

Convective heat transport in stratified atmospheres at low and high Mach number

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Here we study stratified convection in the context of plane-parallel polytropically stratified atmospheres. We hold the density stratification (n_ρ) and Prandtl number (Pr) constant while studying the dependence of the Nusselt number (Nu), which quantifies the efficiency of convective heat transport, on variations of the Mach number (Ma) and the Rayleigh number (Ra). At large but still subsonic values of Ma, a scaling $Nu \propto Ra^{0.45}$ for low values of Ra. As Ra increases and flows become supersonic, this scaling gives way to a much weaker $Nu \propto Ra^{0.20}$. In the regime of $Ma \ll 1$, a consistent $Nu \propto Ra^{0.3}$ is retrieved, reminiscent of the $Nu \propto Ra^{2/7}$ seen in Rayleigh-Bénard convection.

INTRODUCTION

Convection is essential to heat transport in the cores of high mass stars, the envelopes of low mass stars, and the atmospheres of terrestrial and jovian planets. In such systems, convection occurs in the presence of the atmospheric stratification, which can be small but extends up to 14 density scale heights in the Sun’s convective envelope. A basic understanding of the properties of compressible convection in stratified media is important to understanding systems in astrophysics and planetary sciences. Numerical constraints typically have restricted studies of stratified convection to moderately high Mach numbers, appropriate to regions near the Sun’s surface. As such, we know little about the fundamental properties of low-Mach number stratified convection, which occurs in the deep solar interior.

Some of the earliest numerical experiments on stratified convection were performed in two [1–4] and three [5, 6] dimensions and revealed a number of basic properties in the moderate-to-high Mach number regime. In the widely-studied Rayleigh-Bénard (hereafter RB) problem, upflows and downflows are symmetrical and the temperature gradient approaches zero in the convective interior causing the conductive flux to similarly disappear. Highly stratified convection, on the other hand, exhibits narrow downflow lanes broad upflow regions. Furthermore, the *entropy* gradient is flattened by convection rather than the temperature gradient, which means that even in the presence of perfectly efficient convection a significant component of the flux can be carried by conduction.

In RB convection, there exist two primary control parameters: the Rayleigh number (Ra), the ratio of buoyant driving to diffusive damping, and the Prandtl number (Pr), the ratio of viscous to thermal diffusivity. Along with the aspect ratio of the physical domain, these two numbers entirely control the dynamics of the convection. In stratified atmospheres, in addition to specifying the equation of state and fundamental properties of the gas, the two control parameters of RB convection are joined by the degree of stratification across the domain and the

characteristic Mach number (Ma) of the convective flows. Polytropically stratified atmospheres, such as those used in early studies, are an ideal extension of RB convection into the stratified realm as the two additional control parameters are directly linked to basic properties of the atmosphere. The density stratification is set by the number of density scale heights the atmosphere spans (n_ρ), and Ma is controlled by the superadiabatic excess (ϵ), the deviation of the polytropic index from the adiabatic polytropic index [1].

In this letter we study the behavior of convective heat transport, quantified by the Nusselt number (Nu), in polytropically stratified atmospheres. We vary ϵ and Ra while holding n_ρ and Pr constant across all simulations and convective structures which arise in a domain with an aspect ratio of 4. In section II, we describe the construction of atmospheres, our equations, and our method for numerical time evolution. We describe our findings in section III and discuss implications in section IV.

EXPERIMENT

In order to compare our results with previous studies and in an effort to examine a simplest case, we study a fluid composed of monatomic ideal gas particles, with an adiabatic index of $\gamma = 5/3$ and whose equation of state is $P = R^* \rho T$. The initial stratification is polytropic and the gravitational acceleration and conductive heat flux are invariant throughout the depth of the atmosphere. We specify that both the thermal conductivity, κ , and the temperature gradient, ∇T_0 , are constant while constructing our atmosphere such that $\mathbf{F}_{\text{cond},0} = -\kappa \nabla T_0 = \text{constant}$. Under these assumptions, solving the equation of hydrostatic equilibrium produces an atmosphere defined by

$$\begin{aligned}\rho_0(z) &= \rho_{00}(z_0 - z)^m \\ T_0(z) &= T_{00}(z_0 - z),\end{aligned}\tag{1}$$

where z increases upwards within the bounds $z = \{0, L_z\}$. We specify the number of density scale heights

