

Convection in a rapidly rotating spherical shell with an imposed laterally varying thermal boundary condition

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We investigate thermally driven convection in a rotating spherical shell subject to inhomogeneous heating on the outer boundary, extending previous results to more rapid rotation rates and larger amplitudes of the boundary heating. The analysis explores the conditions under which steady flows can be obtained, and the stability of these solutions, for two boundary heating modes: first, when the scale of the boundary heating corresponds to the most unstable mode of the homogeneous problem; second, when the scale is larger. In the former case stable steady solutions exhibit a two-layer flow pattern at moderate rotation rates, but at very rapid rotation rates no steady solutions exist. In the latter case, stable steady solutions are always possible, and unstable solutions show convection rolls that cluster into nests that are out of phase with the boundary anomalies and remain trapped for many thermal diffusion times.

Key words: absolute/convective < instability, convection, rotating flows < geophysical and geological flows

1. Introduction

The problem of thermally driven convection in a rapidly rotating spherical shell with homogeneous boundary conditions (Chandrasekhar 1961) has been extensively studied owing to its application to planetary cores. The governing non-dimensional parameters are defined in full in (2.4) below. They are the Rayleigh number R , measuring the strength of the applied temperature difference across the shell; the Ekman number E , measuring the rate of rotation and the Prandtl number Pr , the ratio of viscous and thermal diffusivities. Initial work on this problem, herein called the *homogeneous problem*, focused on the onset of convective instability (e.g. Roberts 1968; Busse 1970; Zhang 1992*b*; Jones, Soward & Mussa 2000; Dormy *et al.* 2004; Zhang & Liao 2004; Zhang, Liao & Busse 2007), which involves determining the most unstable azimuthal wavenumber, m_c , at which convective instability occurs, and the corresponding critical value of the Rayleigh number, $R = R_c$, as a function of E and Pr . Subsequent work has considered the sequence of bifurcations that occur as R is increased (e.g. Zhang 1991, 1992*a*; Sun, Schubert & Glatzmaier 1993; Tilgner & Busse 1997; Grote & Busse 2001).

An important modification to the homogeneous problem in the geophysical context is the effect of inhomogeneities at the outer boundary (OB) of the spherical shell.

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This problem, called the *inhomogeneous problem* here, may also be applicable to other bodies such as extrasolar planets. As well as E , R and Pr , the inhomogeneous problem is governed by a new non-dimensional parameter ϵ , which describes the ratio of peak-to-peak thermal amplitude variations on the OB to the temperature difference between the boundaries. The azimuthal scale of the boundary variations, denoted by the wavenumber m_b , is also significant.

Zhang & Gubbins (1993) investigated the inhomogeneous problem for low ϵ and high E . In the Earth E is very small. Estimating ϵ for the Earth is problematic due to difficulties in estimating the superadiabatic heat flux; it is unlikely to be small, with mantle convection studies predicting anything from $\epsilon = O(1)$ to $\epsilon = \infty$ (Nakagawa & Tackley 2007). Furthermore, seismic tomography shows large scale anomalies at the base of the mantle (Masters *et al.* 1996) making $m_b < m_c$ in the Earth. In this paper we extend the results of Zhang & Gubbins (1993) by investigating the inhomogeneous problem at lower E and higher ϵ for various scales of the boundary anomalies.

The inhomogeneous problem gives rise to two basic flows. One is convective, setting in as an instability of the basic state that is driven by a difference in temperature between the boundaries (Chandrasekhar 1961). This flow will be termed *buoyancy-driven* in this work. The other flow is a thermal wind maintained by the boundary variations, which will be called *boundary-driven*. While a stable equilibrium is possible for the convective flow alone (when $R < R_c$), there is no such state for a finite value of ϵ as the lateral variations drive a thermal wind (Zhang & Gubbins 1992). Thus no critical value of R exists for the inhomogeneous problem, where these two flows occur together.

The inhomogeneous problem was first studied in the infinite Pr limit by Zhang & Gubbins (1992), who simulated flows driven by the lateral variations. They observed a phase shift between the flow pattern and boundary anomalies due to a dominant thermal wind balance. Zhang & Gubbins (1993) studied the effect of an inhomogeneous fixed temperature outer boundary condition on convective instability at infinite Pr , high E and low ϵ . Using a boundary condition with $m_b = m_c$, which, following Zhang & Gubbins (1993) we call *primary convection* to emphasize the wavenumber matching, they found stable (in the linear sense) steady solutions, where the homogeneous flow pattern attached, or ‘locked’, to the boundary anomalies. In this paper the term ‘steady’ describes a time-independent solution, while the term ‘locked’ is used to describe flows that attach to the boundary anomalies: locked solutions may be unsteady. In the neighbourhood of R_c , Zhang & Gubbins (1993) observed a resonance, with kinetic energy reaching a peak at a value of R , R_r , just above critical. Stable steady solutions bifurcated to time-dependent solutions at a value of R just above R_r , with the type of bifurcation dictated by the value of ϵ . For low ϵ the bifurcation is of Hopf type, leading to an oscillatory solution. For larger values of ϵ the bifurcation is of saddle-node type, and sets in as a wave of infinite period. With $m_b < m_c$, called *secondary convection* here to emphasize the wavenumber mismatch, Zhang & Gubbins (1993) obtained stable steady solutions, albeit at higher values of ϵ . The flow patterns were dominated by the scale of the boundary forcing but modulated by the scale of the most unstable mode. Unstable solutions exhibited wavenumber vacillations between the competing m_b and m_c modes.

In this work a large value of Pr is used that makes inertial forces negligible, an important ingredient for obtaining locked solutions (Zhang & Gubbins 1996). It has been found by experiment that $Pr = 100$ gives a good approximation to the limit $Pr \rightarrow \infty$ considered by Zhang & Gubbins (1993) and this value is used throughout this paper. We demonstrate our findings using four values of the Ekman

number, $E = 8 \times 10^{-4}$, 6×10^{-4} , 3×10^{-4} and 1.2×10^{-4} . Supplementary calculations have been undertaken with E as low as 2×10^{-5} , however these larger values of E demonstrate new features more clearly. The free parameters are thus ϵ and R . We keep $R - R_c$ small to clearly demonstrate the effects of the inhomogeneous boundary condition.

Like Zhang & Gubbins (1993) we employ fixed temperature boundary conditions in the problem formulation. The alternative prescription is fixed heat flux at the OB, the relevant boundary condition for the geophysical application. With homogeneous heat flux boundary conditions Gibbons, Gubbins & Zhang (2007) showed that m_c is not a simple function of E (see also Takehiro *et al.* 2002), while fixed temperature boundary conditions yield a smooth monotonic dependence of m_c on E . In the moderately low E regime, which is the focus of this paper, the two boundary conditions yield very similar results (Gibbons *et al.* 2007). We therefore use fixed temperature boundary conditions to facilitate comparison with previous studies. We shall only consider equatorially symmetric boundary anomalies, where boundary temperature variations are mirrored in the equatorial plane, because the most unstable mode is always equatorially symmetric (Busse 1970) and both Zhang & Gubbins (1993) and Gibbons & Gubbins (2000) found that equatorially antisymmetric boundary conditions have little effect on the solutions compared to the corresponding symmetric case. The inhomogeneous boundary condition used will be time-independent throughout.

The mathematical and numerical formulation are given in §2. The analysis is conducted in four stages: (i) linear onset of convective instability is studied in §3 to determine $R_c(E)$ and $m_c(E)$, which guide choices for the scale of the boundary inhomogeneities; (ii) steady nonlinear inhomogeneous solutions for various ϵ in the neighbourhood of R_c using different scales of the boundary variations are found in §4; (iii) linear stability analysis of the steady solutions found in stage (ii) is presented in §5; (iv) time integration of the full nonlinear equations for solutions found to be linearly unstable are presented in §6. Conclusions are presented in §7.

2. Mathematical formulation

Consider a Boussinesq fluid of constant thermal diffusivity κ , constant coefficient of thermal expansion α and constant viscosity ν confined to a rapidly rotating spherical shell of radial extent $d = r_o - r_i$. Here r_i corresponds to the inner boundary (IB) and r_o to the OB. Spherical polar coordinates, (r, θ, ϕ) , will be used throughout. The fluid rotates about the vertical z -axis with angular velocity Ω . With no flow the basic temperature gradient, $\nabla T_0 = -\beta \mathbf{r}$, is produced by a uniform distribution of internal heat sources, where T_0 is the basic state temperature and β is the temperature gradient of the internally heated basic state. Scaling length by the shell thickness, d , time by the thermal diffusion time, d^2/κ , and using βd^2 as the unit of temperature, the non-dimensional equations are

$$\frac{E}{Pr} \left(\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} \right) + \mathbf{z} \times \mathbf{u} = -\nabla P + Ra T \mathbf{r} + E \nabla^2 \mathbf{u}, \tag{2.1}$$

$$\frac{\partial T}{\partial t} + (\mathbf{u} \cdot \nabla) T = \nabla^2 T + 3, \tag{2.2}$$

$$\nabla \cdot \mathbf{u} = 0. \tag{2.3}$$

Here \mathbf{u} is the fluid velocity, P is the pressure and T is the temperature. The aspect ratio, r_o/r_i , of the shell is set to 0.4 for comparison with Zhang & Gubbins (1993).

The Ekman number, Prandtl number and modified Rayleigh number are

$$E = \frac{\nu}{2\Omega d^2}, \quad Pr = \frac{\nu}{\kappa}, \quad Ra = \frac{g\alpha\beta d^3}{2\Omega\kappa} = RE, \quad (2.4)$$

where g is the gravitational acceleration at the outer boundary. Gravity varies linearly with radius.

The velocity field is decomposed into toroidal and poloidal components

$$\mathbf{u} = \nabla \times (\mathcal{T}\mathbf{r}) + \nabla \times \nabla \times (\mathcal{P}\mathbf{r}). \quad (2.5)$$

Substituting (2.5) into (2.1) and applying $\mathbf{r} \cdot \nabla \times$ and $\mathbf{r} \cdot \nabla \times \nabla \times$ gives

$$E \left(\nabla^2 - \frac{1}{Pr} \frac{\partial}{\partial t} \right) L_2 \mathcal{T} = \frac{E}{Pr} \mathbf{r} \cdot \nabla \times [(\mathbf{u} \cdot \nabla)\mathbf{u}] - \mathbf{r} \cdot \nabla \times (\mathbf{z} \times \mathbf{u}), \quad (2.6)$$

$$E \left(\nabla^2 - \frac{1}{Pr} \frac{\partial}{\partial t} \right) L_2 \nabla^2 \mathcal{P} + Ra L_2 T = -\frac{E}{Pr} \mathbf{r} \cdot \nabla \times \nabla \times [(\mathbf{u} \cdot \nabla)\mathbf{u}] - \mathbf{r} \cdot \nabla \times \nabla \times (\mathbf{z} \times \mathbf{u}), \quad (2.7)$$

where the operator L_2 is defined as

$$L_2 = - \left[\frac{1}{\sin\theta} \frac{\partial}{\partial\theta} \left(\sin\theta \frac{\partial}{\partial\theta} \right) + \frac{1}{\sin^2\theta} \frac{\partial^2}{\partial\phi^2} \right]. \quad (2.8)$$

To compare with Zhang & Gubbins (1993), impenetrable and stress-free boundary conditions are applied to \mathbf{u} at r_i and r_o , which become

$$\mathcal{P} = \frac{\partial^2 \mathcal{P}}{\partial r^2} = \left(\frac{\partial}{\partial r} - \frac{1}{r} \right) \mathcal{T} = 0, \quad (2.9)$$

upon using (2.5). We consider fixed temperature at the IB and OB, corresponding to

$$T(r_i) = 0, \quad T(r_o) = \epsilon f(r_o, \theta, \phi), \quad (2.10)$$

where $f(r_o, \theta, \phi)$ defines the pattern of temperature variations on the OB, $r = r_o$. The spherical harmonic expansion of $f(r_o, \theta, \phi)$ is normalized as

$$\int_0^{2\pi} \int_0^\pi f^2(r_o, \theta, \phi) \sin\theta \, d\theta \, d\phi = 1. \quad (2.11)$$

In this paper f is expressed by a single spherical harmonic, written Y_l^m . Only equatorially symmetric boundary variations are considered and so $l = m$ in all the cases that follow. The parameter ϵ in (2.10), the inhomogeneity parameter, is a measure of the amplitude of boundary variations and is defined as

$$\epsilon = \frac{T_B}{\beta d^2}, \quad (2.12)$$

where T_B is a measure of the peak-to-peak temperature variations on the boundary.

