Elsevier Science Publishers B.V., Amsterdam

Pulsating diapiric flows: Consequences of vertical variations in mantle creep laws

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Received November 21, 1991; revision accepted June 17, 1992

ABSTRACT

Recent laboratory work has suggested that the rheology of the lower mantle may be Newtonian. We have studied the time-dependent dynamics of plumes interacting with a rheological interface separating an upper non-Newtonian mantle and a Newtonian lower mantle. Pulsating diapiric structures with fast time scales are promoted by the interaction of the rising plumes with this rheological boundary. Surface heat flow signals are discernible as pulses, which remain relatively stationary. They correlate well with the localized upwellings just below the surface. The presence of a mobile lithosphere from increasing the non-Newtonian power-law index helps to produce a large-scale circulation in the upper mantle which draws large hot patches away from nearby upwellings. In our calculations the resultant averaged effective viscosity of the non-Newtonian upper mantle is about two orders of magnitude lower than that of the Newtonian lower mantle. A viscously stratified Newtonian model produces more incoherent and broad-scale diapiric structures and is less efficient for generating sharply varying time-dependent thermal signatures.

1. Introduction

Mantle convection has traditionally been studied numerically or experimentally on the assumption that the creep law relating the stress to strain rate does not change in character with depth [1-4]. However recent experiments on perovskite analogues by Karato and Li [5] have opened up the possibility that the rheology of the lower mantle is Newtonian. This finding would, then, suggest vertical variations in rheological laws with depth, as both olivine [6,7] and garnet in the transition zone [8] may have strong non-linear dependence of the strain rate on the stress in the flow law. The dynamics of such types of rheological stratification, a non-Newtonian upper mantle overlaying a Newtonian lower mantle, have recently been examined [9] within the framework of a steady-state model. In this model it was found that concentrated surface streamlines and plate-like characteristics were promoted by this type of rheological stratification, in which the upper mantle had an effective viscosity which was lower than the underlying Newtonian mantle. In this work we will demonstrate that allowing for time-dependence in convection with a rheological layering will bring about some interesting phenomena in the development of pulsating diapiric flows that are not found in models with homogeneous rheology. Hitherto, studies on modelling of plume dynamics have been focussed on a fluid medium with a single flow law with continuously varying parameters, whether it be temperature-dependent [10-12], strictly Newtonian [13,14], non-Newtonian [15,16], or with phase changes [17]. Thus there have been no investigations on the time-dependent behavior of plumes interacting with such a rheological boundary. From physical considerations one would expect a strong nonlinear coupling to be developed at this rheological boundary where the flow law changes abruptly.
from a linear to a nonlinear rheology in the course of the plume's ascent. Also, from the mathematical point of view, the partial differential equations governing the momentum equations change suddenly from a linear to a strongly nonlinear elliptic equation. There are also many circumstances in nature in which the transport property of the medium changes its character all of a sudden, such as radiation passing through composite material with vastly contrasting dielectric constants. Therefore, this problem of time-dependent mantle convection in a rheologically composite medium is very interesting from many viewpoints.

In this work we will present results based on three different models, and show the importance of vertical variations in the creep law parameters on generating pulsating diapiric flow structures. The surface expression of discrete volcanoes along linear volcanic chains suggests that mantle plumes are in the form of diapirs [18] instead of continuous conduits.

2. Model, equations, and methods

The model mantle to be focused on here is an incompressible medium with a non-Newtonian upper mantle overlying a Newtonian lower mantle. The olivine upper mantle is considered to be non-Newtonian because of the presence of dislocations in mantle xenoliths and the existence of seismic anisotropy in the upper mantle [19]. Recent laboratory work on garnet [8] has shown that it also has a rather strong nonlinear rheology. We solve the Boussinesq equations for infinite Prandtl number convection without internal heating. The momentum and temperature equations are given in dimensionless form for temperature $T$, dynamic pressure $p$ and velocity $u$ by:

$$\frac{\partial \tau_{ij}}{\partial t} - Ra T \hat{z} - \nabla p = 0$$

$$\frac{\partial T}{\partial t} = \nabla^2 T - u \cdot \nabla T$$

The dimensionless depth-varying creep law used in the computations is given by:

