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Is the 670 km phase transition able to layer the Earth's convection in a mantle with depth-dependent viscosity?

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Abstract The effect of a viscosity stratification on phase change dynamics have been investigated with axi-spherical convection models. As in previous studies with a constant viscosity mantle an intermittent layering appears for a Clapeyron slope from -2 MPa/K to -3 MPa/K. A viscosity increase in lower mantle requires a more negative Clapeyron slope to produce the layering. This shift is sensitive to the mechanical boundary condition. With a viscosity contrast of 30, a no-slip top condition does not lead to layering in the range of the possible values for the Clapeyron slope. With a free-slip condition, the threshold is at -4 MPa/K. Just below this threshold, a whole mantle circulation driven by a cylindrical hot plume coexists with layered mantle domains over several billion years.

Introduction

For a long time, the seismic discontinuities observed inside the Earth's mantle have been considered to be solid phase changes. The 670 km discontinuity is particularly important because it is associated with an endothermic phase change which can break the convection into two layers. Numerous numerical experiments of mantle convection were devoted to this phenomenon. Some of them, mostly done in two-dimensional Cartesian geometry, focused on basic effects such as the influence of the Rayleigh number [Christensen and Yuen, 1985], the form of governing equations [Ita and King, 1994] or the variations of mantle properties [Steinbach and Yuen, 1994]. Other models attempted to model mantle convection more directly by introducing a spherical geometry and considering phase changes at their actual depth [Machetel and Weber, 1991; Tackley *et al.*, 1993; Solheim and Peltier, 1994]. In these cases, an intermittent behavior between a one layer and a two layer convection occurs for Clapeyron slopes in the range of -2 MPa/K to -4 MPa/K. This transitional mode occurs as cold instabilities accumulate above the phase transition before their rapid discharge into the lower mantle by an avalanche-like effect. Experimental estimates of the Clapeyron slope of the transition from γ -spinel into perovskite+magnesiostite, range from -2 MPa/K to -3 MPa/K [e.g. Bina and Helffrich, 1994]. Accordingly, this pattern may be expected in the Earth.

Among the spherical models mentioned above, only Tackley *et al.* [1994] take the depth dependence of viscosity into account, but they do not systematically study its influence. Furthermore, in this model, the viscosity varies smoothly over one order of magnitude without a step at the phase change. In reality, viscosity depends on temperature, pressure and shear stress, but probably also on the mineral structure. The aim of the present paper is to study the influence of viscosity stratification caused by phase changes. A

comparison between axisymmetrical and 3-D spherical cases with constant viscosity by Machetel *et al* [1995] shows that the threshold for layering occurs for the same Clapeyron slope in both situations. This result is somewhat surprising in view of cartesian experiments where the mass flux across the phase change interface varies slightly between the 2-D and 3-D cases [Yuen *et al*, 1994]. This is explained by the fact that the axisymmetrical geometry offers the possibility to test the stability of the stratification to 3-D waves triggered along the polar axis and to 2-D ones elsewhere. Accordingly, we conduct this study in axisymmetrical geometry which allows to dramatically increase our computing speed to follow the evolution of the convective circulation with sharp phase changes over period of time several times longer than the Earth's age.

Formalism

We employ a set of convection equations which corresponds to the framework of anelastic compressible fluids with infinite Prandtl number described by Jarvis and McKenzie [1980], extended to include depth-dependent thermal expansivity coefficient α and to depth dependent bulk modulus K_s .

The equation of state

The linearized expression of density as a function of the temperature T and the pressure P is:

$$\rho(T, P) = \rho_r(1 - \alpha(T - T_s) + K_s^{-1}(P - P_h)) \\ + \Delta\rho_{400}(\Gamma_{400}(T, P) - \Gamma_{400}(T_{400}, P)) \\ + \Delta\rho_{670}(\Gamma_{670}(T, P) - \Gamma_{670}(T_{670}, P)) \quad (1)$$