Equations (2.2), (2.6) and (2.7), together with the boundary conditions (2.9) and (2.10), form the inhomogeneous problem. The manifold of solutions is explored in four stages by the following sets of calculations.

2.1. Onset of convection with homogeneous boundary conditions

For the linear onset of convective instability with homogeneous boundaries ($\epsilon = 0$) all nonlinear and inhomogeneous terms in (2.6) and (2.7) are neglected. An exponential time-dependence, $\partial/\partial t = \sigma$, is assumed, where $\sigma = \sigma_r + i\sigma_i$ is the complex growth rate.

The equations thus become linear and form an eigenvalue problem for σ , which is solved by the implicitly restarted Arnoldi method (IRAM, Lehoucq, Sorensen & Yang 1998). The largest real part of σ represents the rate of growth (decay) of the disturbance if $\sigma_r > 0$ (< 0). The imaginary part of σ determines whether the solution is steady ($\sigma_i = 0$) or oscillatory ($\sigma_i \neq 0$).

For each azimuthal wavenumber m and an initial guess for Ra , the eigensystem is solved for σ . An iterative method is used to determine the value of Ra that gives $\sigma_r = 0$. The most unstable mode, m_c is the mode having the lowest such Ra , which is the critical Rayleigh number Ra_c .

2.2. Locked inhomogeneous solutions

To account for the inhomogeneous outer boundary, the temperature is decomposed as

$$T(r, \theta, \phi, t) = \Theta(r, \theta, \phi, t) + \epsilon f(r_o, \theta, \phi), \tag{2.13}$$

making the boundary conditions on the unknown temperature Θ homogeneous. The nonlinear equations (2.6) and (2.7) become inhomogeneous,

$$E\nabla^2 L_2 \mathcal{T} = \frac{E}{Pr} \mathbf{r} \cdot \nabla \times [(\mathbf{u} \cdot \nabla) \mathbf{u}] - \mathbf{r} \cdot \nabla \times (\mathbf{z} \times \mathbf{u}), \tag{2.14}$$

$$E\nabla^2 L_2 \nabla^2 \mathcal{P} + Ra L_2 \Theta = -\epsilon Ra L_2 f - \frac{E}{Pr} \mathbf{r} \cdot \nabla \times \nabla \times [(\mathbf{u} \cdot \nabla) \mathbf{u}] - \mathbf{r} \cdot \nabla \times \nabla \times (\mathbf{z} \times \mathbf{u}), \tag{2.15}$$

$$\nabla^2 \Theta - (\mathbf{u} \cdot \nabla) \Theta + 3 = \epsilon (\mathbf{u} \cdot \nabla - \nabla^2) f. \tag{2.16}$$

It is anticipated that the wavenumber m_b will dominate the flow pattern in the neighbourhood of Ra_c (Zhang & Gubbins 1993). We therefore neglect azimuthal wavenumbers that are not multiples of m_b in some numerical computations. This allows thousands of nonlinear calculations to be performed and enables the use of Floquet theory (Verhulst 1993) in the instability analysis (§2.3). Many computations have been performed with full resolution in azimuth to test the lower resolution calculations and results from these computations will be presented in the following sections. The assumption will be shown to be adequate for stable solutions.

The code used to compute locked, inhomogeneous solutions is documented in Gibbons *et al.* (2007) and reproduces the results of Zhang & Gubbins (1993), who used a different code. The code does not time step the equations; solutions are obtained by an iterative method.

2.3. Instability analysis

A perturbation analysis is used to explore the stability of the locked solutions computed in §2.2. Let the locked solutions be denoted by $(\mathcal{T}_0, \mathcal{P}_0, \Theta_0)$, and general perturbations by \tilde{A} , where $\tilde{A} \in \{\tilde{\mathcal{T}}, \tilde{\mathcal{P}}, \tilde{\Theta}\}$. Substituting $\mathcal{T} = \mathcal{T}_0 + \tilde{\mathcal{T}}$, $\mathcal{P} = \mathcal{P}_0 + \tilde{\mathcal{P}}$, $\Theta = \Theta_0 + \tilde{\Theta}$, into (2.1) and (2.2) and neglecting products of the perturbations gives the following set of linear equations

$$E \left(\nabla^2 - \frac{\sigma}{Pr} \right) L_2 \tilde{\mathcal{T}} = \frac{E}{Pr} \mathbf{r} \cdot \nabla \times [(\mathbf{u}_0 \cdot \nabla) \tilde{\mathbf{u}} + (\tilde{\mathbf{u}} \cdot \nabla) \mathbf{u}_0] - \frac{E}{Pr} \mathbf{r} \cdot \nabla \times (\mathbf{z} \times \tilde{\mathbf{u}}), \tag{2.17}$$

$$E \left(\nabla^2 - \frac{\sigma}{Pr} \right) L_2 \nabla^2 \tilde{\mathcal{P}} + Ra L_2 \tilde{\Theta} = -\mathbf{r} \cdot \nabla \times \nabla \times [(\mathbf{u}_0 \cdot \nabla) \tilde{\mathbf{u}} + (\tilde{\mathbf{u}} \cdot \nabla) \mathbf{u}_0] - \frac{E}{Pr} \mathbf{r} \cdot \nabla \times \nabla \times (\mathbf{z} \times \tilde{\mathbf{u}}), \tag{2.18}$$

$$(\nabla^2 - \sigma) \tilde{\Theta} - (\tilde{\mathbf{u}} \cdot \nabla) \Theta_0 - (\mathbf{u}_0 \cdot \nabla) \tilde{\Theta} + 3 = \epsilon (\tilde{\mathbf{u}} \cdot \nabla) f. \tag{2.19}$$

E	m_c	Ra_c	σ_i
3.16×10^{-3}	6	17.1	-10.7
8×10^{-4}	7	21.7	-45.0
6×10^{-4}	8	22.6	-53.6
3×10^{-4}	10	26.6	-90.8
1.2×10^{-4}	14	34.7	-173.5

TABLE 1. Results of linear stability analysis for the onset of convection with homogeneous boundary conditions and $Pr = 100$. For each value of E the most unstable mode m_c , the corresponding critical modified Rayleigh number Ra_c and the drift rate σ_i are given. Negative values of σ_i indicate prograde drift. Data for the lowest value of E considered by Zhang & Gubbins (1993), $E = 3.16 \times 10^{-3}$, are included for comparison.

The system is linear and \mathbf{u}_0, T_0 are periodic in ϕ , so Floquet theory may be applied to the ϕ -dependence of the perturbations. We therefore define the perturbations \tilde{A} as

$$\tilde{A} = \sum_{l=1}^L \sum_{m=0}^l \tilde{A}_l^m(r) Y_l^m(\theta, \phi) \exp\{iM\phi + \sigma t\}, \quad (2.20)$$

where $\exp\{iM\phi\}$ is the Floquet factor, and the Floquet exponent M is an integer. A specified value of M defines the subclasses of perturbations that are retained in the solutions. When $M = 0$ the azimuthal wavenumbers of the perturbation include

$$m = 0, m_c, 2m_c, 3m_c, \dots,$$

and for general M ,

$$m = 0, m_c - M, m_c + M, 2m_c - M, 2m_c + M, \dots,$$

and so on for $M = 0, 1, \dots, m_c - 1$. Often the complete set of M s does not need to be computed as certain values of m_c can make different subclasses equivalent: for example, with $m_c = 7$ subclasses defined by $M = 1$ and $M = 6$ are identical, as are $M = 2$ and $M = 5$, and $M = 3$ and $M = 4$.

The system (2.17)–(2.19) is an eigensystem for $\sigma = \sigma_r + i\sigma_i$, which is solved by IRAM. The most unstable mode is the mode with largest σ_r , which is found by testing all values of M . The code used for stability analysis is documented in Gibbons *et al.* (2007) and reproduces the results of Zhang & Gubbins (1993).

2.4. Time integration

To determine the properties of unstable solutions predicted by the linear analysis it is necessary to integrate the full set of nonlinear equations, (2.1)–(2.3), by time stepping. The details of the numerical method used here can be found in Willis, Sreenivasan & Gubbins (2007); the code differs from that of Gibbons *et al.* (2007) and has been used as an independent verification of our results.

Initial conditions for the inhomogeneous solutions are taken from the corresponding homogeneous solutions. This allows us to track the initial response of the system to the imposed inhomogeneous boundary condition as well as its long-term evolution.

3. Onset of convection with homogeneous boundary conditions

Results of linear stability analysis for the onset of convection with homogeneous boundaries are presented in table 1 for each value of E considered. The most unstable

E	m_c	Ra_c	σ_i
3.16×10^{-3}	6	17.1	-10.7
8×10^{-4}	7	21.7	-45.0
6×10^{-4}	8	22.7	-53.6
3×10^{-4}	10	26.7	-90.8
1.2×10^{-4}	14	34.8	-173.7

TABLE 2. Same as table 1 but with $Pr = \infty$.

mode is always symmetric with respect to the equator and the value of m_c increases as E decreases. The radial length scale of the motions also decreases with E , however the axial length scale is large for all E : the flow is organized into rolls (see figure 3). The frequency is negative for every E , indicating prograde drift of the rolls. The travelling wave is steady in a frame of reference co-rotating with the rolls.

Table 2 shows the same information as table 1 but for $Pr = \infty$. Values of m_c and σ_i are identical to the case with $Pr = 100$. Values of Ra_c differ at the fourth significant figure between the two cases. It therefore appears that $Pr = 100$ is a good approximation to the limit $Pr \rightarrow \infty$ near the onset of convection. Calculations in the marginally supercritical regime undertaken with both $Pr = 100$ and $Pr = \infty$ further verify this.

4. Locked inhomogeneous solutions

4.1. Primary convection

Figure 1 shows kinetic energy as a function of Ra for $E = 8 \times 10^{-4}$, 6×10^{-4} and 3×10^{-4} . For $E \leq 1.5 \times 10^{-4}$, there is no value of ϵ that forces a steady locked solution for $Ra \geq Ra_c$. This interesting result will be discussed separately in §6.