the atmosphere spans, n_ρ , to determine L_z . Throughout this letter, we set $n_\rho = 3$ such that the density at the bottom of the atmosphere is larger than at the top by a factor of 20. Thermodynamic variables are nondimensionalized at the top of the atmosphere as $P_0(L_z) = \rho_0(L_z) = T_0(L_z) = 1$, requiring $z_0 \equiv L_z + 1$ and $R^* = T_{00} = \rho_{00} = 1$. The polytropic index is set by the adiabatic index and the superadiabatic excess, ϵ , such that $m = m_{ad} - \epsilon$ where $m_{ad} \equiv (\gamma - 1)^{-1}$ is the adiabatic polytropic index. The subsequent entropy gradient at the top of the atmosphere is $\nabla S(L_z) = -\epsilon$. The characteristic timescale of such an atmosphere is related to the characteristic atmospheric buoyancy time, $t_b = \sqrt{L_z/g\epsilon}$. We will utilize buoyancy time units throughout this letter.

Atmospheric diffusivities are set by the Rayleigh number and the Prandtl number. We define the non-dimensional Rayleigh number as

$$\text{Ra} = \frac{gL_z^3(\Delta S_0/c_P)}{\nu\chi}, \quad (2)$$

where ΔS_0 is the entropy difference between the top and bottom of the atmosphere, ν is the kinematic viscosity (viscous diffusivity), and χ is the thermal diffusivity. The relationship between the thermal and viscous diffusivities is set by the Prandtl number, $\text{Pr} = \nu/\chi$. We relate the dynamic viscosity, μ , and the thermal conductivity, κ , to their corresponding diffusivities such that $\nu \equiv \mu/\rho$ and $\chi \equiv \kappa/\rho$. As a result, $\text{Ra} \propto (\nu\chi)^{-1} \propto \rho^2$, such that for our atmospheres with $n_\rho = 3$, the Rayleigh number increases by a factor of 400 across the domain. This formulation leaves Pr constant throughout the depth of the atmosphere, and in this letter we impose $\text{Pr} = 1$.

At the constant values of n_ρ and Pr used, the primary control parameters of convection are ϵ and Ra . We decompose our atmosphere into the background polytrope ($\ln \rho_0, T_0$) and the fluctuations about that background ($\mathbf{u}, \ln \rho_1, T_1$). The scaling of the entropy gradient with ϵ is reflected in the evolved values of these fluctuations, which follow the scaling of $\text{Ma}^{1/2} \propto T_1/T_0 \propto \rho_1/\rho_0 \propto \epsilon$, and which scale as roughly $\text{Ra}^{0.25}$, as in Fig. 1.

We evolve the Fully Compressible Navier-Stokes equations, which take the form:

$$\frac{D \ln \rho}{Dt} = -\nabla \cdot (\mathbf{u}) \quad (3)$$

$$\begin{aligned} \frac{Du}{Dt} &= -\nabla T - T\nabla \ln \rho + \mathbf{g} \\ &\quad - \nu \nabla \cdot (\bar{\boldsymbol{\sigma}}) - \bar{\boldsymbol{\sigma}} \cdot \nabla \nu - \nu \bar{\boldsymbol{\sigma}} \cdot \nabla \ln \rho \end{aligned} \quad (4)$$

$$\begin{aligned} \frac{DT}{Dt} &= -(\gamma - 1)T\nabla \cdot (\mathbf{u}) + \frac{\chi}{c_V} \nabla^2 T + \\ &\quad \nabla T \cdot \nabla \chi + \chi \nabla T \cdot \nabla \ln \rho + \nu (\bar{\boldsymbol{\sigma}} \cdot \nabla) \cdot \mathbf{u} \end{aligned} \quad (5)$$

where $D/Dt \equiv \partial_t + \mathbf{u} \cdot \nabla$ and the viscous stress tensor is

