal value close to $Ro = 0.18$ being about 15% larger than in the non-rotating case $Ro = \infty$. Since the LSC is no longer present at this Rossby value we conclude that the peak heat transfer is independent of the LSC.

Introduction. – The flow in a fluid layer confined between two horizontal plates and driven by a destabilising temperature gradient is commonly known as Rayleigh–Bénard convection [1]. The dimensionless parameters that govern this flow are the Rayleigh and Prandtl numbers, defined respectively as

$$Ra \equiv \frac{g\alpha\Delta T H^3}{\nu\kappa}, \quad \sigma \equiv \frac{\nu}{\kappa}.$$  

(1)

Here $g$ is the gravitational acceleration, $\Delta T$ the temperature difference between the plates, $H$ their vertical separation, and $\nu$, $\kappa$ and $\alpha$ the kinematic viscosity, thermal diffusivity and thermal expansion coefficient of the fluid, respectively. In most experimental investigations of Rayleigh–Bénard convection a cylindrical geometry is used. The cylinder axis is aligned vertically with gravity. This introduces an additional parameter, the diameter-to-height aspect ratio $\Gamma \equiv D/H$. In this work we will use $\Gamma = 1$ geometries. An interesting addition to this flow problem is the presence of a background rotation. In this case the rotation vector points antiparallel to gravity. The rotation rate $\Omega$ is incorporated in the dimensionless Rossby number $Ro$, which states the relative importance of buoyancy and rotation:

$$Ro \equiv \frac{1}{2\Omega \sqrt{\frac{g\alpha\Delta T}{H}}}.$$  

(2)

Thus, without rotation $Ro = \infty$ while rotation dominates for $Ro \ll 1$. Rotating convection is a simple model system relevant for a wide variety of geophysical and astrophysical flows, e.g., open-ocean convection in the polar regions [2], the outer convective layer of our Sun [3] and the interior of the giant gaseous planets [4]. Hence, rotating convection is a relevant topic for meteorologists, climatologists, oceanographers and astrophysicists alike.

It is known from linear stability analysis [1] that rotation stabilises convection, i.e., the critical Rayleigh number $Ra_c$ for onset of convective motion increases when rotation is applied. Paradoxically, the convective heat transfer through the fluid can actually increase by the addition of a moderate rotation. The heat transfer is expressed as the Nusselt number $Nu$:

$$Nu \equiv \frac{qH}{k\Delta T}.$$  

(3)
where the total heat-current density \( q \) is normalised by the conductive flux \( k \Delta T/H \) \((k\) is the thermal conductivity\) that would be present in absence of fluid motion. Increased \( \text{Nu} \) under rotation was reported in several experimental and numerical studies [5–9]. Generally, the effect is ascribed to an additional vertical transport term introduced in the boundary layers (BLs) by the rotation: the so-called Ekman pumping [10]. At higher rotation rates, however, the stabilising effect becomes dominant and a rapid decrease to the conductive state \( \text{Nu} = 1 \) is found, with all fluid motion suppressed.

Turbulent convection in a cylinder possesses a remarkable feature: the organisation of the flow into a domain-filling large-scale circulation (LSC), first described by Krishnamurti and Howard [11]. The warm plumes gather near the sidewall to rise on one side of the container, while the cold plumes sink on the opposite side. In this way one large ‘flywheel’ structure is formed [12]. The LSC is a starting point of the heat transfer scaling theory by Grossmann and Lohse [13–16]. The LSC is known to describe interesting temporal dynamics, such as a coherent oscillation in the azimuthal direction with a time scale of order one minute, and sudden reversals and cessations happening irregularly with time scales of many hours or even days (characteristic time scales based on \( \text{Ra} \times 10^3 \)). For an extensive overview we refer to Xi et al. [17] and references therein. The LSC is still present when rotation is added [18]. But, given the restraints that rotation causes, the LSC must be expected to break down at a certain rotation rate. The flow phenomenology is then best described as a collection of columnar vortices that concentrate the vertical transport in their interior.