It is clear that no spatial resonance is observed in figure 1. The presence or otherwise of a resonance is a complicated issue that depends on both ϵ and m_b . Figure 2 shows kinetic energy as a function of Ra for various E at three values of ϵ . Solutions with $E = 3.16 \times 10^{-3}$ and $\epsilon = 0.18$ correspond to solutions obtained by Zhang & Gubbins (1993). As E decreases with $\epsilon = 0.18$, the primary resonance observed by Zhang and Gubbins disappears at $E = 2.3 \times 10^{-3}$. Increasing ϵ to 0.3 results in an observed resonance when $E = 2.5 \times 10^{-3}$, however no other value of E considered shows resonance. When ϵ is further increased to 0.5 resonance is not observed for any E considered. Hence large ϵ and large m_b appear to destroy resonance, with no resonance observed for $E < 2 \times 10^{-3}$. At low E , resonance is prohibited because the boundary effects and the buoyancy-driven flow do not reinforce each other, as described below.

A steady locked solution is shown in figure 3 with its homogeneous counterpart. The homogeneous solution exhibits rolls aligned with the rotation axis and centred about the mid-radius of the shell (figures 3*b* and 3*f*). The inhomogeneous solution, while also being nearly two-dimensional (figure 3*e*), contains two layers of rolls, one located near the IB and the other near the OB (figure 3*a*).

The two-layer flow pattern forms because, at low E , the radial length scale of the convection rolls is small, resulting in weak motion near the boundaries (figure 3*b*). Fluid near the IB is locked by the boundary anomalies acting near the poles of the outer spherical shell, forming the inner layer in figure 3(*a*). Near the OB boundary

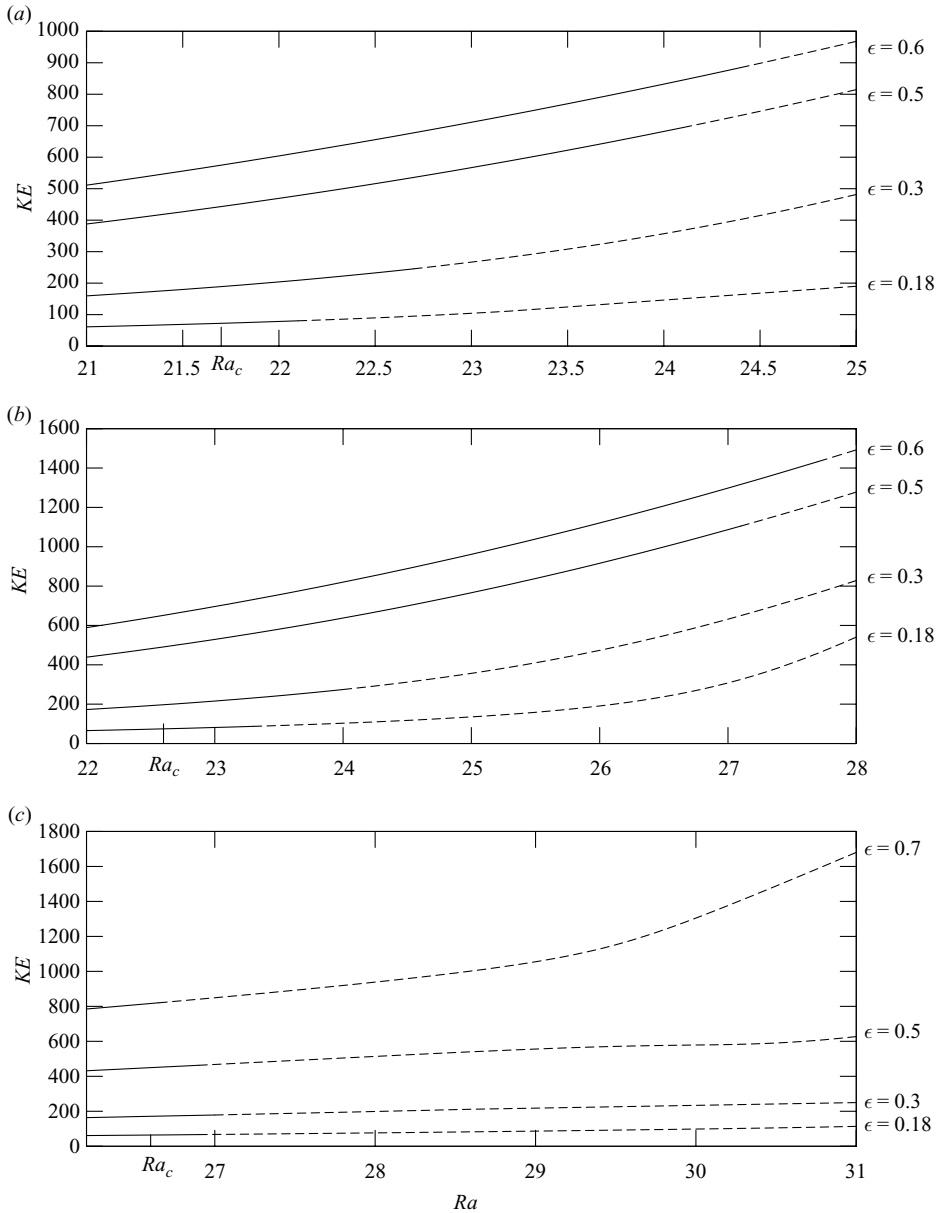


FIGURE 1. Kinetic energy as a function of Ra for selected values of ϵ . Cases are $E = 8 \times 10^{-4}$ with a Y_7^7 boundary condition (a), $E = 6 \times 10^{-4}$ with a Y_8^8 boundary condition (b), and $E = 3 \times 10^{-4}$ with a Y_{10}^{10} boundary condition (c). Stable solutions are shown as thick lines, unstable solutions as broken lines. Lines are continued until the stability boundary is found. All stable solutions are steady. Solutions are obtained at intervals of $Ra = 0.1$.

effects dominate, creating a second layer of rolls that are out of phase with the boundary temperature anomalies (figure 3a), as would result if convection was absent (Zhang & Gubbins 1992). Convection is suppressed near the mid-radius of the shell where it is preferred in the homogeneous case (compare figures 3b and 3a).

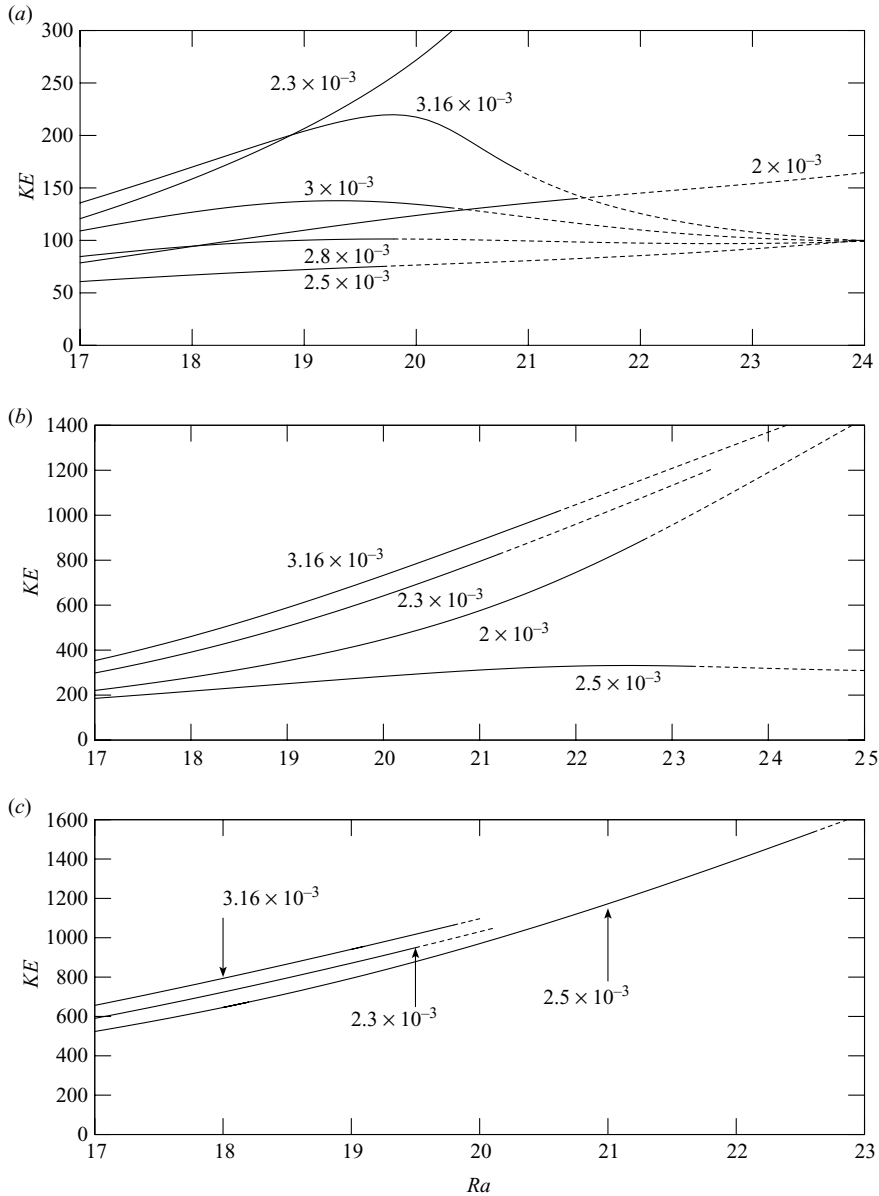


FIGURE 2. Kinetic energy as a function of Ra for various E with $\epsilon = 0.18$ (a), $\epsilon = 0.3$ (b) and $\epsilon = 0.5$ (c) with $m_b = m_c$. Stable solutions are shown as thick lines, unstable solutions as broken lines. Lines are continued until the stability boundary is found. All stable solutions are steady. Solutions are obtained at intervals of $Ra = 0.1$. The solution with $E = 3.16 \times 10^{-3}$ and $\epsilon = 0.18$ corresponds to a solution in figure 1 (a) of Zhang & Gubbins (1993); it is included for comparison with the earlier work.

Figure 4 shows that the two-layer pattern persists at lower values of E , with both layers becoming more closely confined near the boundaries. As E decreases the interior flow is less affected by the boundary anomalies and hence larger values of ϵ are required for locking (figure 1). However, larger ϵ does not cause the two-layer pattern as it does not alter the imposed scale of the boundary anomalies.