$$\tau_{ij} = A(z) \epsilon^{1/n(z)} - \epsilon_{ij}$$

where $Ra$ is the Rayleigh number based on the Newtonian viscosity of the lower mantle and the depth of the entire layer, $\tau_{ij}$ and $\epsilon_{ij}$ are the elements of, respectively, the deviatoric stress and strain-rate tensor, $n(z)$ represents the depth variation of the power-law index $n$ in the mantle creep law, $\epsilon$ is the second invariant of the strain-rate tensor, defined in [16], and $T$ is the dimensionless temperature based on the temperature difference imposed across the layer. The important control variables of this problem are $Ra$, $n(z)$ and $A(z)$. It is to be noted that $A(z)$ in eqn. (3) is a dimensionless parameter and is not coupled to $n(z)$. In dimensional units $A(z)$ would

![Fig. 1. Finite-element grid used for capturing diapiric structures. This grid is used for the global snapshot in Fig. 2. Quadratic triangular elements are used for the velocity field and linear triangular elements for the temperature field. There are six grid points per velocity element and three grid points per temperature element. Vertically there are 19 velocity elements in the upper mantle and 18 in the lower mantle. Along $x$ are 14 velocity elements between 0 and 0.8, 24 between 0.8 and 1.5, and 21 between 1.5 and 3.5. Total number of grid points used exceeds 8500.](image)
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depend on \( n(z) \) in eqn. (3) [20]. For this composite rheology we find it computationally convenient to base the Rayleigh number on the linear viscosity of the lower mantle. The power-law index for the lower mantle is taken to be \( n = 1 \), representing a diffusion creep mechanism. For the upper mantle we take either \( n = 1 \) (for the purely Newtonian layered case) or \( n = 3 \), representing a dislocation creep mechanism [20]. In the model with a lithosphere present we use \( n = 5 \) to describe Dorn creep, which is used for describing lithospheric deformation [21].

The coordinates are given by \( x \) and \( z \), with the \( z \)-axis aligned with the gravity vector. Isothermal boundary conditions \( (T = 0 \text{ and } T = 1) \) are imposed at the top \( (z = 0) \) and bottom \( (z = 1) \) respectively. We integrate eqns. (1), (2), (3) in time with a penalty-function, finite-element method.

\[
Ra = 10^4 \quad n = (3,1) \quad A = (1,1)
\]

Fig. 2. Global shots of the stream function and temperature \( T \) for model A. This is a model with a non-Newtonian upper mantle above a Newtonian lower mantle with the same prefactors \( A \). \( Ra \) based on the lower mantle viscosity is \( 10^4 \). The boundary between the upper and lower mantle is at \( z = 0.23 \). The box indicates the section where the zoom snapshots are taken. Time is non-dimensionalized by the thermal diffusion across the layer \( (d^2/\kappa) \). For a model mantle 1800 km deep (accounting for the spherical to cartesian transformation), \( t = 0.001 \) would correspond to 102 Myr. Grey scales are linear and evenly spaced.
Fig. 3. Sequence of zoom-in snapshots of T and \( \psi \) for case A. Time advances upward. The period of the pulsation is around 0.0008, which is shorter than the overturn time of the large-scale flow. Overturn times scale as \( O(d/V_{\text{rms}}) \). For this case is \( O(0.005) \).
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[22], imposing stress-free impermeable boundary conditions along all sides and reflecting boundary conditions at the two vertical boundaries.

A predictor-corrector method [23] is used in advancing the set of time-dependent convection equations with a penalty-function parameter of $10^6$ and with linear temperature and quadratic velocity triangular elements. We have benchmarked for aspect-ratio one this primitive-variable code against two codes [1, 16] based on bi-cubic splines for non-Newtonian ($\mu = 3$) rheology and found good agreement [16, 23]. For this problem with fast diapirc instabilities coming off a rheological boundary, grid refinement is needed at the interface ($z = 0.23$). Figure 1 shows the finite-element grid configuration with aspect ratio 3.5 which is used for monitoring a time-dependent solution, shown in Fig. 2, in which 49 and 59 unevenly spaced velocity elements are used respectively along the vertical and horizontal directions. There are six grid points in one velocity element and three points in one temperature element, because of the different order [23]. The vertical column with a dense grid is moved along $x$ adaptively by the user in order to follow and capture the diapiric structure rising from the central plume shown in Fig. 2.