In this work we investigate the flow phenomenology with experiments using stereoscopic particle image velocimetry (SPIV) and with direct numerical simulation (DNS). The long time scales associated with the azimuthal dynamics of the LSC are easier resolved in an experiment than in DNS. However, DNS is well-suited for calculation of the Nusselt number. A systematic variation of the rotation rate shows its influence on the heat transfer and the LSC. We wish to address whether the LSC plays a role in the convective heat transfer under rotation.

**Experimental procedure.** — The experimental setup is the same as in our previous works [19, 20]; here we briefly repeat the most important parts. The convection cell is schematically depicted in fig. 1(a). A cylinder of dimensions \( H = D = 230 \text{ mm} \) is closed from below by a copper block with electric heater underneath. At the top, cooling water is circulated through a transparent chamber. Bottom and top temperatures are controlled with temperature sensors and controllers; the applied temperature difference is \( \Delta T = 5 \text{ K} \) which corresponds to \( \text{Ra} = 1.11 \times 10^9 \) (at \( \sigma = 6.37 \), based on the mean temperature of 24°C). The temperature at the top is constant up to \( \pm 0.04 \text{ K} \); at the bottom up to \( \pm 0.02 \text{ K} \). The cylinder is filled with water that is seeded with 50-\( \mu \text{m} \)-diameter polyamid seeding particles. It is enclosed in a square container, with the volume in between also filled with water. This allows the crossing of a horizontal laser light sheet at 45 mm from the top without too much refraction. Two cameras at different viewpoints record the particle images. A stereoscopic particle image velocimetry (SPIV) algorithm [21] processes the images into three-component two-dimensional velocity vector fields consisting of 53 × 55 vectors. The rectangular measurement area covers roughly 120 × 150 mm\(^2\), thus not the full circular cross-section of the cylinder. The typical measurement error is about 5%. All equipment is placed on a rotating table. Experiments are conducted at different rotation rates; see table 1. For the cases \( \text{Ro} \geq 3.85 \) one velocity field per second is recorded for a measurement duration of more than one hour each (approximately 4000 fields). At smaller \( \text{Ro} \leq 2.89 \) fifteen velocity fields are measured per second, with a duration of over 11 minutes (approximately \( 10^4 \) fields) per experiment. With this distinction we attempt to capture the relevant time scales in each regime. Inspired by Xi et al. [17], fig. 1(b) illustrates our definition of the LSC orientation \( \theta \): the orien-
tation of the mean horizontal velocity vector in a circular region. Since the measurement area is near the top plate the horizontal branch of the LSC is detected; its orientation $\theta$ is thus conveniently found in this way. The spatial separation of upward (red) and downward (blue) motion is another indication of the LSC. A velocity snapshot at $Ro = 0.72$ in fig. 1(c) shows that the LSC is not present at that $Ro$.

An unfortunate side effect of the requirement for optical accessibility is the inability to measure $Nu$: due to lack of thermal insulation heat is lost and the heat input into the fluid cannot be measured accurately.

**Measurements of the LSC dynamics.** — The LSC orientation $\theta$ is monitored in time, see fig. 2(a). The initial orientations have been subtracted and the different curves are shifted apart vertically for clarity. Several interesting features can be observed. First, in the non-rotating case ($Ro = \infty$) there is on average no change in $\theta$. It oscillates rather coherently. This is a well-known feature of the LSC which we documented for this particular experiment in [19]. The oscillation period is $\tau_0 = 133 \pm 2$ s, which can, according to Brown et al. [22] and Xi et al. [17], be used to define a Reynolds number $Re_0 \equiv 2H^2/\nu \tau_0 = 862 \pm 13$ for the LSC. Second, when rotation is added the LSC describes an anticyclonic (counter to the background rotation) precession in the co-rotating reference frame (this was also observed by Hart et al. [18]). There is still an oscillation noticeable from the signals, but it is found to be more erratic than at $Ro = \infty$. Third, at $Ro = 1.44$ the last coherent trace of the LSC is detected, while at $Ro = 0.72$ and smaller only a strongly fluctuating signal for $\theta$ in time is found (not shown). This points at the disappearance of the LSC for these $Ro \lesssim 0.72$, which is confirmed by a visual inspection. The flow field is then dominated by vortical structures [6,8].