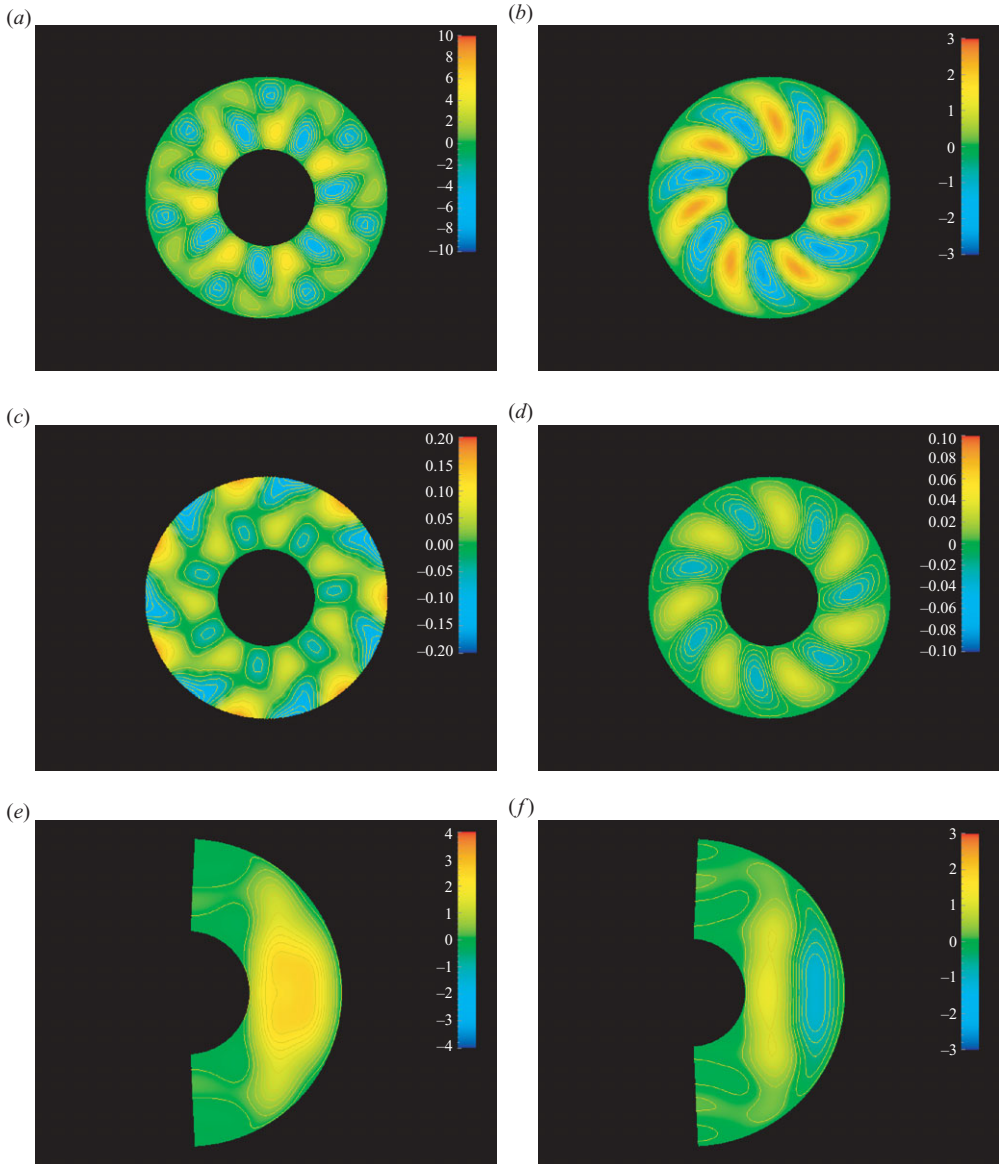


FIGURE 3. Properties of homogeneous (right) and inhomogeneous (left) solutions for $Ra = 21.7$, $E = 8 \times 10^{-4}$, $\epsilon = 0.3$ with an imposed Y_7^7 boundary condition: radial velocity (u_r) in the equatorial plane (*a*, *b*); temperature field in the equatorial plane (*c*, *d*); u_r in the meridional plane (*e*, *f*), taken at $\phi = \pi$. $\phi = 0$ is at the rightmost edge of the equatorial section. The homogeneous solution drifts in the prograde sense while the inhomogeneous solution is steady.

This pattern does not arise at larger E because the radial length scale of the convection rolls is greater. The effects of the boundary anomalies penetrate deeper into the shell (compare figures 4*d*, 4*c* and 3*c*) as large E gives low $m_c (= m_b)$. Hence there is no region of the shell in which convection or boundary effects dominate and the resulting interaction leads to steady solutions that retain the single-layer pattern seen in homogeneous convection, as found by Zhang & Gubbins (1993). Calculations

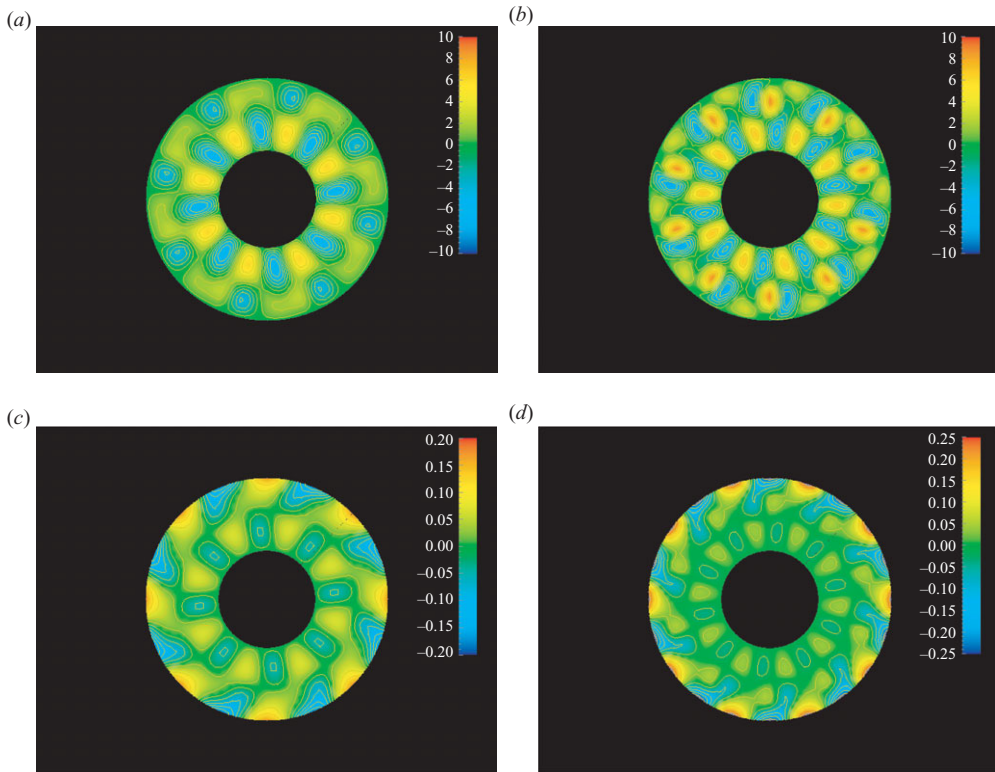


FIGURE 4. u_r in the equatorial plane (top) and T in the equatorial plane (bottom) for the inhomogeneous solutions with $E = 6 \times 10^{-4}$, $Ra = 22.6$ and $\epsilon = 0.3$ with an imposed Y_8^8 boundary condition (left), and $E = 3 \times 10^{-4}$, $Ra = 26.6$ and $\epsilon = 0.5$ (right) with an imposed Y_{10}^{10} boundary condition. Meridional sections have been omitted as they are qualitatively identical to the case $E = 8 \times 10^{-4}$ above. Both solutions are steady.

in figure 2 indicate that a transition from the single-layer to two-layer pattern occurs at $E \approx 1.5 \times 10^{-3}$, close to the value of E where m_c changes from $m_c = 6$ to $m_c = 7$.

4.2. Secondary convection

Figure 5 shows kinetic energy as a function of Ra for $E = 6 \times 10^{-4}$, 3×10^{-4} and 1.2×10^{-4} , when $m_b = m_c/2$. Larger values of ϵ are required to obtain steady locked solutions than for the case $m_b = m_c$ as found by Zhang & Gubbins (1993). The spatial resonance phenomenon does not occur for any of the cases considered. We show that this may be explained by the dominance of the boundary-driven flows.

Steady locked solutions exist for $E = 1.2 \times 10^{-4}$ when a Y_7^7 boundary condition is imposed, indicating that boundary conditions with $m_b < m_c$ are the most effective in obtaining steady locked solutions as E decreases. At larger E , boundary conditions with $m_b = m_c$ were found to be most effective (Zhang & Gubbins 1993). The pattern of convection is shown in figure 6 for $E = 1.2 \times 10^{-4}$. The boundary mode dominates the convection pattern; it is modulated by the $m_c = 14$ mode. High-amplitude regions of u_r are out of phase with the boundary temperature anomalies because of the dominance of the inhomogeneous boundary condition. Figure 7 shows that the boundary mode dominates at higher E . With $m_b = m_c/2$ the penetration depth of boundary anomalies

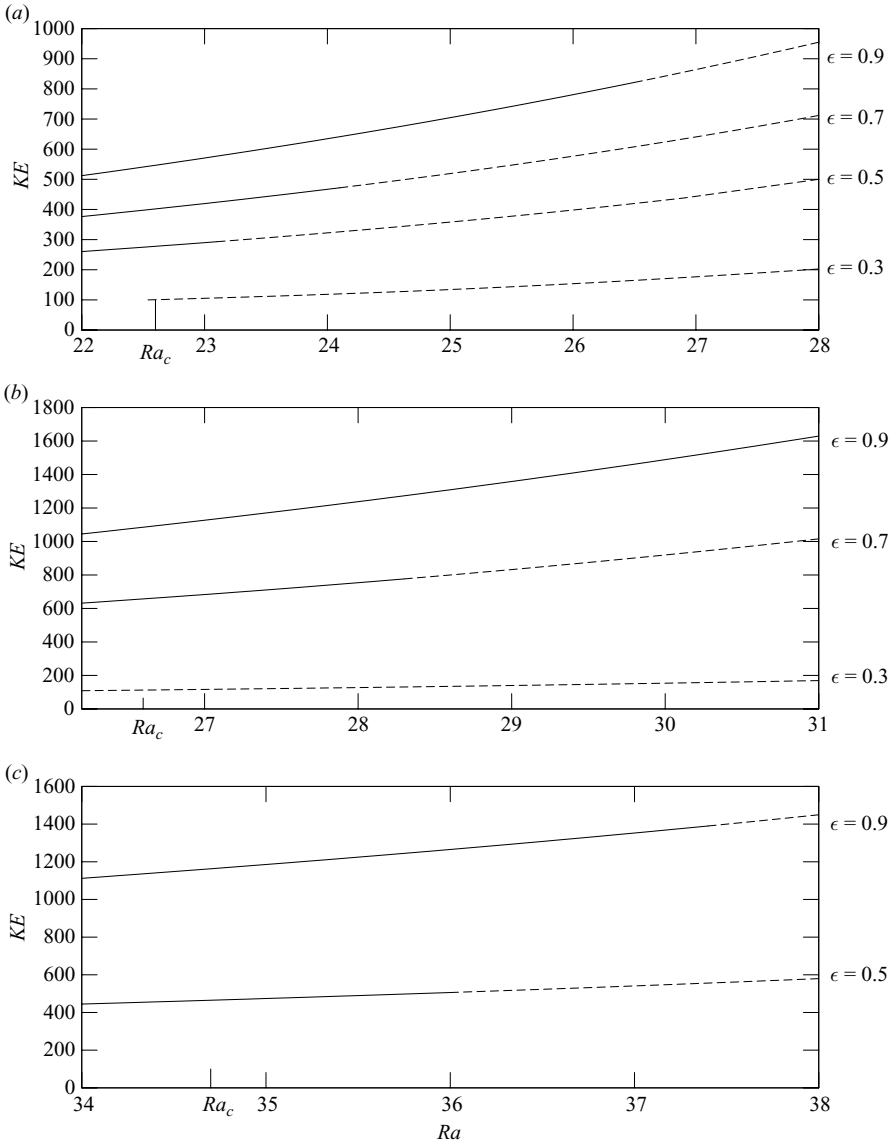


FIGURE 5. Kinetic energy as a function of Ra for selected values of ϵ . Cases are $E = 6 \times 10^{-4}$ ($m_c = 8$) with an imposed Y_4^4 boundary condition (a), $E = 3 \times 10^{-4}$ ($m_c = 10$) with an imposed Y_5^5 boundary condition (b), and $E = 1.2 \times 10^{-4}$ ($m_c = 14$) with an imposed Y_7^7 boundary condition (c). Filled lines denote stable solutions; dashed lines denote unstable solutions. Curves are continued until the stability boundary has been determined. All stable solutions are steady. Solutions are obtained at intervals of $Ra = 0.1$.

is greater than for $m_b = m_c$ and this can promote steady locked solutions. Decreasing m_b further requires a larger ϵ in order to obtain steady solutions because the amplitude of boundary effects weakens. ϵ cannot be increased indefinitely as the steady state becomes unstable to boundary-driven instabilities. Hence steady locked solutions are unlikely to exist for all values of m_b , especially when the difference between m_c and m_b is large.