Here, T_s and P_h stand for the adiabatic temperature and the hydrostatic pressure respectively; $\Delta\rho_{400}$ (180.53 kg/m^3) and $\Delta\rho_{670}$ (388.57 kg/m^3) are the density jumps given by P.R.E.M.; and $\Gamma(P, T)$ represents the phase state function. Its analytical expression is :

$$\Gamma_i = \frac{1}{2} \left(1 + th\left(\frac{P - P_i - \gamma_i(T - T_i)}{\delta P} \right) \right), \quad (2)$$

where γ_i is the Clapeyron slope, P_i and T_i the pressure and the temperature at the reference depth ($P_{400}=13.1 \cdot 10^9 \text{ Pa}$, $T_{400}=1725 \text{ K}$, [Katsura and Ito, 1989], $P_{670}=23.8 \cdot 10^9 \text{ Pa}$, $T_{670}=1733 \text{ K}$, [Ito and Takahashi, 1989]). Here, δP is the half-width of the phase loop, assumed to be 0.4 GPa (about 10 km thick).

To build the reference profiles $\alpha(r)$, $\rho_r(r)$ and $K_s(r)$, we first took an exponential form for ρ_r and K_s which fits the P.R.E.M. [Diewonsky and Anderson, 1981] profiles with the smoothed discontinuities at 400 km and 670 km. To these analytical expressions, we then add the PREM density and compressibility steps at the discontinuities. As in Glatzmaier [1988], the Gruneisen parameter is assumed to be constant and equal to 1. This allows to

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relate the radial dependence of α to $\rho_r(r)$ and $K_s(r)$, such as α increases from 10^{-5} K^{-1} at the CMB to $4 \cdot 10^{-5} \text{ K}^{-1}$ at the surface.

The momentum equation

Because of the axisymmetrical geometry, the velocity field can be deduced from a poloidal scalar potential field. Taking the curl of the momentum conservation equation, we obtain a fourth-order elliptical equation relating the scalar potential to the temperature. The radial variation of the viscosity is simulated using two constant-viscosity layers, within which the continuity of both the velocity field and the vertical stress are prescribed [Cserepes *et al.*, 1988]. We restrict the convective domain to a half sphere. The solution is obtained by a Legendre polynomial horizontal decomposition up to the 512th harmonic, and by a vertical finite-difference decomposition on 600 regular spaced radial levels.

The energy conservation equation

We handle the latent heat release using the formalism of *Christensen and Yuen* [1985] to write the energy conservation equation. To solve this equation, we use a control volume method [Patankar, 1980] on a regular grid consisting of 600 radial and 256 horizontal points under isothermal boundary conditions. A reflecting condition is applied at the equator. Comparisons with analytical and numerical published results validate the computations [Machetel and Weber, 1991; Solheim and Peltier, 1994]. In addition, we internally check computations for global energy conservation. We find that the energy conservation equation integrated over the whole fluid domain is satisfied to within about 0.05%.

Models

Modeling of the global convection can follow two contrasting approaches. Either the rheology of the lithosphere is neglected and the top boundary is stress free with a 0°C temperature condition [e.g. Solheim and Peltier, 1994], or the top boundary is taken as the base of the lithosphere [e.g. Machetel and Weber, 1991]. In the latter case, a model simulates a possible secondary mode of convection unrelated to plate motion. So the top boundary of the mantle model may be seen as the base of the mechanical lithosphere. This situation corresponds to a strain rate threshold at a sharp temperature interval around 900°C, and to a rigid condition. Both approaches account for some basic aspects of the global mantle convection. Hence we have considered both situations, a whole and sublithospheric convection, hereafter referenced as model A and B, respectively. We set the temperature of the CMB to 3000 K in both models and perform models A and B with an internal heating rate of 20 TW and of 5 TW, respectively.

Figure 1 presents characteristic snapshots of the flow and temperature fields of six experiments. The Clapeyron slope of the transition at 400 km depth is set to 3 MPa/K [Katsura and Ito, 1989]. The experiments commence with a thermal field calculated from a previous model, except for the first one, where the initial condition is a conductive profile with a small random perturbation. Each experiment ran several billion years (some 10^5 time steps) until reaching a quasi-steady regime, when input heat sources balance the output heat flow.

