Finite-Size Effects Lead to Supercritical Bifurcations in Turbulent Rotating Rayleigh-Bénard Convection

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(Received 19 July 2010; published 23 November 2010)

In turbulent thermal convection in cylindrical samples with an aspect ratio \( \Gamma = D/L \) \((D \text{ is the diameter and } L \text{ the height})\), the Nusselt number \( \text{Nu} \) is enhanced when the sample is rotated about its vertical axis because of the formation of Ekman vortices that extract additional fluid out of thermal boundary layers at the top and bottom. We show from experiments and direct numerical simulations that the enhancement occurs only above a bifurcation point at a critical inverse Rossby number \( 1/\text{Ro}_c \), with \( 1/\text{Ro}_c \approx 1/\Gamma \). We present a Ginzburg-Landau–like model that explains the existence of a bifurcation at finite \( 1/\text{Ro} \), as a finite-size effect. The model yields the proportionality between \( 1/\text{Ro} \) and \( 1/\Gamma \) and is consistent with several other measured or computed system properties.

Turbulence, by virtue of its vigorous fluctuations, is expected to sample all of phase space over wide parameter ranges. This viewpoint implies that there should not be any bifurcations between different turbulent states. Contrary to this, several cases of discontinuous transitions have been observed recently in turbulent systems [1]. When they occur, they are likely to be provoked either by changes in boundary conditions or boundary-layer structures or by discontinuous changes in the large-scale structures, as a parameter is varied.

Recently, some of us [2,3] reported on the effect of rotation about a vertical axis at a rate \( \Omega \) on turbulent convection in a fluid heated from below and cooled from above (known as Rayleigh-Bénard convection; for recent reviews, see [4–6]). For a cylindrical sample with an aspect ratio \( \Gamma = D/L = 1.00 \) \((D \text{ is the diameter and } L \text{ the height})\) a supercritical bifurcation was found, both from experiments and from direct numerical simulation (DNS) of the Boussinesq equations of motion. At a finite \( \Omega \), as expressed by the inverse Rossby number \( 1/\text{Ro} \approx \Omega \) (to be defined explicitly below), there was a sharp transition from a state of nearly rotation-independent heat transport (as expressed by the Nusselt number \( \text{Nu} \) to be defined explicitly below) to one in which \( \text{Nu} \) was enhanced by an amount \( \Delta \text{Nu}(1/\text{Ro}) \). This is illustrated by the data shown in Fig. 1. The increase of \( \text{Nu} \) was attributed to Ekman pumping [2,3,7–13], i.e., to the formation of (cyclonic) vertical vortex tubes (“Ekman vortices”), which extract and vertically transport additional fluid from the boundary layers (BLs) and thereby enhance the heat transport. The bifurcation was located at a critical value \( 1/\text{Ro}_c \approx 0.40 \) [3]. The reason for the existence of the bifurcation at \( 1/\text{Ro}_c > 0 \) hitherto had not been understood. While such bifurcations are common near the onset of Rayleigh-Bénard convection in the domain of pattern formation [14], their existence in the turbulent regime implies a paradigm shift.

In this Letter, we report on further experiments for samples with \( \Gamma = 2.00, 1.00, \) and 0.50 which (i) all show bifurcations between different turbulent states and (ii) reveal that \( 1/\text{Ro}_c \) varies approximately in proportion to \( 1/\Gamma \). We offer an explanation of these and other phenomena in terms of a phenomenological Ginzburg-Landau–like description which predicts a finite-size effect upon the vortex density \( A \).