In fig. 2(b) the anticyclonic precession frequency $|\omega|$ and its dependence on $Ro$, on a log-log scale for the current study (red circles; $Ra = 1.11 \times 10^9$, $\sigma = 6.37$) and for Hart et al. [18] (black crosses; $Ra = 2.9 \times 10^{11}$, $\sigma = 8.4$). The error bars in the red symbols are based on the rms spread of $\theta(t)$ around the linear fit. The following lines are also included (red for this work and black for Hart et al. [18]): dash-dotted lines based on expression (4); solid lines based on (5); vertical dashed lines indicating $Ro$ at which the BL thickness $\delta_{BL}$ is equal to the theoretical Ekman layer thickness $\delta_E = \sqrt{\nu/\Omega}$.  

### Table 1: Rotation rates $\Omega$ and corresponding Rossby numbers $Ro$ used in the experiments. For reference, Earth’s rotation $\Omega = 7.3 \times 10^{-5}$ rad/s amounts to $Ro = 1.6 \times 10^3$ in the current setup.

<table>
<thead>
<tr>
<th>$\Omega$ (rad/s)</th>
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<tr>
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<td>11.5</td>
<td>0.080</td>
<td>1.44</td>
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<td>7.69</td>
<td>0.16</td>
<td>0.721</td>
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<td>0.64</td>
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<tr>
<td>0.030</td>
<td>3.85</td>
<td>1.28</td>
<td>0.090</td>
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### Fig. 2: (a) LSC orientation $\theta$ versus time $t$ for several $Ro$. (b) Absolute precession frequency $|\omega|$ and its dependence on $Ro$, on a log-log scale for the current study (red circles; $Ra = 1.11 \times 10^9$, $\sigma = 6.37$) and for Hart et al. [18] (black crosses; $Ra = 2.9 \times 10^{11}$, $\sigma = 8.4$). The error bars in the red symbols are based on the rms spread of $\theta(t)$ around the linear fit. The following lines are also included (red for this work and black for Hart et al. [18]): dash-dotted lines based on expression (4); solid lines based on (5); vertical dashed lines indicating $Ro$ at which the BL thickness $\delta_{BL}$ is equal to the theoretical Ekman layer thickness $\delta_E = \sqrt{\nu/\Omega}$.  

In fig. 2(b) the anticyclonic precession frequency $|\omega| = d\theta/dt$, obtained with linear fits to the data of fig. 2(a), is plotted as a function of $Ro$ (red circles). The results of Hart et al. [18] for $Ra = 2.9 \times 10^{11}$ and $\sigma = 8.4$ are also included (black crosses). Simple models to explain this dependence have been proposed [18, 23]. Here we follow the model of Brown and Ahlers (BA) [23]. Define co-rotating coordinates $(x, y, z)$ with the origin at the centre of the cylinder and the unit vector $\hat{z}$ pointing vertically upward. BA suggest a modelling of the LSC in terms of a prescribed velocity distribution: within a conical region inside the cylinder (indicated in grey in fig. 3) a horizontal velocity $u_x = 2\Omega z/H$, $u_y = 0$ is assumed, closed by a matching vertical velocity outside the conical region. Here $U$ is a typical velocity for the LSC and $r \equiv \sqrt{x^2 + y^2}$. The horizontal velocity causes a Coriolis acceleration $a_C = -2(\Omega + \omega)u_x$, directed to the right of $u_x$, with $\omega$ the precession rate of the LSC. The Coriolis acceleration is balanced by friction in the viscous BLs near the bottom and top plates and the sidewall, due to the precession of the LSC. BA use as estimate for the thickness of the BL $\delta_{BL} \equiv 0.5\nu H \nu^{1/2}$. This relation is independent of rotation and thus need not hold for smaller $Ro \lesssim 1$ [24], but for $Ro \gg 1$ it is applicable. The viscous term, act-
Also, we refine the friction near \( \sigma \) formed there.