Constant viscosity cases

In Models A1 and B1, the Clapeyron slope at the 670 km discontinuity is set to -3 MPa/K. As in the Solheim and Peltier's [1994] experiments the convective flow is partially stratified:

closed streamlines delimit both domains restricted to the upper mantle and domains encompassing the whole mantle. A weak thermal boundary layer occurs at the depth of the endothermic phase change, as it is apparent on the geotherm. Most of cold downwellings, which result from instabilities of the upper boundary layer, remain confined to the upper mantle. Nevertheless, the strongest ones pass episodically into the lower mantle, but their associated upward return flow remains diffuse and the hot plumes cannot cross the phase change. This pattern is reminiscent of the avalanche effect depicted by Machetel and Weber [1991] and Solheim and Peltier [1994]. The rate of heat advected at the surface by model A1 and B1 is 50 TW and 15 TW, respectively, and the velocity of sinking flows during avalanche events are of the order of 30 cm/year and 10 cm/year, respectively.

With a more negative Clapeyron slope of -4 MPa/K in cases A2 and B2, the stratification becomes quite complete and forms a large thermal boundary layer at 670 km depth. The step in temperature reaches 750 K in model A2 and 540 K in model B2 (i.e. 30% of the temperature drop across the entire depth in both cases). Here, none of the instabilities initiated at the top boundary and the core-mantle boundary cross the phase change. Nonetheless, this restriction does not preclude significant material exchange between the upper and lower mantle. On snapshot A2, we observe the simultaneous onset of cold and hot instabilities within the thermal boundary layer of the phase change. The development of these instabilities leads to a rapid discharge of the upper mantle into the lower mantle (not shown). This phenomenon has already been described by Steinbach *et al.* [1993]. It happens in a very short time (50 My) compared to its recurrence interval (around 0.6 By). So, we have to distinguish a partially stratified convection, as in case A1, where the endothermic phase change imposes a weak resistance against a whole mantle convection, from an intermittent circulation, as in case A2, where little exchange occurs between layers except during catastrophic mantle overturns. This difference is apparent in the difference of the instantaneous thermal profile (in black) and the temporally averaged one (in red). In case A1, their similarity indicates that the mean thermal state of the mantle is not altered by the avalanches, while case A2 shows clear fluctuations over time. In the no-slip experiment B2 no such an overturn occurred in 10 By. In B2 the mantle stratification seems very stable in the temperature profiles. The contrast between model A2 and B2 shows that the layering induced by endothermic phase change is also sensitive to the boundary conditions, as previously shown by Ita and King [1994].

Cases with a stratified viscosity

Models A3 and B3 show the flow pattern when the Clapeyron slope is strongly negative, -4 MPa/k. The lower mantle is now 10 times more viscous than in the constant-viscosity cases and the upper mantle 3 times less viscous, corresponding to a viscosity contrast of 30 across the 670 km discontinuity. This viscosity difference lies in the range of acceptable values to explain the relationships between the geoid and mantle tomography at low harmonic degrees [e.g. Ricard *et al.*, 1989]. In both cases, we start the calculation with the previous isoviscous and partially layered models, and observe the layering collapsed in a very short time. In contrast to models A1 and B1, only one hot plume at the polar axis crosses the whole mantle, and the downwelling return flow is now diffuse. Clearly the whole circulation is driven by a strong upwelling sweeping away cold instabilities which seem unable to penetrate into the lower mantle. Indeed, here the Rayleigh number based on the lower mantle viscosity is one order of magnitude