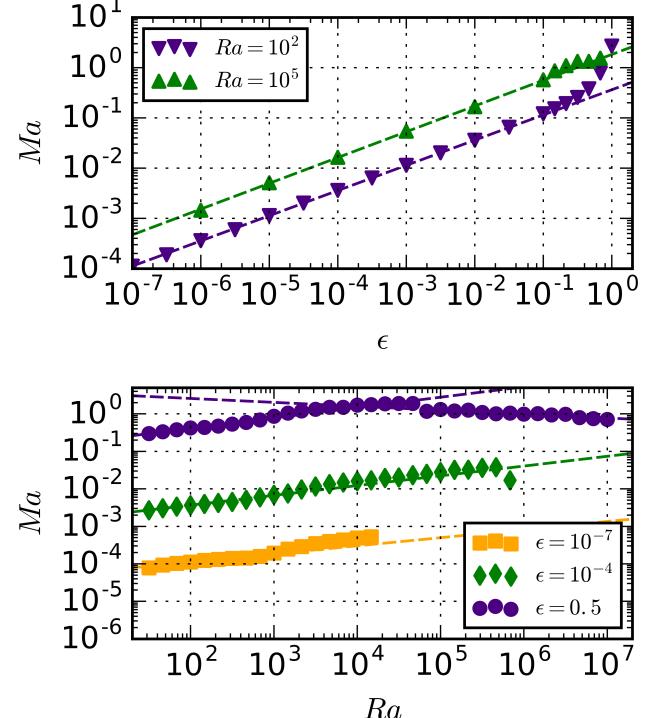


FIG. 1. (a) Characteristic horizontally averaged maximum Mach numbers which have been time averaged for over $100t_b$ starting roughly $50t_b$ after the start of simulations. For $\epsilon \leq 0.1$, an evident scaling of $\text{Ma} \propto \{\epsilon^{0.5015}, \epsilon^{0.5130}\}$ at $\text{Ra} = \{10^2, 10^5\}$ is retrieved. When $\epsilon \rightarrow m_{ad}$, large deviations from this power law are seen and the system quickly approaches the supersonic regime. (b) At high epsilon, Ma scales as $\text{Ra}^{0.2606}$ until it reaches the supersonic regime, at which point it turns over and follows a power law of $\text{Ra}^{-0.1042}$. At low epsilon, consistent power laws are achieved throughout all values of Ra studied, where $\text{Ma} \propto \{\text{Ra}^{0.2606}, \text{Ra}^{0.2165}\}$ for $\epsilon = \{10^{-4}, 10^{-7}\}$. All error bars are negligible. To study dynamics in the regime of deep solar convection ($\text{Ma} \approx 10^{-4}, \text{Ra} \approx 10^{20}$), a superadiabatic excess of $\epsilon \approx 10^{-17}$ is appropriate.

defined as

$$\sigma_{ij} \equiv \left(\frac{\partial u_i}{\partial x_j} + \frac{\partial u_j}{\partial x_i} - \frac{2}{3} \delta_{ij} \nabla \cdot (\mathbf{u}) \right). \quad (6)$$

In such stratified systems, the total convective flux can be defined as

$$\mathbf{F}_{\text{conv}} \equiv \mathbf{F}_{\text{enth}} + \mathbf{F}_{\text{KE}} + \mathbf{F}_{\text{PE}} + \mathbf{F}_{\text{visc}}, \quad (7)$$

where $\mathbf{F}_{\text{enth}} \equiv \rho \mathbf{u} (c_V T + P/\rho)$ is the enthalpy flux, $\mathbf{F}_{\text{KE}} \equiv \rho |\mathbf{u}|^2 \mathbf{u}$ is the kinetic energy flux, $\mathbf{F}_{\text{PE}} \equiv \rho \mathbf{u} \phi$ is the potential energy flux (with $\phi \equiv -gz$), and $\mathbf{F}_{\text{visc}} \equiv -\rho \nu \mathbf{u} \cdot \bar{\boldsymbol{\sigma}}$ is the viscous flux. Taking an inner product of Eq. 4 with \mathbf{u} and adding it to Eq. 5, the full energy equation in conservation form is retrieved,

$$\frac{\partial}{\partial t} \left(\rho \left[\frac{|\mathbf{u}|^2}{2} + c_V T + \phi \right] \right) + \nabla \cdot (\mathbf{F}_{\text{conv}} + \mathbf{F}_{\text{rad}}) = 0 \quad (8)$$

where $\mathbf{F}_{\text{rad}} = -\kappa \nabla T$. An understanding of the flux terms is essential to characterizing the convective heat transport in our systems.

The atmosphere is bounded above and below by impenetrable, stress free, fixed temperature boundary conditions such that

$$w = \partial_z u = T_1 = 0 \quad (9)$$

at the boundaries.