With both contributions are averaged over the cylinder volume. With \( U \) based on the Reynolds number \( \text{Re} = U H / \nu \), BA find for \( \omega \):

\[
\frac{\omega}{\Omega} = \frac{1}{1 + 12 \text{Re}^{-1/2}}.
\]

This dependence is plotted in fig. 2(b) with the dash-dotted lines. (Hart et al. did not report the Reynolds number: \( \text{Re} = 1.0 \times 10^4 \) is found with the Grossmann–Lohse theory [13–16].)

We propose some refinements to the BA model. First, previous measurements of the velocity field due to the LSC have shown that the LSC has a rounded shape in a vertical cross-section, see, e.g., the sketches in [25]. In order to better mimic this shape we propose for the horizontal velocity in the conical region (marked grey in fig. 3):

\[ u_\phi = 2 U z [1 - (r/H)^2] / H. \]

Also, we refine the friction near the bottom and top plates. Since \( u_\phi \) is dependent on the radial coordinate the friction needs to be integrated radially. Applying these refinements in the model, we find for \( \omega \):

\[
\frac{\omega}{\Omega} = \frac{1}{1 + \frac{2 \pi}{\delta_z} \text{Re}^{-1/2}}.
\]

This relation is depicted in fig. 2(b) with the solid lines. The difference between the two formulations (4) and (5) is only minor. The exact velocity profile is found to be relatively unimportant. The nonlinear dependence of \( \omega \) on \( \text{Re} \) is not captured with these simple models. We expect that there is additional friction, possibly due to a sustained azimuthal oscillation at large \( \text{Re} \). A more sophisticated formulation of the friction term is required to replicate the dependence more accurately. The large discrepancy between model and the Hart et al. results is mostly due to an overestimated \( \text{Re} \); for high \( \text{Re} \) such as that of Hart et al. (\( \text{Re} = 2.9 \times 10^{11} \)) the Grossmann–Lohse theory shows significant deviations from recent experiments [22].

The breakdown of the LSC is at a \( \text{Re} \) close to where the BL structure to the rotation-dominated Ekman layer is a possible mechanism leading to the breakdown of the LSC.

With numerical simulations we can describe the transition between LSC and vortex state in more detail; a quantitative criterion for the existence of the LSC will be introduced.

**Numerical arrangement.** – The governing equations for the current flow problem are the incompressible Navier–Stokes and temperature equations with rotation and in the Boussinesq approximation [1]:

\[
\frac{\partial u}{\partial t} + (u \cdot \nabla)u + \frac{1}{
\text{Ro}} \frac{\partial T}{\partial z} = -\nabla p + \sqrt{\frac{\sigma}{\text{Ra}}} \nabla^2 u + T \hat{z},
\]

\[
\frac{\partial T}{\partial t} + (u \cdot \nabla)T = \frac{1}{\sqrt{\text{Ra}}} \nabla^2 T,
\]

\[ \nabla \cdot u = 0, \]

with \( u \) the velocity vector, \( p \) the pressure, \( T \) the temperature and \( \hat{z} \) the vertical unit vector pointing counter to gravity. These equations are shed into this dimensionless form by using the scaling variables \( H \) for length, \( \Delta T \) for temperature and \( \tau = H / U \) for time, based on the so-called free-fall velocity \( U \equiv \sqrt{g \Delta T H} \) [26]. The equations, written in cylindrical coordinates \((r, \phi, z)\) with corresponding velocity vector \((u_r, u_\phi, u_z)\), are solved in a cylindrical volume with \( \Gamma = 1 \). All walls are no-slip. The sidewall is adiabatic, while the top and bottom walls are kept at constant temperatures: \( T = 1 \) at \( z = -0.5 \) (bottom wall) and \( T = 0 \) at \( z = 0.5 \) (top wall).

The governing equations are discretised with second-order accurate finite-difference approximations. Time-integration is done with a third-order Runge–Kutta scheme. The spatial and temporal discretisation is described in more detail by Verzicco and Orlandi [27] and Verzicco and Camussi [26, 28].