We assumed that the relative Nusselt enhancement \( \delta \text{Nu}(1/\text{Ro})/\text{Nu}(0) \) is proportional to the average \( A \) of \( A \) over a horizontal cross section of the sample near the BLs. Consistent with the DNS that we report here, we assumed that \( A \) vanishes at the sample side wall. For the infinite system the model predicts that \( A \), and thus \( \delta \text{Nu}/\text{Nu}(0) \), increases linearly from zero starting at \( 1/\text{Ro} = 0 \). For the finite system the model gives a threshold shift proportional to \( 1/\Gamma \) as found in the experiment and by DNS. The shift is predicted to be followed by a linear increase of \( A \) in proportion to \( (1/\text{Ro}) - (1/\text{Ro}_c) \) which yields \( \delta \text{Nu}/\text{Nu}(0) = S_1(\Gamma)(1/\text{Ro} - 1/\text{Ro}_c) \). The model gives an initial slope \( S_1(\Gamma) \) that decrease with decreasing \( \Gamma \), again consistent with DNS and measurements. From DNS we show that \( A \) decreases to zero near the side wall over a length that is consistent with an estimate of a healing length \( \xi \) based on the model. Thus, we found consistency between the model predictions and all properties that we were able to either measure or compute from DNS.

Before proceeding, we define the relevant dimensionless parameters. The inverse Rossby number is given by

\[
\frac{1}{\text{Ro}} \approx \Omega \sqrt{\frac{\nu}{\alpha}} \frac{\delta T}{H} \frac{H}{L}
\]
1/\( \text{Ro} \) = \((2\Omega)/(\sqrt{\beta g \Delta T/L})\), where \( \Omega \) is the rotation rate in rad/s, \( \beta \) the isobaric thermal expansion coefficient, \( \Delta T \) the temperature difference between the bottom and top plate, and \( g \) the gravitational acceleration. The Rayleigh number is \( \text{Ra} = (\beta g \Delta TL^3)/(\nu \kappa) \), where \( \kappa \) and \( \nu \) are the thermal diffusivity and the kinematic viscosity, respectively. The Nusselt number is given by \( \text{Nu} = (QL)/(\Delta T \lambda) \), where \( Q \) is the heat-current density and \( \lambda \) is the thermal conductivity. Finally, the Prandtl number is \( \text{Pr} = \nu/\kappa \).

In Fig. 1, we show experimental and numerical data [15] for \( \text{Nu}(1/\text{Ro})/\text{Nu}(0) \) as a function of \( 1/\text{Ro} \) for several values of \( \text{Ra} \). From top to bottom, the three panels are for \( \Gamma = 0.50 \), 1.00, and 2.00, respectively [17]. One sees that there are substantial variations even below the bifurcation, particularly at the larger \( \text{Ra} \). To our knowledge the origin of this structure is not known in detail. One sees that there are clear breaks in the curves, e.g., for \( \Gamma = 1.00 \) [Fig. 1(b)] at \( 1/\text{Ro} = 0.4 \), indicating the bifurcation to a different state. The location of this transition is well within our resolution independent of \( \text{Ra} \).

In Fig. 2, we plotted all the available data for \( 1/\text{Ro}_c \) for \( \text{Pr} = 4.38 \) (and different \( \text{Ra} \)) as a function of \( 1/\Gamma \). The line shown there is a fit of

\[
\frac{1}{\text{Ro}_c} = a \left( \frac{1}{\Gamma} \right) + b
\]

(1)

to the data. Its coefficients are \( a = 0.381 \) and \( b = 0.061 \). One sees that the data are consistent with an initial linear increase from zero of \( 1/\text{Ro}_c \) with \( 1/\Gamma \), with a small quadratic contribution becoming noticeable as \( 1/\Gamma \) becomes larger.