These calculations have been conducted for an aspect ratio of 3.5 in order for the plumes not to be influenced by edge effects. In contrast to Newtonian convection the computational time for these types of mixed Newtonian–non-Newtonian fluids is much longer than for simple Newtonian fluids, as the viscosity distribution changes with time, and subiteration is needed at each time step. The cost in CPU time is at least five times greater than for the Newtonian fluids (see case H below in Figs. 11–16). We note that it is essential to examine the dynamics of this type of flow first in two-dimensional configurations, which allow much better—and necessary—resolution for the diapiric regimes than is now feasible in 3D [24]. Moreover, the spatial and temporal resolutions needed for monitoring sharp diapirc instabilities in 3D would impose a severe restriction on an extensive investigation of strongly convecting regimes in which a new physical mechanism has to be understood.

In contrast to the study of steady-state situations [9], visualization of the time-dependent events on graphics workstations plays an indispensable role in this investigation. Many of the phenomena shown below (see for example, Figs. 2 and 3) may not have been perceived had we not employed visualization techniques to monitor in detail the temporal development of these flows.

3. Results

In this section we present results taken from large-scale numerical simulations of three cases. These computations are lengthy and require over 100 hours of supercomputer (CRAY-1MP) time. The first is called model A, which consists of a non-Newtonian upper mantle above a Newtonian lower mantle. The second model, model C, is similar to model A, but with a lithosphere. The third, model H, has a viscously stratified Newtonian rheology throughout.

3.1 Model A (non-Newtonian upper mantle and Newtonian lower mantle)

In this first case, A, we will examine a two-layer model with $\mu(z) = (3, 1)$ for the upper and lower mantle, respectively; the pre-factors are $A = (1, 1)$ for the two layers and the Rayleigh number is $10^4$. The initial condition is taken from a steady-state solution [9] together with a perturbation in $T$ with an amplitude of 0.05 and fundamental sinusoidal wavelengths along both the $x$ and $z$ directions.

Figure 2 shows a global snapshot at diffusional $t = 0.0164$, which takes place well after the initial transients have died away. The smaller box represents the region to be focussed on in Fig. 3. This is the site of the broad upwelling impinging on the rheological interface. At the interface there is a drop in effective viscosity owing to the stresses produced by the large-scale rapid circulation in the upper mantle. We notice that to the left the upwelling does not undergo a sharp change in the horizontal dimension. From the grey scale contours one can see that in the upper mantle there is a large patch of medium-temperature material being advected around by the large-scale flow. The descending currents are sluggish in the lower mantle because of the increase in viscosity, as the flow descends into a stiffer Newtonian lower mantle.
At the interface above the upwelling lower mantle plume a boundary layer has been formed. In Fig. 3 we show the sequential development of this boundary layer as it becomes unstable. The hot material is transported very efficiently in the upper mantle as a consequence of the lower effective viscosity of the upper mantle and the drop in viscosity as a consequence of the stresses produced by the local thermal instability. As the supply of hot material from the lower mantle is at a much lower rate, it takes time for the thermal boundary layer to form again, after most of the hot material has been extracted from the interface. From comparing the first and last snapshot, one obtains a distinct impression of the periodic nature of this flow. The sharp decrease in viscosity as a consequence of the thermal instability is indicated in the extrema in the stream function, of which the amplitude increases by a factor of nearly four. The diapiric events take place on a much shorter time scale than the overturn time associated with the large-scale circulation. The periodicity of this pulsating mechanism is more evident in the plots shown below (Fig. 6).

An effective $Ra$ of between $10^4$ and $10^5$ for the lower mantle may not be all that unreasonable when one takes into consideration the combined effects of a viscosity increase of 10 to 50 [25,26] a decrease in thermal expansivity of 6 [27] and an increase of lattice thermal conductivity in the lower mantle of around 4 to 6 [28,29].