All simulations adopt \( \text{Ra} = 1 \times 10^9 \) and \( \sigma = 6.4 \). Different Rossby numbers are used: \( \text{Ro} = 0.045, 0.068, 0.090, 0.18, 0.36, 0.72, 1.08, 1.44, 1.80, 2.16, 2.52, 2.88, 5.76, 11.52 \). The grid resolution is \( N_r \times N_\phi \times N_z = 385 \times 193 \times 385 \). The points in the azimuthal direction are distributed evenly. Close to the bottom and top walls, as well as close to the sidewall, there is a denser grid in order to resolve the BLs formed there.

An important result of the simulations for this work is the Nusselt number \( \text{Nu} \). It is calculated in three ways [28]: (i,ii) the time-and-area averaged wall-normal temperature gradient \( \text{Nu} = \langle \partial T / \partial z \rangle_{A,t} \) at top and bottom, and (iii) the time-and-volume averaged convective heat flux \( \text{Nu} = 1 + \sqrt{\text{Ra}} \langle u_z T \rangle_{V,t} \). The final \( \text{Nu} \) value is the weighted mean of the three results. Consistency of all three indicates that both bulk and near-wall resolution are adequate [28]. Furthermore, a set of 64 numerical probes is evenly distributed over a circle close to the sidewall \((r = 0.45)\) at half-height \((z = 0)\). These probes register local velocities, temperature, etc. at every time step.
This arrangement will be used to detect the LSC just as BA do [23]; the vertical-velocity signal as a function of \( \phi \) has a sin-like shape.

The duration of the averaging in each simulation was at least 300 dimensionless time units. This compares roughly to 150 large-eddy turnover times based on a fluid particle traversing an elliptic trajectory of circumference \( \sim 2H \) (cf. the LSC) at a velocity \( U \) [28].

LSC detection in simulations. – We now wish to address the existence of an LSC at the different Ro. The numerical probe data are used to formulate a quantitative criterion. At each time step the probes record an azimuthal profile of vertical velocity \( u_z(\phi) \). A Fourier transform is applied to the \( u_z(\phi) \) profile to determine the relative importance of each azimuthal mode. The \( n = 1 \) Fourier mode can be associated with the LSC, as is shown in fig. 4(a). In this plot the circles represent \( u_z(\phi) \) from the numerical probes for a single time, while the solid line is drawn by using the amplitude and phase for the \( n = 1 \) mode as obtained with the Fourier transform. To quantify whether an LSC exists, the mean energy of each Fourier mode is calculated for each Ro value. Our criterion is that there is an LSC when the \( n = 1 \) mode is the most energetic individual Fourier mode. The results are shown in fig. 4(b) as mean energy content of mode \( n \) (\( E_n \)) normalised by the mean total energy \( E \). They are \( n = 0 \), the constant-value contribution (crosses), \( n = 1 \), the signature of the LSC (circles), \( n = 2 \) (squares), \( n = 3 \) (triangles) and the sum of modes \( n \geq 4 \) (pluses). A rough division into three Ro ranges is observed. For \( Ro > 2.5 \) the LSC mode \( n = 1 \) contains more than half of the energy, a clear indication of the presence of an LSC. For \( Ro \lesssim 1.2 \) the \( n = 1 \) mode is buried under higher-order modes; no LSC is present. The intermediate transition range \( 1.2 \lesssim Ro \lesssim 2.5 \) shows dramatic changes in \( E_1 \); it rapidly decreases with decreasing \( Ro \). Yet the \( n = 1 \) mode is still the strongest individual mode. We still consider an LSC to be present, but it is weaker than for \( Ro > 2.5 \).

The conclusion that an LSC is present for \( Ro \gtrsim 1.2 \) and absent for \( Ro \lesssim 1.2 \) matches well with the experimental results: in the experiment at \( Ro = 1.44 \) the LSC was detected, while it was not found at \( Ro = 0.72 \).

The simulation method described here is well-suited for the detection of an LSC. Its azimuthal dynamics, however, possesses long time scales that are not easily resolved in simulations. An example: the azimuthal oscillation period \( \tau_0 = 133 \pm 2 \) s of the LSC covers about 30 dimensionless time units \( \tau \). Our current simulations only cover the time required for ten LSC oscillations. The addition of a small rotation adds an even larger time scale. Therefore, it is impractical to study the azimuthal dynamics of the LSC with these simulations.