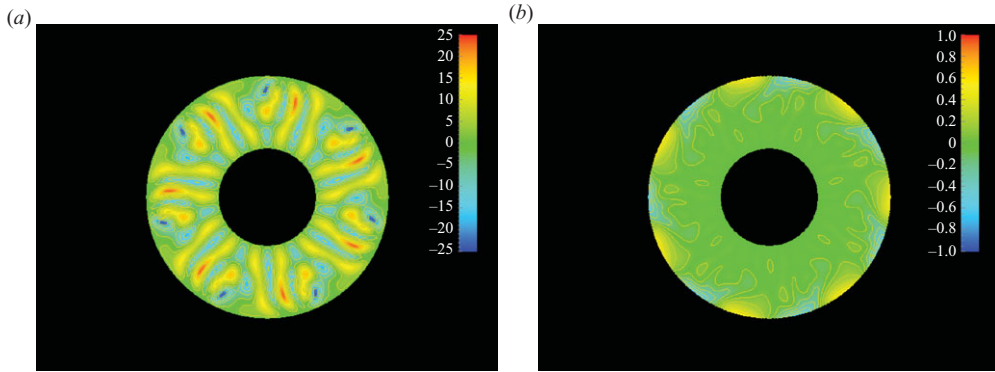


FIGURE 6. u_r (a) and T (b) in the equatorial plane for $E = 1.2 \times 10^{-4}$, $Ra = 34.7$ and $\epsilon = 0.9$ with an imposed Y_7^7 boundary condition. The solution is steady, equatorially symmetric and dominated by the boundary mode, $m_b = 7$.

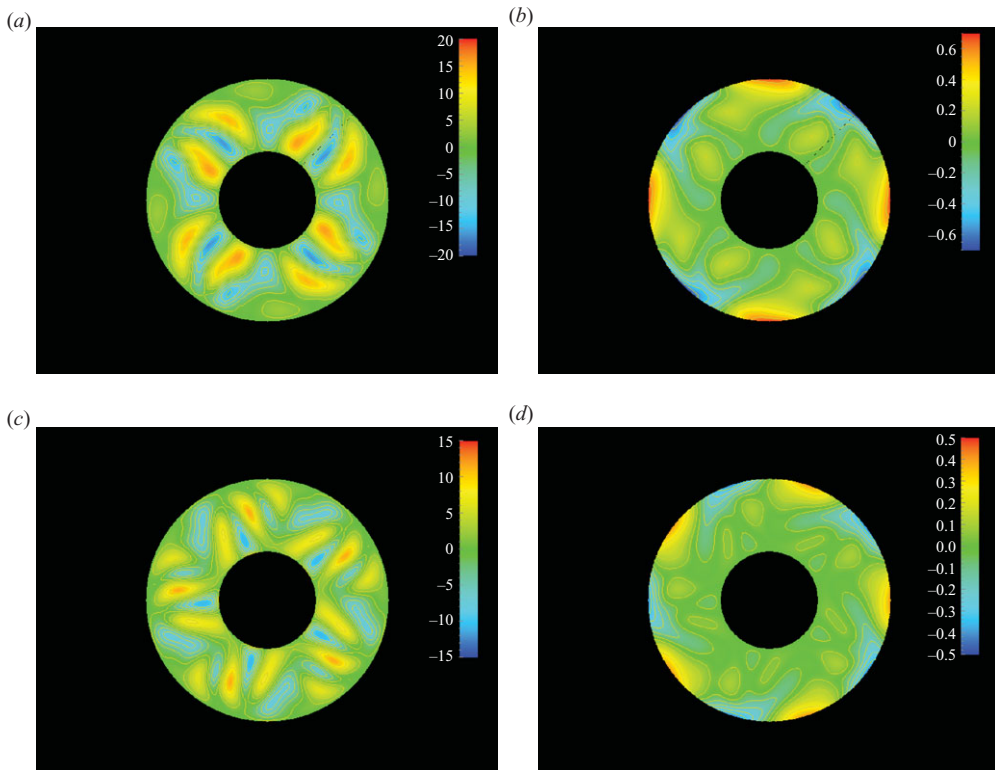


FIGURE 7. u_r (left) and T (right) in the equatorial plane for $E = 6 \times 10^{-4}$, $Ra = 22.6$, $\epsilon = 0.9$ with an imposed Y_4^4 boundary condition (top) and $E = 3 \times 10^{-4}$, $Ra = 26.6$, $\epsilon = 0.9$ with an imposed Y_5^5 boundary condition (bottom). Both solutions are steady and equatorially symmetric.

5. Instability analysis

5.1. Stability of primary convection

Figure 8 displays instability curves for each value of E when $m_b = m_c$. Curves for $E = 1.2 \times 10^{-4}$ are not shown as no steady solution exists from which to perturb.

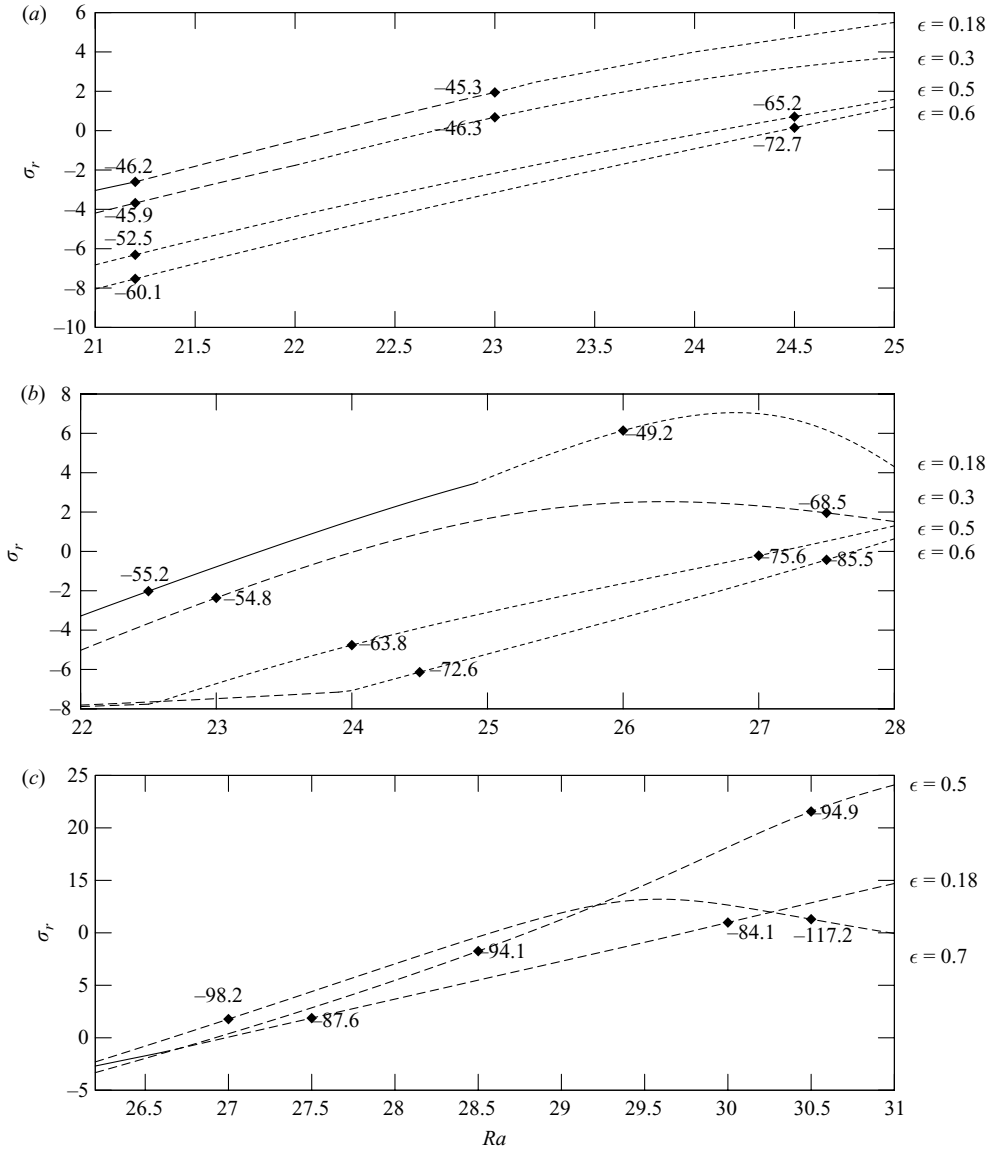


FIGURE 8. σ_r as a function of Ra for (a–c) $E = 8 \times 10^{-4}$, 6×10^{-4} and 3×10^{-4} with $m_b = m_c$. Different curves represent different values of ϵ . Numbers next to the curves denote the frequency σ_i of the most unstable disturbance at that point on the curve. Full lines show that the most unstable disturbance is from the $M = 0$ subclass; long dashed lines show the $M = 1$ subclass; short dashed lines show the $M = 2$ subclass (see § 2.3 for details).

Solutions at higher values of ϵ cannot be calculated as too much lateral heating induces an instability: no steady state exists to perturb. The drift rate σ_i is non-zero at all points on all curves, suggesting that the bifurcations are of Hopf-type, although the precise nature of any bifurcation cannot be determined with certainty without integrating the full nonlinear equations.