Using a linear eigenvalue solve, we found the critical Ra number to be blahblahblah

We utilize the novel Dedalus pseudospectral framework (<http://dedalus-project.org/>) to time-evolve Eqs. 3-5 using an implicit-explicit third-order four-step Runge-Kutta timestepping scheme [7]. Variables are time-evolved on a dealiased Chebyshev (vertical) and Fourier (horizontal, periodic) domain in which the physical grid dimensions are 3/2 the size of the coefficient grid. Our physical grid sizes range from 96x384 grid points at the lowest values of Ra to 1152x4608 grid points at $\text{Ra} \geq 10^7$. By using IMEX timestepping, we are able to study flows at moderate ($Ma \approx 1$) and very low ($Ma \approx 10^{-4}$) Mach number (Fig. 1b).

RESULTS

The efficiency of convection is quantified by the Nusselt number. While the Nusselt number is well-defined in RB convection as the amount of total flux divided by the steady-state background conductive flux [8, 9], a well-defined Nusselt number is more elusive in stratified convection. A traditional definition of the Nusselt number in stratified convection is [1, 3]

$$Nu \equiv \frac{F_{\text{conv}, z} + F_{\text{rad}, z} - F_A}{F_{\text{ref}} - F_A}, \quad (10)$$

where $F_{\text{conv}, z}$ and $F_{\text{rad}, z}$ are the z-components of \mathbf{F}_{conv} and \mathbf{F}_{rad} , respectively. F_A is the adiabatic conductive flux, defined as $F_A = -\kappa \partial_z T_{\text{ad}}$. For an ideal gas atmosphere in hydrostatic equilibrium, $\partial_z T_{\text{ad}} \equiv -g/c_P$. $F_{\text{ref}} = \Delta T/L_z$, where $\Delta T = T_u - T_\ell$, is the conductive flux of a linear profile connecting the upper and lower plates.

We contend that this is the general form of the Nusselt number. To illustrate this, we consider a few limiting cases. Convection works to suppress entropy stratification and create isentropic atmospheres. Under the Boussinesq approximation, density variations are ignored and entropy stratification is directly proportional to temperature stratification, such that $\nabla S \rightarrow 0$ only when $\nabla T \rightarrow 0$. Thus, for RB convection, $\nabla T_{\text{ad}} \equiv 0$ and the familiar form of Nu is retrieved. In the case of stratified convection, as $\epsilon \rightarrow m_{ad} + 1$, $\nabla P \propto g \rightarrow 0$ and the resulting $\nabla T_{\text{ad}} \rightarrow 0$. In such a case, $F_A \rightarrow 0$ and the familiar

definition of the RB Nusselt number is appropriate to use, as convection carries all of the atmospheric flux in such a case [10]. As $\epsilon \rightarrow 0$, $\nabla T_{\text{ad}} \rightarrow \nabla T_0$, the atmospheric initial temperature profile. This causes increasingly smaller velocity, and thermodynamic perturbations (e.g. Fig. 1), but the removal of the $O(1)$ background term in the numerator and denominator of Eq. 10, makes the numerator and denominator both $O(\epsilon)$.

We solved initial value problems which start in hydrostatic and thermal equilibrium and experienced infinitesimal kicks compared to ϵ in T_1 . Solutions were time-evolved until a long-time average of Nu showed little dependence on height. At high values of ϵ , shock systems form in the upper atmosphere near downflow lanes (see e.g. Fig. 2a) and propagate towards upflow lanes. Such systems were reported in both two [4] and three [11] dimensional polytropic simulations previously. It is unsurprising that [4] discovered that as $\gamma \rightarrow 1$, shocks are easier to find and produce, as in such a regime $m_{ad} \rightarrow \infty$ and, as a result, $\epsilon \rightarrow \infty$. Low mach number flows, such as those in an $\epsilon = 10^{-4}$ atmosphere (e.g. Fig. 2b) have similar bulk thermodynamic structure but lack the complicating dynamics of shock heating. This is interesting in light of the current literature, as low Mach number flows are in pressure equilibrium with the background and so pressure forces can no longer be the cause of narrow downflow lanes and broad upflow regions, as suggested by [3]. As Ra is increased to very large values (e.g. Fig. 2c), thermodynamic structures break up into small packets which traverse through the domain many times before diffusing rather than spanning the whole domain. The complicated nature of high Ra dynamics, especially in the low Ma regime where shocks are absent, has barred us from sufficiently converging any solutions in the regime of $Ra > 10^6$ at low ϵ .