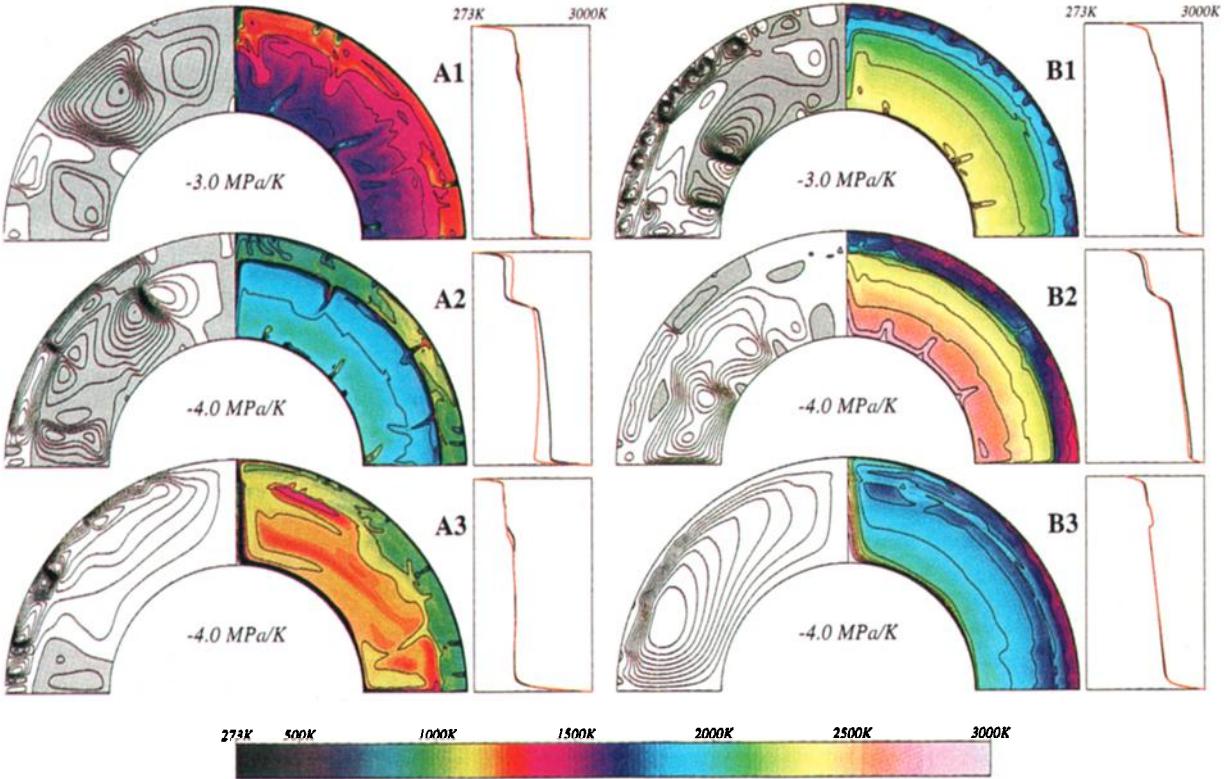


Figure 1. Snapshots of six experiments drawing the stream function, the temperature field and the horizontally averaged temperature profile (in black the instantaneous one, in red the timely averaged over the last 30 000 timesteps, i.e. at least .5 By.). Cases A are run with a free slip and a 0°C top condition, and B with a no slip and a 900°C top condition. The internal heating rate is set to 20 TW in cases A and 5 TW in cases B. Cases 1 and 2 are viscosity constant models with a viscosity set to 10^{21} Pas. In case 3 the viscosity is set to 10^{22} Pas below 670 km and $3 \cdot 10^{20}$ Pas above.

smaller than in experiments A1 and B1. This modification might explain the observed whole mantle convective pattern since layering tends to grows with increasing vigor of convection [Christensen and Yuen, 1985; Steinbach *et al* 1993]. However, the viscosity profile leads to convective flows which are just as efficient as those developed in cases A1 and B1. The rate of heat advected by convection in cases A3 and B3, 45 TW and 13 TW respectively, remain similar to those yielded by cases A1 and B1. Also, in the ascending plume the velocity ranges from 10 to 15 cm/y in the high-viscosity layer and reaches a maximum of 40 cm/y in the upper low-viscosity layer.

Intuitively, one expects the increase of viscosity with depth to reinforce the layering. However, the opposite holds. A similar paradoxical behavior was found by Hansen and Yuen [1994]: the decrease of thermal expansivity helps penetrate a compositional barrier. They explain this observation by the ability of depth-dependent properties to produce focused plumes and thus predict an equivalent result with depth-dependent viscosity. Actually, one of the most spectacular effects of the viscosity increase with depth is to significantly cool the mantle [Gurnis and Davies, 1986; Hansen *et al.*, 1993]. The average temperature at mid-depth is 200 K lower in cases A3 and B3 than in cases A1 and B1. The cooling enhances the temperature contrast between the adiabatic mantle and the core of hot plumes. On the one hand, the increased temperature contrast enhances the resistive forces at the phase boundary. On the other hand it allows strong upwellings to develop over the lower mantle thickness, i.e. 4 times the height of cold downwelling in the upper mantle. To trigger an avalanche cold instabilities gather up and accumulate enough buoyancy to overcome the stabilizing effect of the phase change, as documented by Machetel and Weber [1991] and Weinstein [1993]. In case of hot plumes ascending from the CMB, this accumulation of buoyancy naturally results from their great vertical extent. Thus