Solutions in figure 8 become unstable to perturbations from the $M = 1$ and $M = 2$ subclasses, in contrast to Zhang & Gubbins (1993) who found most unstable modes

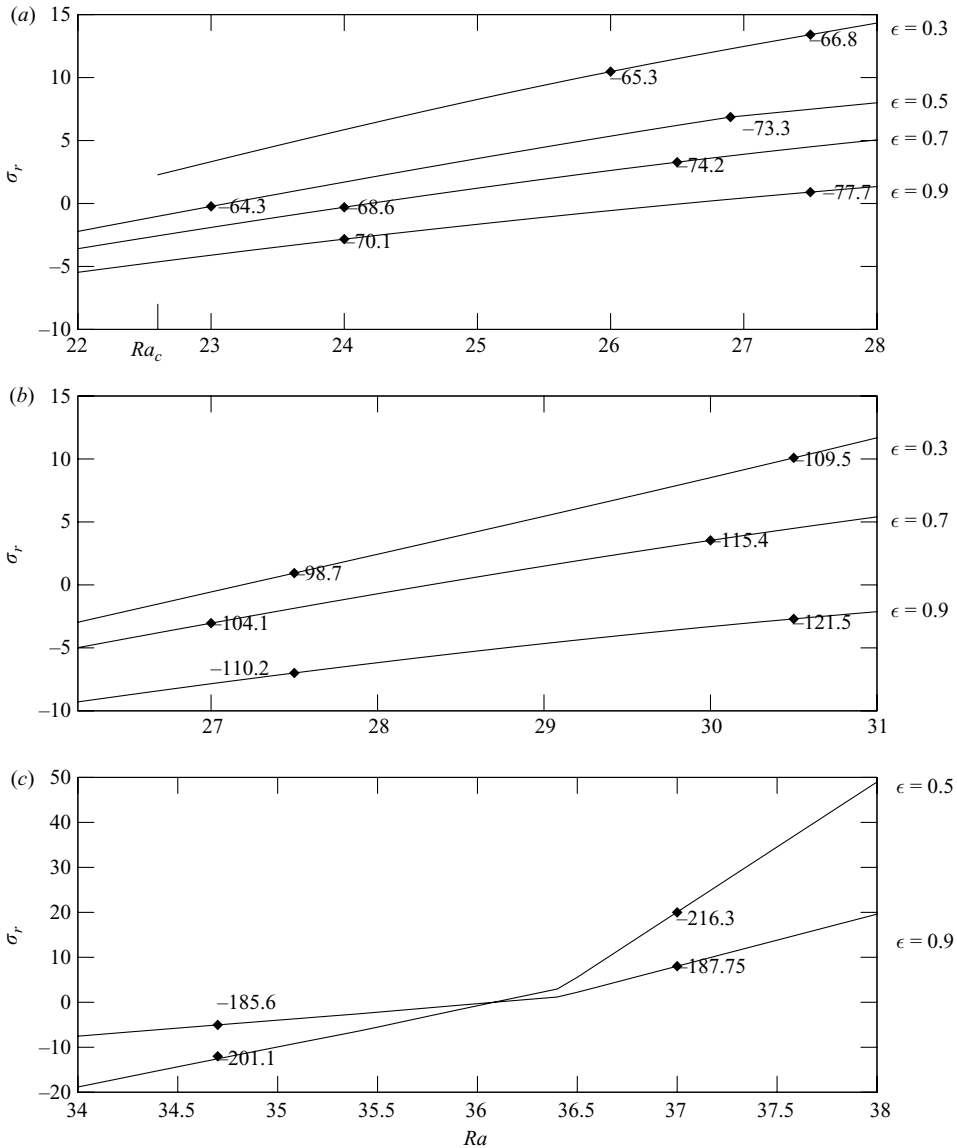


FIGURE 9. Stability curves for (a–c) $E = 6 \times 10^{-4}$, 3×10^{-4} and 1.2×10^{-4} with $m_b = m_c/2$. Different curves represent different values of ϵ . Numbers next to the curves denote the frequency, σ_i , of the most unstable mode at that point on the curve. Filled lines show that the most unstable disturbance is from the $M = 0$ subclass.

with $M = 0$. The explanation is likely to lie with the resonance observed by Zhang and Gubbins near Ra_c ; no resonance is observed here and instability sets in by a different mode breaking the constraint imposed by the boundary condition.

5.2. Stability of secondary convection

Figure 9 shows instability diagrams for $E = 6 \times 10^{-4}$, 3×10^{-4} and 1.2×10^{-4} when $m_b = m_c/2$. The stability boundaries for all E and ϵ are characterized by $\sigma_i \neq 0$. Unstable solutions bifurcate from steady solutions with $M = 0$ for all cases considered. This contrasts with the results of § 5.1 where $M \neq 0$ is preferred. The change is due to

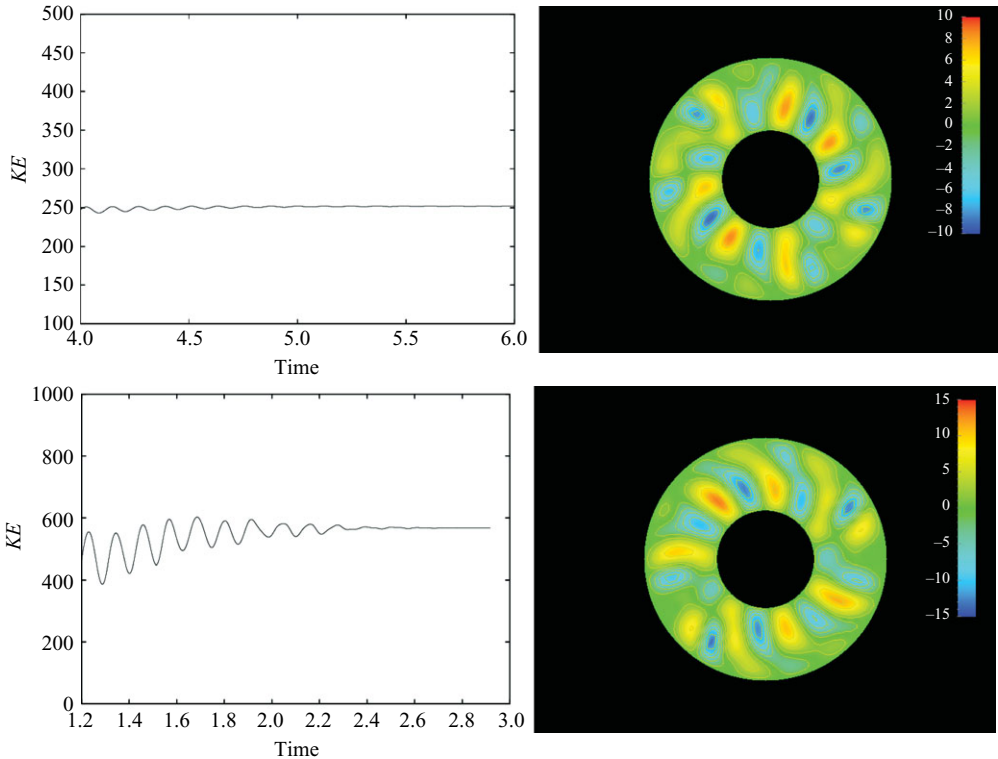


FIGURE 10. Kinetic energy time-series (left) and snapshots of u_r in the equatorial plane (right) for $E = 8 \times 10^{-4}$, $\epsilon = 0.3$ and $m_b = m_c = 7$, with $Ra = 23$ (top) and $Ra = 25$ (bottom). The former case is at the stability boundary, while the latter is above the stability boundary. Both solutions are periodic, drifting in the prograde sense and symmetric with respect to the equatorial plane.

the introduction of a new competing scale. With $m_b \neq m_c$ the interaction between the two dominant modes results in the loss of stability of steady locked solutions.

6. Time integration

6.1. Unstable primary solutions

Figure 10 shows a comparison between two nonlinear solutions at, and just above, the stability boundary where the linear analysis predicts $M = 2$ as the most unstable perturbation subclass. Both solutions are periodic in time with prograde drift near the mid-radius of the shell, but they differ markedly in planform. At the stability boundary modes from the $M = 2$ subclass are barely visible; the equatorial pattern is dominated by seven pairs of convection cells, unevenly organized in azimuth. Above the stability boundary the pattern is dominated by the $m = 9$ mode (a member of the $M = 2$ subclass) and modulated by a large $m = 7$ mode. Stationary flows near the boundaries are observed as in previous sections. Steady locked solutions that become unstable to perturbations from the $M = 1$ subclass (e.g. for $E = 3 \times 10^{-4}$ in figure 8) display similar properties. Clearly the effect of the imposed boundary condition diminishes quickly once Ra is increased past the stability boundary for steady solutions and locking is quickly lost.

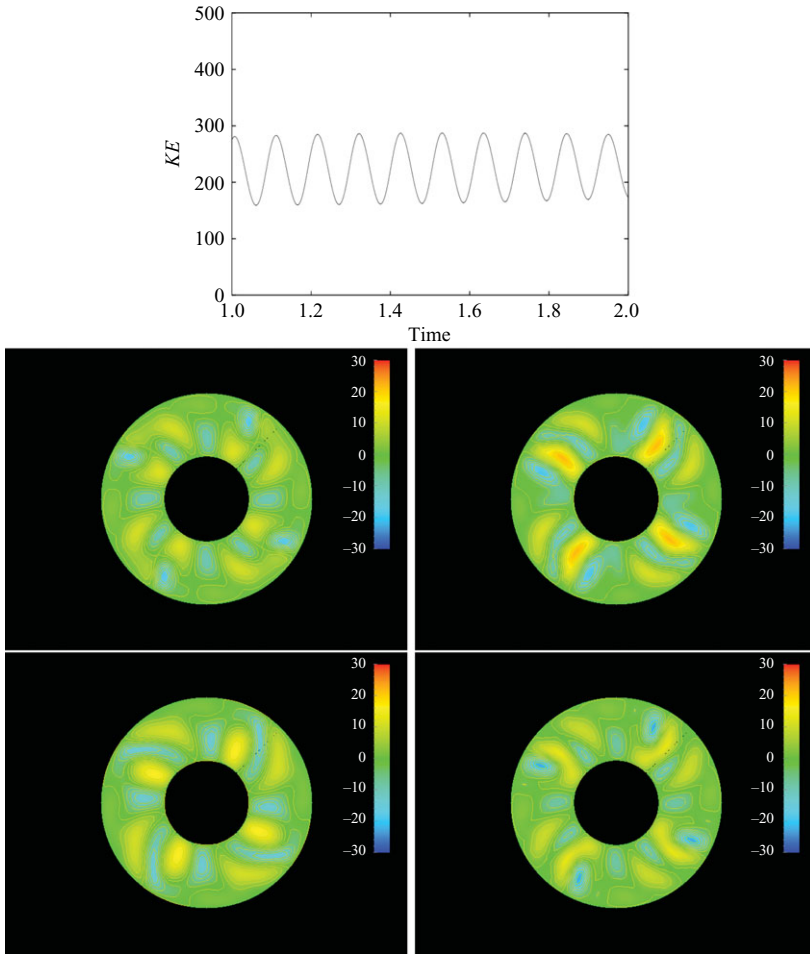


FIGURE 11. Kinetic energy time series (top) and snapshots of u_r in the equatorial plane for $Ra = 24.1$, $E = 6 \times 10^{-4}$, $\epsilon = 0.7$ with $m_b = 4$. Snapshots are evenly spaced in time, spanning one period of the oscillation and should be left–right and top–bottom. The solution is equatorially symmetric, so meridional sections have been omitted.