Heat transfer. – DNS is eminently suited for calculation of the heat transfer. In fig. 5 the Nusselt number \( Nu \) is presented as a function of \( Ro \) and scaled with its corresponding value at \( Ro = \infty \) (no rotation). This visualisation shows the relative increase of \( Nu \) under rotation, and makes it easier to include results from other studies that may not match exactly in terms of Rayleigh and Prandtl numbers, or domain aspect ratio. The current results are depicted with squares. Other studies that are included: (i) Rossby [5] (dash-dotted line), experiments in water at \( Ra = 2.5 \times 10^6 \), \( \sigma = 6.8 \), with variable aspect ratio, but always \( \Gamma > 6 \); (ii) Liu and Ecke [7] (triangles), experiments in water at \( Ra = 1 \times 10^8 \) (extrapolated from their correlations), \( \sigma = 6.3 \), cuboid cell with side-to-height aspect ratio \( \Gamma = 0.78 \); (iii) Oresta et al. [9] (circles), DNS at \( Ra = 2 \times 10^4 \), \( \sigma = 0.7 \) in a cylinder with \( \Gamma = 0.5 \). Indeed, in all included results a rise of \( Nu \) is found in a range of \( Ro \). But the Rossby numbers for the peak Nusselt values from the various studies disagree; there are dependencies on all parameters \( Ra \), \( \sigma \) and even \( \Gamma \), plus other possible influences like the geometry (cylinder vs. cuboid).

The increased \( Nu \) is found for \( 0.05 \lesssim Ro \lesssim 2.5 \). The upper limit of this range is the highest Rossby number at which rotation affects the LSC. The maximal heat transfer is found at \( Ro \approx 0.2 \), for which definitely no LSC is present. Thus the LSC takes no part in the peak heat

Fig. 4: (a) Azimuthal profile \( u_z(\phi) \) (circles) and the corresponding \( n = 1 \) Fourier mode (solid line), as an indication of the LSC, at \( Ro = 5.76 \). (b) Normalised average energy content \( E_n/E \) of Fourier modes: \( n = 0 \) (crosses), \( n = 1 \) (circles), \( n = 2 \) (squares), \( n = 3 \) (triangles), and the sum of modes \( n \geq 4 \) (pluses). The vertical dashed lines indicate \( Ro = 1.2 \) and \( Ro = 2.5 \).
The dominant flow structuring and heat transfer under rotation. At $\text{Ro} = 0.045$ the Nusselt number is still 95% of the non-rotating value, so this situation is still quite far from the conductive state with $\text{Nu} = 1$.

**Conclusion.** — The dominant flow structuring and heat transfer in rotating turbulent convection in a cylinder have been investigated numerically and experimentally. At Rossby numbers $\text{Ro} \gtrsim 1.2$ the domain-filling large-scale circulation (LSC) is observed. The azimuthal dynamics of the LSC in time are described: without rotation the LSC exhibits an oscillation with a well-defined frequency, while under rotation the LSC describes an anticyclonic precession with an intricate dependence on $\text{Ro}$. For $\text{Ro} \lesssim 1.2$ it is seen in both simulation and experiment that the LSC is no longer the dominant flow feature; it is fully suppressed by the rotation. In a future paper we will use the DNS results to investigate the BLs in detail. Assumptions in the Brown & Ahlers model for the LSC precession [23] can then be tested directly.

The heat transfer through the fluid layer as a function of $\text{Ro}$ attains a maximal value close to $\text{Ro} = 0.18$. In spite of the LSC being a crucial ingredient of heat transfer scaling theory in non-rotating turbulent convection [13–16], it is found that under rotation the LSC can play only a minor role in the increased heat transfer: the LSC is damped by rotation when $1.2 \lesssim \text{Ro} \lesssim 2.5$ and fully suppressed when $\text{Ro} \lesssim 1.2$. Therefore it is anticipated that a theory for the heat transfer under strong rotation needs another starting point; it cannot be based on the existing theory for non-rotating convection [13–16].

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