At high values of Ra, the heat transport properties of the systems become increasingly complex and time-dependent. Large ϵ flows exhibit two local maxima in the enthalpy flux and kinetic energy flux: one in the upper atmosphere caused by the shocks, and one in the lower atmosphere caused by the deep mixing of convective motions. At low ϵ , only the deep maximum is present. Furthermore, for low values of ϵ , the presence of fixed-temperature boundary conditions allows the flux at the boundaries to vary. Many runs at $Ra > 10^5$, $\epsilon = 10^{-4}$ exhibit time-dependent states of stratified convection (such as those shown in Fig. 2), in which the flux entering the system at the bottom of the atmosphere exceeds that at the top. These states are punctuated by states of vigorous shearing which flatten the bottom temperature gradient towards adiabatic, allowing the excess energy to exit through the upper boundary. These shearing states will be covered in more detail in a future paper. Regardless, a proper long-term average over shearing and non-shearing states retrieves an invariant flux profile (and therefore Nu profile) throughout the depth of the atmosphere (After

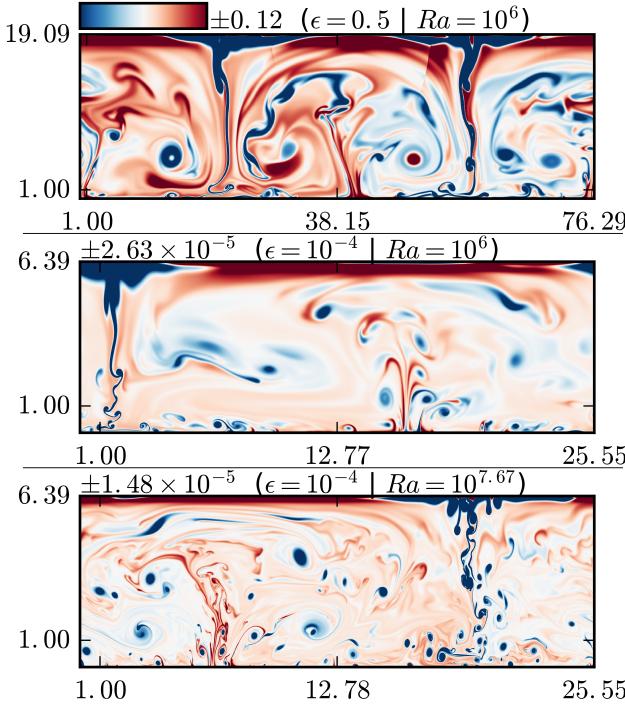


FIG. 2. Characteristic entropy. The time- and horizontally-averaged profile is removed in all cases. At high ϵ (a), shock systems form near the upper downflow lanes and propel shock-heated material deep within the atmosphere at sufficiently high Ra . At low ϵ but at the same Ra (b), shock systems are absent, but otherwise the dynamics are similar. As Ra is increased (c), downflow lanes no longer span the entirety of the domain and individual small blobs are responsible for carrying the flux.

about 24 more hours, it'll be flat! I don't have enough data at the parameters shown in Fig. ?? yet.

After appropriately time-averaging the fluxes, a sensible Nusselt number is retrieved. Nusselt numbers for all simulations at low and high Ma are plotted in Fig. 4. At $\epsilon = \{10^{-4}, 10^{-7}\}$, scaling laws of $Nu \propto Nu^{0.3090, 0.3087}$ are retrieved. At $\epsilon = 0.5$, in the near-sonic regime ($Ra \leq 10^4$), the scaling of Nu with Ra is inflated, with $Nu \propto Ra^{0.4489}$. As simulations pass into the supersonic regime and shocks start to form near the downflows, propelling warm fluid deep into the atmosphere, that scaling drops to $Nu \propto Ra^{0.1936}$. Error on all power law scalings is negligible.