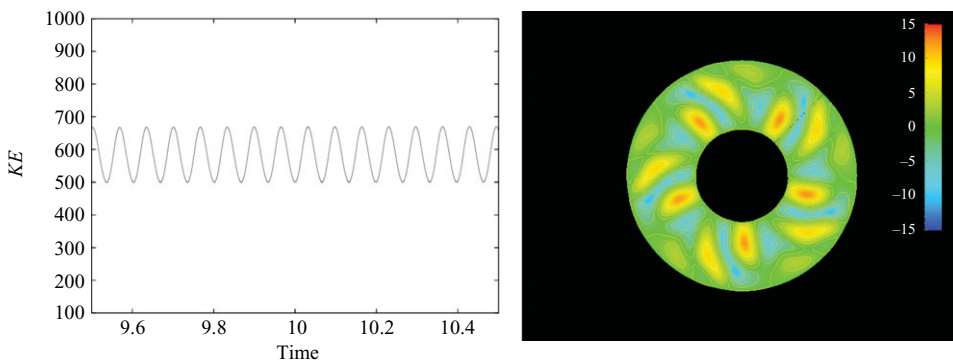


FIGURE 12. Kinetic energy time series (a) and a snapshot of u_r in the equatorial plane (b) for $E = 3 \times 10^{-4}$, $Ra = 28.5$, $\epsilon = 0.7$ with $m_b = 5$. Linear analysis predicts $M = 0$ as the most unstable perturbation subclass. The solution is periodic, equatorially symmetric and contains five persistent nests of convection rolls.

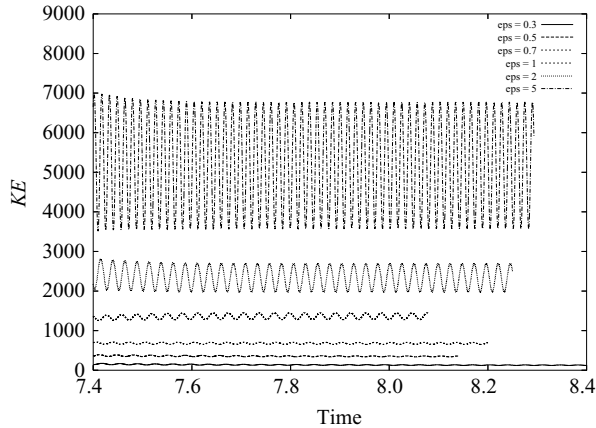


FIGURE 13. Kinetic energy time series for selected values of ϵ at $E = 1.2 \times 10^{-4}$ and $Ra = 34.7$ with $m_b = m_c = 14$. All solutions are periodic.

6.2. Unstable secondary solutions

Figure 11 shows a periodic solution of secondary convection at the stability boundary where linear analysis predicts $M = 0$ as the most unstable perturbation subclass.

The equatorial flow pattern contains four persistent centres at $\phi = \pi/4, 3\pi/4, 5\pi/4$ and $7\pi/4$ where convection rolls become trapped, forming a nest. Rolls confined within the nests are time-dependent but do not break out of the nests; the nests are locked. Nests are out of phase with the boundary anomalies, illustrating the dominance of the boundary effects. Another example of this type of flow pattern is shown in figure 12, which is above the stability boundary, showing that this behaviour persists to higher values of Ra .

6.3. Solutions for $E \leq 1.2 \times 10^{-4}$

6.3.1. Primary convection

No values of ϵ in the range $0.1 \leq \epsilon \leq 10$ force a steady solution for $E < 1.5 \times 10^{-4}$ when $Ra \geq Ra_c$. Kinetic energy time series for selected values of ϵ are shown in figure 13. Low ϵ solutions are dominated by the buoyancy-driven flow and slightly modulated by boundary effects. Solutions with $\epsilon \geq 5$ are unstable because of large temperature variations on the boundary.

A periodic solution for $\epsilon = 0.7$ is shown in figure 14 at evenly spaced time points spanning one period of oscillation in figure 13. The pattern maintains a striking three-layer structure with stationary layers near the boundaries and drift confined to a relatively thin region between these layers. The middle layer is where convection rolls establish in the homogeneous problem. Because of the small radial length scales of the buoyancy-driven and boundary-driven flows, the boundary anomalies are unable to overcome the convective driving, even at the stability boundary of homogeneous convection, and therefore no stable equilibrium exists. Figure 15 shows that such a flow occurs at $E = 2 \times 10^{-5}$.

Increasing ϵ increases the amplitude of the boundary-driven flow, but the radial scale of this flow is determined by m_b . Hence a larger ϵ has little effect on the middle layer and an $\epsilon \geq 5$ produces vigorous motions due to the large boundary temperature variations. Steady locked solutions can only be achieved with $m_b < m_c$ (§4.2).

Zhang & Gubbins (1993) considered an analogous Landau equation as a framework for understanding steady locked solutions. The framework equation implies that there always exists an ϵ that will yield a steady solution for any amplitude of the

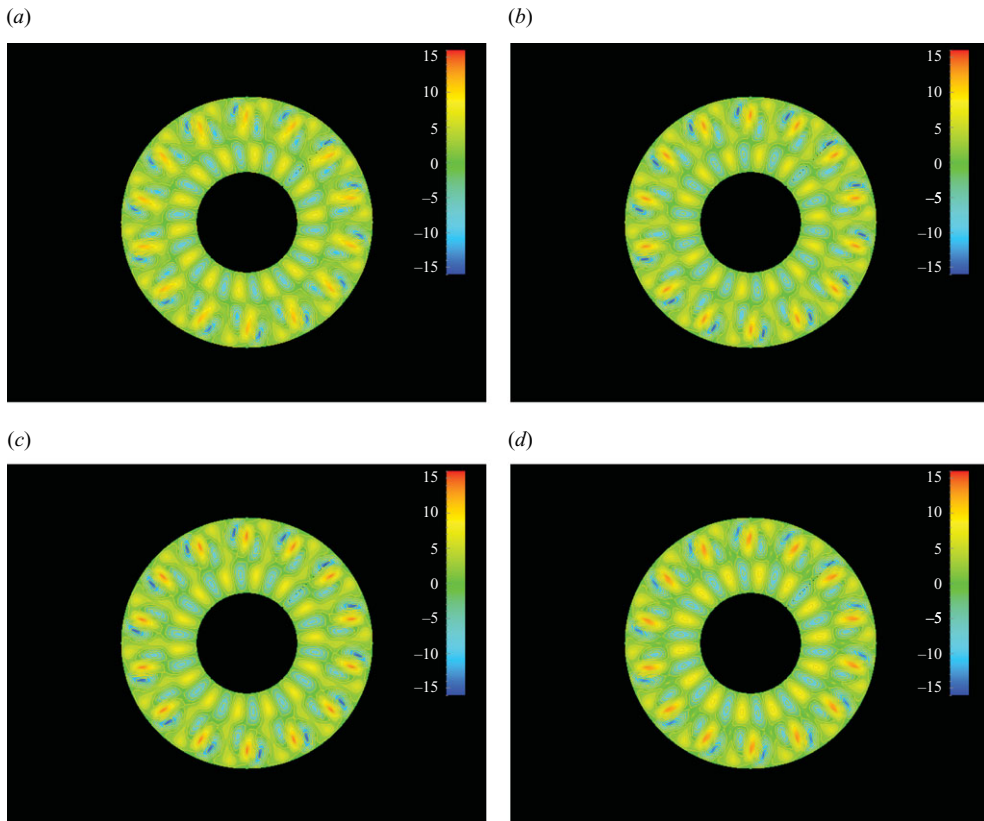


FIGURE 14. Snapshots of u_r in the equatorial plane, at equally spaced time points spanning one period, for $E = 1.2 \times 10^{-4}$, $Ra = 34.7$, $\epsilon = 0.7$ with an imposed Y_{14}^{14} boundary condition, showing the oscillation in the interior of the domain. Snapshots should be viewed in alphabetical order.

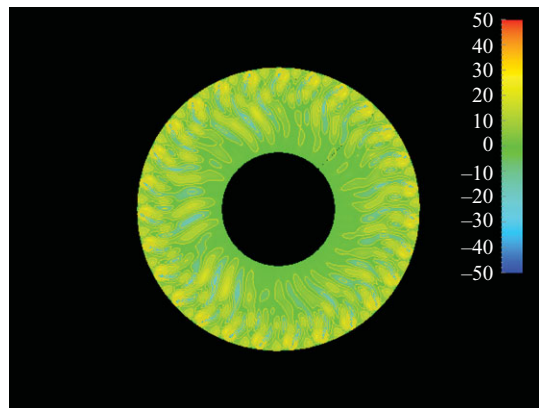


FIGURE 15. Snapshot of u_r in the equatorial plane for $E = 2 \times 10^{-5}$, $Ra = 62$, $\epsilon = 1$, with an imposed Y_{27}^{27} boundary condition. The most unstable mode $m_c = 27$, and $Ra_c = 61.1$. The solution is equatorially symmetric.

driving force. The result in this section shows the shortcomings of this analogue. The framework equation does not account for the Coriolis force, yet it is the rotation that dictates the selection of the mode at the onset of convective instability. The dominant

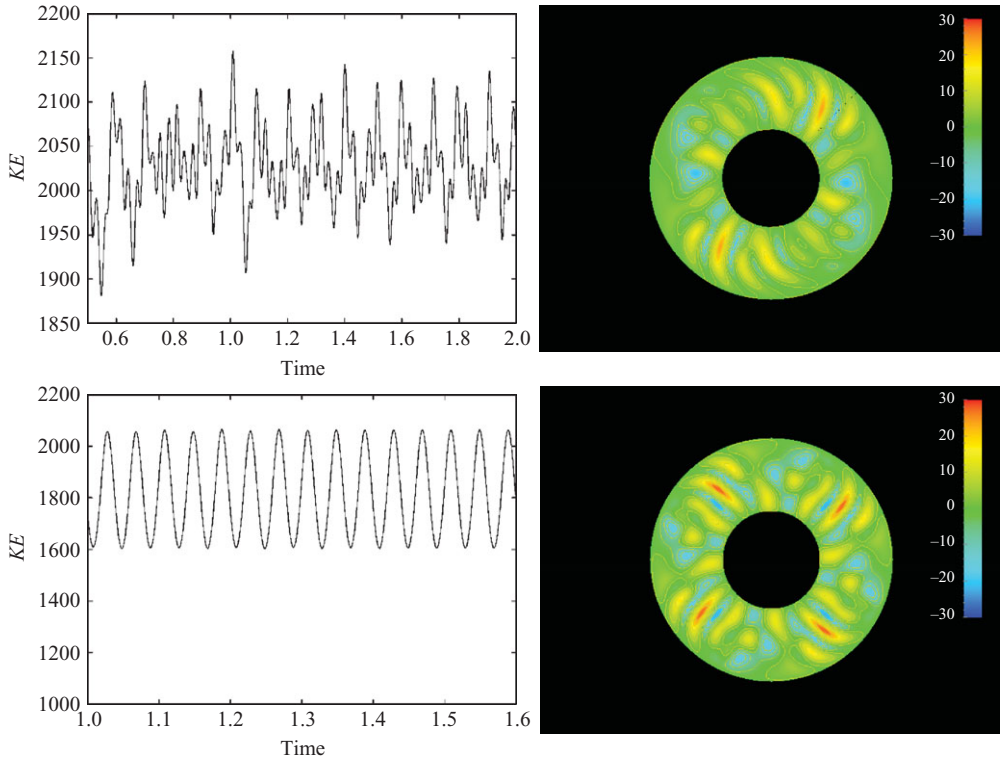


FIGURE 16. Kinetic energy time series (left) and snapshots of u_r in the equatorial plane (right) for $E = 1.2 \times 10^{-4}$, $Ra = 34.7$, $\epsilon = 0.9$ with $m_b = 2$ (top) and $m_b = 4$ (bottom). The solutions are equatorially symmetric so meridional sections have been omitted.