In order to understand the \( \Gamma \) dependence of \( 1/\text{Ro}_c \), we studied the vortex statistics by using data obtained from DNS. We used the so-called \( Q \) criterion [9,18–20] to determine the fraction \( \hat{A} \) of the horizontal area that was covered by vortices. Using this criterion implies that the quantity \( Q_{2D} \) [3,21], which is a quadratic form of various velocity gradients, was calculated in a plane of fixed height. An area is then identified as a “vortex” when \( Q_{2D} < -0.1 \langle |Q_{2D}| \rangle \), where \( \langle |Q_{2D}| \rangle \) is the volume-averaged value of the absolute values of \( Q_{2D} \) [3]. The result of this procedure is shown for different \( 1/\text{Ro} \) in Fig. 3 for \( \text{Pr} = 6.26 \). In Fig. 4, we plot \( \hat{A} \) as a function of \( 1/\text{Ro} \) at the edge of the kinetic BL (which depends on \( \text{Ro} \); see [22]) and at the fixed distance 0.023L (the kinetic BL thickness without rotation) from the plates. Although there is quite a bit of scatter, the data are consistent with a linear increase of \( \hat{A} \) for \( 1/\text{Ro} > 1/\text{Ro}_c \), with a small constant background \( \hat{A} = A_0 \) below \( 1/\text{Ro}_c \). The azimuthally averaged vortex density \( \langle \hat{A} \rangle_\phi \) is given in Fig. 5. It shows that the Ekman vortices are inhomogeneously distributed: While in the bulk their fraction is roughly constant, there are almost
phenomenological Ginzburg-Landau–like model for the

distance with vortices at the edge of the kinetic BL (circles) and at a

no vortices at all close to the side wall, signaling a strong boundary effect.

In an effort to understand the existence of a finite onset (see Fig. 1) of the Ekman-vortex formation and the dependence of the critical inverse Rossby number on $\Gamma$ (see Fig. 2), to elucidate the linear rise and initial slope of $\langle A \rangle_\phi(r/L)$ near the wall (see Fig. 5), we propose a phenomenological Ginzburg-Landau–like model for the local vortex density:

$$
\dot{A} = (1/\text{Ro}^2)A - gA^3 + \xi_0^2 \nabla^2 A. \tag{2}
$$

Here $\dot{A}$ is the time derivative of $A$. We chose the coefficient of the linear term as $1/\text{Ro}^2$ because for the time-independent infinitely extended spatially uniform system it yields a stable solution $A = g^{-1/2}(1/\text{Ro})$, which implies a vortex density proportional to the rotation rate. The term with $\nabla^2 A$ represents the lowest-order term of a gradient expansion since terms proportional to $\nabla A$ would lead to an unphysical propagating mode.

When spatial variations are allowed, the ground state $A = 0$ can be shown to be stable (i.e., to have a growth rate $\sigma < 0$) to disturbances with wave number $k$ when $1/\text{Ro}$ falls below a neutral curve given by

$$
1/\text{Ro}_c(k) = \xi_0 k. \tag{3}
$$

For the finite system it is necessary to introduce appropriate boundary conditions. Here we shall consider a one-dimensional system over the range $-\Gamma/2 \leq x \leq \Gamma/2$ for simplicity and illustrative purposes. The two-dimensional system with circular boundaries and no azimuthal variation was treated in detail in Ref. [23] and yields the same result for $1/\text{Ro}_c$. Since there can be no vortices at the side wall of the sample (see Fig. 5, where we verified this based on the numerical data), we chose $A(-\Gamma/2) = A(\Gamma/2) = 0$. For the wave number $k_0$ of the lowest mode this yields $k_0 = \pi/\Gamma$. This in turn gives

$$
1/\text{Ro}_c = 1/\text{Ro}_c(k_0) = \pi \xi_0/\Gamma. \tag{4}
$$

Thus, consistent with the data in Fig. 2, the model yields the proportionality between $1/\text{Ro}_c$ and $1/\Gamma$. We note that the curvature indicated by the quadratic contribution to Eq. (1) can be accommodated easily by higher-order gradient terms in Eq. (2). Comparison with experiment [see Eq. (1)] gives $\xi_0 = a/\pi = 0.121$.