no transition period is needed and the whole mantle circulation observed in cases A3 and B3 persists over more than 2 By and 10 By, respectively, even though these plumes pulse because they drag most of the instabilities growing at the CMB. Note that snapshot A3 displays a transient hot elongated sheet ascending along the equatorial plane which is unable to pass into the upper mantle. This observation suggests that only cylindrical plumes can overcome the phase change barrier.

In spite of the similarity between models A3 and B3, some essential differences occur. In the geotherm of model B3, a temperature inversion appears around the endothermic phase transition. It is characteristic of a totally unlayered circulation [Christensen and Yuen, 1985; Solheim and Peltier 1994] and is caused by the latent heat effects. Surprisingly, a small boundary layer appears in geotherm of case A3, indicating a partial stratification despite the persistence of the polar plume crossing the whole mantle. The streamlines of snapshot A3 shows a circulation almost perfectly layered far from the polar plume. Here, the partial layering is not temporal, but spatial. After 2 By, the convection suddenly switches to a stratified regime with mantle overturns occurring every 1.5 By. This may indicate that case A3 is close to bifurcating between a one layer and a two. Accordingly, for a viscosity contrast of 100 (all other parameters and boundary conditions the same as in case A3), we checked that the circulation keeps a one layer regime with a stable heat flow and without any thermal boundary layer at the phase change depth during 4 By. Conversely, with a viscosity contrast of 30, a fully layered convection, stable over several billion years, occurs when the Clapeyron slope is -5 MPa/K. This indicates that the threshold between one layer and two layers is very sharp for a stratified mantle viscosity and that the partially stratified regime pictured in case A1 seems unlikely. The material exchanges between both layers still occur, but only during a few catastrophic avalanches. They are

similar to the ones observed in case A2, but they are less frequent and involve a huge temperature step (around 1000 K) across the phase transition: in 10 By three avalanches leading to a complete mantle reserval are observed. This is in agreement with *Steinbach and Yuen* [1994], who showed on cartesian experiments that the magnitude of avalanches is largest when viscosity increases with depth.

Conclusions

Our constant viscosity models agree with the results obtained by *Solheim and Peltier* [1994] with both an exothermic phase change located at a depth of 400 km and an endothermic one located at a depth of 670 km. We show here that the viscosity increase with depth is likely to shift the threshold value of the Clapeyron slope to more negative values. With a viscosity step by a factor of 30 to 100 across the upper-lower mantle boundary, the threshold would lie in the range of most extreme estimates for the Clapeyron slope of the γ spinel-perovskite transition [*Ito et al.* 1990; *Bina and Helffrich*, 1994]. In this case we observe that the viscosity increase with depth yields a convective pattern dominated by strong cylindrical hot plumes and weak sinking sheets, as previously shown by *Rabinowicz et al* [1990]. These plumes pass more easily through strong endothermic phase transitions than those rising in an isoviscous fluid [*Nakakuki et al.*, 1994; *Schubert et al.*, 1995]. Furthermore, experiments close to the bifurcation threshold depict a new convection picture where a global mode driven by a huge ascending plume is superimposed on a layered one. An equivalent behavior has been observed by *Zhong and Gurnis* [1994] when cold downwellings are stiffened by a temperature-dependent rheology: slab penetration is enhanced without significant alteration of the convective layering. So, it may be expected that both patterns coexist in the Earth, i.e. a stratified mantle where only some plumes and slab cross the interface. Such a geodynamical picture might explain the apparently conflicting message of isotope geochemistry: the simultaneous existence of a shallow depleted reservoir, source of MORB, consistent with a two layer convective mode, and of upwellings deeply rooted in reservoir enriched by a crustal input. Finally, the time scale needed to study mantle convection with phase changes is of the order of the Earth's age: the bifurcation from an unlayered to a layered regime, as well as the stagnation period between two mantle overturns can last several billion years.

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