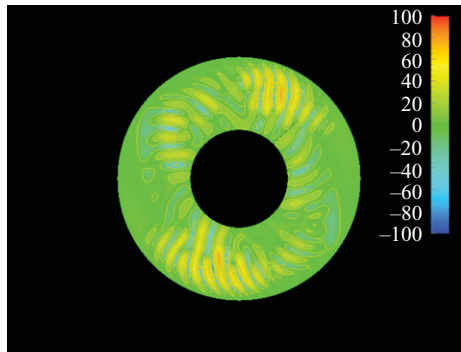


FIGURE 17. Snapshot of u_r in the equatorial plane for $E = 2 \times 10^{-5}$, $\epsilon = 1$ and $Ra = 62$ using a boundary condition with $m_b = 2$. The most unstable mode is $m_c = 27$ and $Ra_c = 61.1$. The solution is equatorially symmetric.

scales of the convective flow and the boundary anomalies, together with ϵ , determine whether a steady locked solution is obtained.

6.3.2. Secondary convection

No steady locked solutions have been found at $E = 1.2 \times 10^{-4}$ for $m_b < 7$. Figure 16 shows a comparison of unstable nonlinear solutions with different applied boundary

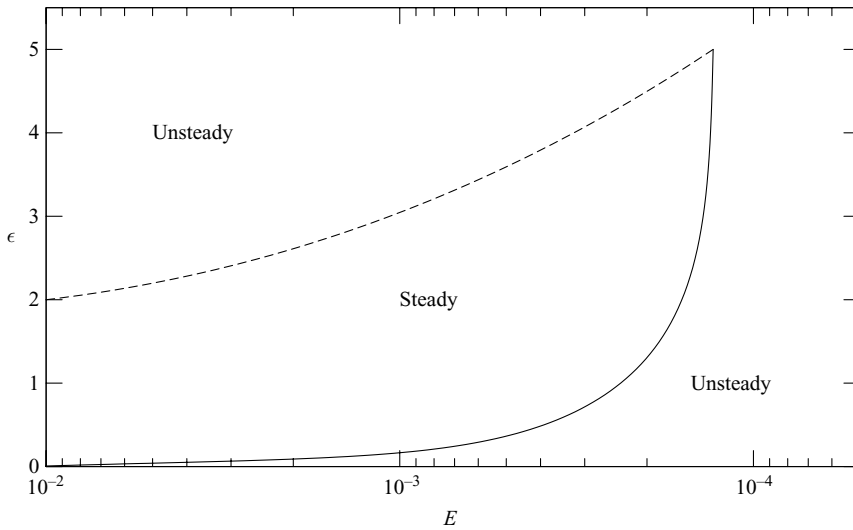


FIGURE 18. Schematic diagram of the boundaries between steady and unsteady solutions in the (E, ϵ) plane for inhomogeneous convection at $Ra = Ra_c$ with $m_b = m_c$. The solid line describes the minimum value of ϵ for a given E that produces a steady locked solution. Above the dashed line large lateral temperature variations on the OB prevent locking.

inhomogeneities. The kinetic energy time series for $m_b = 2$ is highly time-dependent, while for $m_b = 4$ the solution is periodic. This demonstrates the weakening of boundary effects as $m_c - m_b$ grows but flows are still locked. Both solutions display the nesting of convection rolls shown in § 6.2, with the number of nests dictated by the imposed scale of the boundary condition. Nesting becomes more prominent when the wavenumbers m_c and m_b differ substantially: the effect is much weaker for the solutions of § 6.2. Further calculations have shown that such solutions persist to $E = 2 \times 10^{-5}$, shown in figure 17, and up to a value of Ra that is twice supercritical.

7. Conclusion

Our principal conclusion is that boundary conditions having spatial variations larger than the preferred mode of homogeneous convection ($m_b < m_c$) are most effective in obtaining steady solutions as E decreases. When $m_b = m_c$, steady solutions occur, shown schematically in figure 18 for $Ra = Ra_c$. The minimum value of ϵ needed to obtain a steady locked solution increases rapidly as E becomes small (solid line). Large values of ϵ destabilize the system due to lateral heating (dashed line) and the two curves therefore intersect. Such a region was not found for boundary conditions with $m_b < m_c$. This may be due to the larger radial length scale of boundary conditions with $m_b < m_c$ than those with $m_b = m_c$: in the former case fluid motions are suppressed deeper into the shell, with boundary effects dominating if ϵ is high enough.

The spatial resonance phenomenon observed by Zhang & Gubbins (1993) disappears when E is decreased or ϵ is increased. This potentially important result does not have a simple explanation, partly because resonance is a nonlinear phenomenon: it is related to the amplitude of convection, which for a given E is dictated primarily by ϵ in the neighbourhood of Ra_c . We suggest that loss of resonance at low E occurs because of the small characteristic scales of the buoyancy-driven flow and

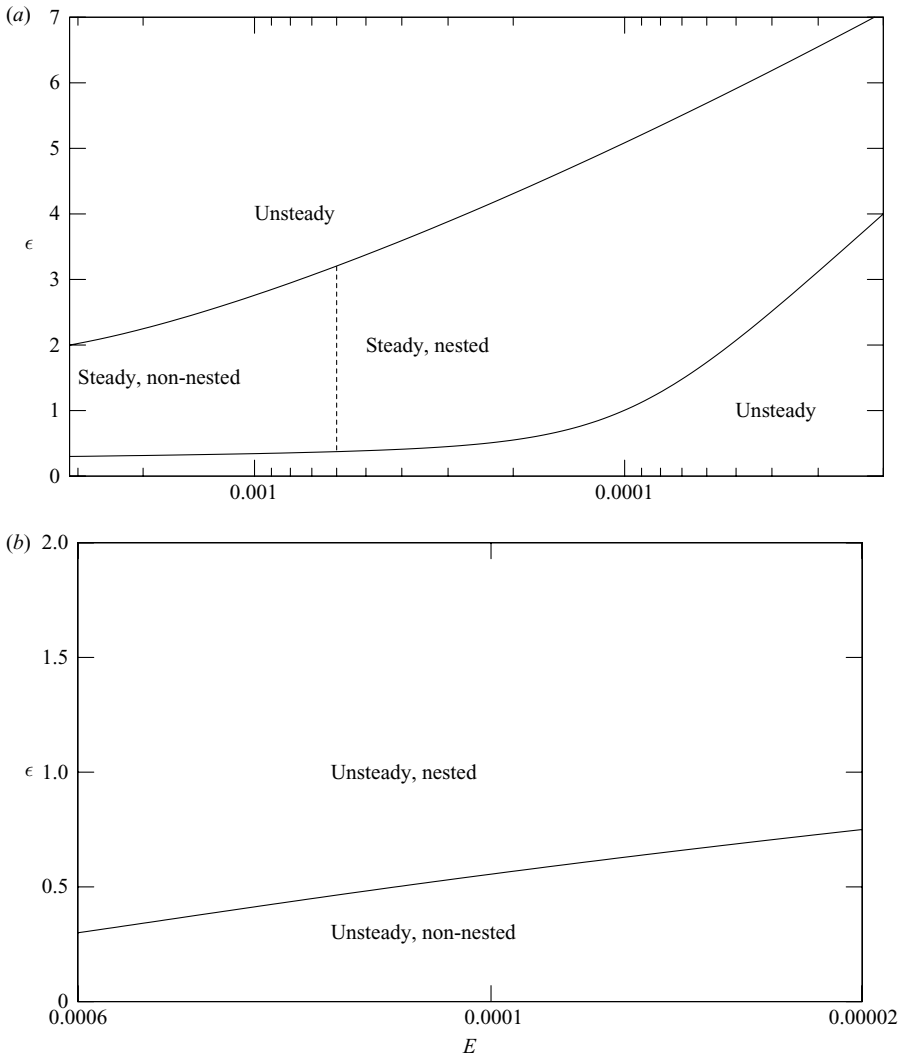


FIGURE 19. Schematic diagram of the boundaries between steady and unsteady solutions in the (E, ϵ) plane for inhomogeneous convection for Ra in the neighbourhood of Ra_c with $m_b = m_c/2$ (a) and $m_b = 2$ (b).

the boundary anomalies: the boundary anomalies no longer reinforce the buoyancy-driven flow and a two-layer flow pattern emerges. Changing ϵ does not affect the characteristic scales of the flow and boundary anomalies; we may expect an absence of resonance at low E for any $\epsilon > 0$, however for large E a different mechanism must occur to produce the behaviour in figure 2. It is the nonlinear interactions between the boundary anomalies and the buoyancy-driven flow that ultimately determine whether a resonance exists at high ϵ and it is this complexity that prevents a more detailed understanding of the resonance phenomenon.

Solutions of secondary convection ($m_b < m_c$) are summarized in figure 19 for two cases: (i) when $m_b = m_c/2$, the case studied in detail in this paper; (ii) when $m_b = 2$, the geophysically relevant case. When $m_b = m_c/2$ steady solutions are always possible for the range of E studied. For $E \leq 6 \times 10^{-4}$, solutions in the neighbourhood of Ra_c

show that convection rolls cluster into nests that are out of phase with the boundary anomalies and remain trapped for many thermal diffusion times. Our calculations indicate that nested solutions appear abruptly as E is decreased and are insensitive to the value of ϵ (shown by a dashed line in figure 19). When $m_b = 2$ solutions are always unstable for the range of E considered. Our calculations indicate that nesting is weakly dependent on ϵ for this case: if ϵ is too small nesting is not observed; however nesting appears to occur for a wide range of ϵ .

Nested solutions may be relevant to the Earth where the magnetic field in the historical record shows four concentrations of intense flux (Jackson, Jonker & Walker 2000), which could be maintained by persistent large-scale fluid downwellings (Willis *et al.* 2007). At low E this is unlikely to occur due to the small scales, which are dynamically preferred, but may be produced by the nests of convection cells found in this paper. Nesting can produce large-scale features and it is possible that the Earth contains two such nests due to a dominant mantle anomaly with $m_b = 2$.

It is interesting to consider the effect of an imposed inhomogeneous heat flux at the OB rather than the imposed fixed temperature considered in this paper. Heat flux boundary conditions tend to favour longer length scales than temperature boundary conditions near the onset of convection (e.g. Glatzmaier & Gilman 1981; Takehiro *et al.* 1999); however the discrepancy disappears at high rotation rates (Gibbons *et al.* 2007). It can therefore be anticipated that nested solutions exist with fixed heat flux at the outer boundary as long as m_b is sufficiently lower than the m_c , as we argue for the case of fixed temperature boundaries.

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