To elucidate the rapid decrease of $\langle A \rangle_\phi(r/L)$ in Fig. 5 near $r/L = 0.5$, we consider Eq. (2) for a semi-infinite system over the range $-\infty < x \leq 0.5$ with the boundary condition $A(x = 0.5) = 0$. It yields the solution

$$
\langle A \rangle_\phi(r/L) = \langle A \rangle_\phi(r/L = 0.5) \left( \frac{r/L}{0.5} \right)^{-1/2}. \tag{5}
$$

In total, the statistics is based on 8 snapshots. The dashed line is the uniform case. The density approaches zero close to the side wall.
\[ A(x) = (Ro^2g)^{-1/2} \tanh[(0.5 - x)/\xi]. \] 

(5)

with

\[ \xi = \sqrt{2}\xi_0 Ro. \] 

(6)

Thus, near the boundaries, the model predicts that the amplitude \( A(x) \) of the one-dimensional model, and thus to a good approximation also the azimuthal average \( \langle A \rangle \rangle \) in Fig. 5, should “heal” to its bulk value over a length \( \xi \). Using a representative \( Ro \approx 0.4 \) for Fig. 5, we estimate from Eq. (6) that \( \xi \approx 0.07 \). This is roughly consistent with the rapid variation of \( \langle A \rangle \rangle \) near \( r/L = 0.5 \) seen in Fig. 5.

Above the bifurcation the model yields [23]

\[ \tilde{A} = g^{-1/2} \left( \frac{1}{Ro} - \frac{1}{Ro_c} \right). \]

(7)

Thus

\[ \frac{\delta Nu}{Nu(0)} = S_1(\Gamma) \left( \frac{1}{Ro} - \frac{1}{Ro_c} \right). \]

(8)

which is consistent with the data in Fig. 1. Numerical solutions of the amplitude equation of Ref. [23] have shown that the renormalized coefficient \( g \) in Eq. (7) is larger than \( g \) in Eq. (2). Thus, the initial slope of \( \tilde{A} \) above \( 1/Ro \), and \( S_1(\Gamma) \) in Eq. (8) are reduced by the finite size of the system. The decrease of \( S_1(\Gamma) \) with decreasing \( \Gamma \) that can be seen in Fig. 1 is also consistent with the model.

At constant \( \Gamma \), \( S_1 \) depends slightly on the Prandtl number. This suggests that the nonlinear coefficient \( g \) in Eq. (2), and thus \( \tilde{g} \) in Eq. (7), is dependent on \( Pr \). Similarly, the bifurcation point \( 1/Ro_c \) depends slightly on \( Pr \). This is accommodated in the model equation (2) by a slightly Prandtl-dependent length scale \( \xi_0 \).

It remains to be seen whether the phenomena reported and explained here in terms of a finite-size effect have analogies in bifurcations between turbulent states in other systems [1]. The more general lesson which is learned is that the Ginzburg-Landau approach, which has been so versatile to understand the spatiotemporal dynamics of patterns, can also be useful in understanding the remarkable bifurcations between turbulent states.

We thank Jim Overkamp for contributing to the experiments with \( \Gamma = 2 \) and Gerald Oerlemans, Chao Sun, and Freek van Uittert for the design and construction of the experimental setup in Eindhoven. The work of S. W., J.-Q. Z., and G.A. was supported by the U.S. National Science Foundation through Grant No. DMR07-02111. We thank the DEISA Consortium (www.deisa.eu), co-funded through the EU FP6 Project No. RI-031513 and the FP7 Project No. RI-222919, for support within the DEISA Extreme Computing Initiative. The simulations were performed on the Huygens cluster (SARA) and the support from Wim Rijks (SARA) is gratefully acknowledged. R. J. A. M. S. was financially supported by the Foundation for Fundamental Research on Matter (FOM).

[15] In the first citation in Ref. [2], we described the numerical method—a finite difference solver for the Boussinesq equations with the Coriolis force added to account for the rotation. For example, the simulations at \( Ra = 2.91 \times 10^8, Pr = 4.38, \) and \( \Gamma = 2 \) were performed on a \( 769 \times 385 \times 289 \) grid in the azimuthal, radial, and axial direction, respectively, which yielded sufficient resolution in the BLs and the bulk [16]. Special care was taken for the azimuthal and radial resolution in order to keep the flow well resolved close to the sidewall boundaries in this large box.
[17] Some of the data for \( \Gamma = 1.00 \) had been published [2, 3].