DISCUSSION

Here we have presented a non-dimensional construction of polytropic atmospheres and the control parameters which determine flow characteristics in such systems. We argue that such systems are the natural extension of

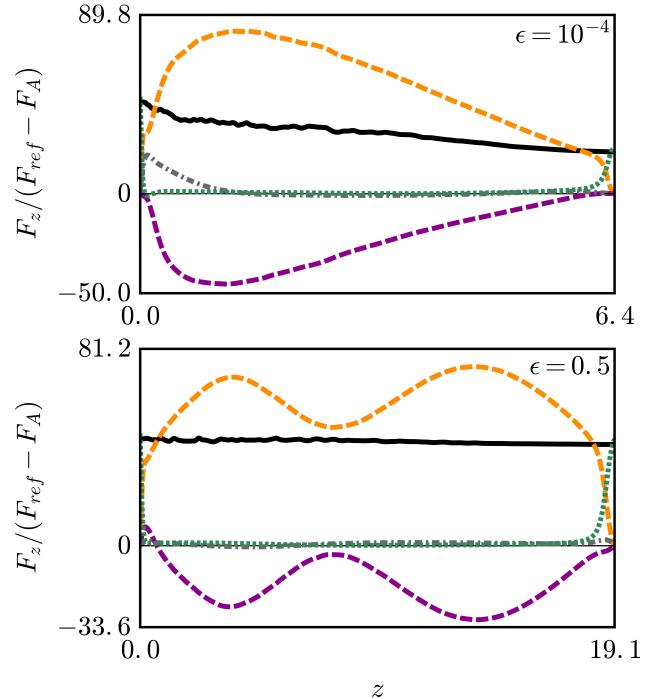


FIG. 3. Time-averaged flux profiles for low (top) and high (bottom) Mach number flows at $Ra = 10^6$. The dashed lines correspond to the enthalpy flux (orange, positive) and kinetic energy flux (purple, negative). The grey dash-dot line is the viscous flux, and the green dotted line is the radiative flux with the adiabatic radiative flux removed. All fluxes are normalized by $F_{ref} - F_A$, as in the calculation of the Nusselt number. The solid black line is the properly normalized sum of all the fluxes, and under this normalization its height-averaged value is the Nusselt number.

the RB problem to stratified systems, and should be used to understand the basic properties of stratified convection. We have argued that the traditional definition of the Nusselt number in stratified atmospheres [1] is the correct one, and shown how it extends in extreme cases.

Furthermore, we have demonstrated that low Ma flows exhibit the same broad-upflow / narrow downflow dynamics as are observed in high Mach number flows [3]. However, unlike at high Ma, pressure perturbations are no longer responsible this effect; rather, high temperature, low density upflows and low temperature, high density upflows causing appropriately signed entropy perturbations are responsible for the entirely buoyantly-driven flow. It is possible that at low Ma, these flows begin to feel the pressure response of the top and bottom walls which leads them to deflect at the far boundary in a way that high Ma flows do not.

The dynamics of these polytropic solutions are complex and highly time-dependent, even in two dimensions. Time-dependent oscillating shear states have developed spontaneously, as seen before in RB convection

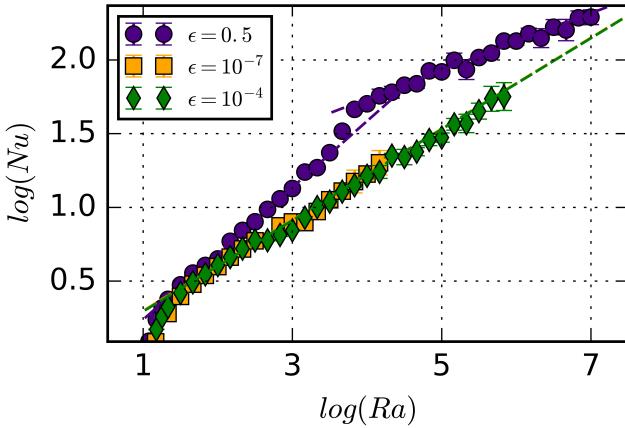


FIG. 4. Variation of Nu as Ra increases is shown for $\epsilon = \{10^{-7}, 10^{-4}, 0.5\}$. At high ϵ , a clear transition from the subsonic to supersonic regime is evident in the scaling of Nu with Ra . In the low ϵ regime Nu vs. Ra scalings collapse onto a similar line which is reminiscent of RB scalings [8].

(citation). Future work will aim to better understand the mechanisms of shearing states and whether or not these states are attainable in three-dimensional, non-rotating atmospheres. While the two-dimensional work studied here offers a basis for comparing the heat transport abilities of stratified convection compared to RB convection [8], it will be informative to study the low Mach number systems presented here in three dimensions

and in atmospheres bounded by stable regions.

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