In Fig. 4 we show the temporal development of the surface heat flux $q_s$. From the grey scale contours one can readily discern the discrete pulses of heat arriving at nearly the same spot. The heat flux associated with these thermal events (around 60) is nearly an order of magnitude greater than the background value of 7. It is also interesting to observe the effects of the excess heat flow being propagated with time. A trail of high heat flux is left along $x$ with time. Surrounding hotspot swells are typically anomalous heat flow patterns. This model shows the phenomenon of basal reheating of the lithosphere [30,31], which is based on field observations of the Hawaiian swell [32].

Figure 5 displays the temporal evolution of the vertical velocity along $x$ at a depth of $z = 0.05$, which is within the thermal boundary layer because the time averaged Nusselt ($Nu$) number at the surface is around 7 (see Fig. 6). The close correlation between peaks of the surface heat flow (Fig. 4) and the vertical upwellings at the base of the lithosphere indicates that thermal advection is the dominating mechanism responsible for the local heat transfer at these 'hot spots'.

Figure 6 summarizes the temporal evolution of the viscosity field, the root-mean-squared velocity $V_{rms}$, and the surface $Nu$. The sharp spikes of the thermal events are well recorded by $V_{rms}$, where amplification by a factor of 4 is found. The surface $Nu$ displays peaks with smaller amplitudes and behaves in a fashion that is similar to a nonlinear oscillator, in the form of a slow relaxation followed by a rapid rise. Temporal variations of the surface $Nu$ are much smaller than those associated with the local heat-flow values $q_s$. 

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Fig. 4. The temporal evolution of the local surface heat flux for case A. Heat flux is non-dimensionalized by $k \Delta T / d$, where $k$ and $\Delta T$ are respectively the thermal conductivity and temperature difference across the layer depth $d$. Time advances downward.
shown in Fig. 4. In strongly time-dependent flows there is a large disparity between the local \(q(x, t)\) and \(Nu(t)\) based on the horizontal average. It is therefore important to recognize this distinction in the application of the usual heat-transfer scaling relationships to local and global geological phenomena. Changes in the viscosity field reflect the stress variations generated during these pulsing events. The average value of the viscosity (middle curve) does not vary as much as the extrema (top and bottom curves). The contrast of the averaged viscosities between the upper and lower mantle is 200 for case A.

3.2 Model C (lithosphere, non-Newtonian upper mantle and Newtonian lower mantle)

More plate-like behavior can be brought out by introducing a thin top layer with a larger power-law index of \(n = 5\) [9]. For the time-dependent case we have used a \(n(z)\) of (5, 3, 1) and a \(A(z)\) of (11, 1, 1) for \(Ra \) of \(2 \times 10^4\). The global snapshot is shown in Fig. 7. From the stream function the more plate-like characteristics of model C compared to case A is clearly manifested. There is now a hot tongue emanating from the edge plume. This hot flat patch is pulled quite strongly by the large-scale circulation produced by the presence of the mobile lithosphere. Its hot tip nearly reaches the middle plume where the diapiric structures are developed. The thermal anomalies developed from the rheological interface for case A and C are diapiric in nature because they consist of detached blobs and do not form a continuous conduit.

Diapiric structures can also be formed in a wide variety of circumstances, such as in the interior of the cell for high Rayleigh number convection with Newtonian and non-Newtonian
rheologies [14,16] and at phase boundaries [17]. Here we have shown another mechanism based on changes in the rheological laws.

Figure 8 shows a sequence of temporal snapshots which have been zoomed around the central plume. Again, the episodic nature of this diapirc shedding process is illustrated by the development of $T$ and $\psi$. Time scales for these diapirc outbursts are short relative to that for the large-scale circulation. The lithosphere acts to control the upstream flow of plume material beneath the plate. The pulsating nature of this flow is further demonstrated by the time histories of $q_s$ and the vertical velocity $w$ in Fig. 9 and 10.

The presence of the lithosphere brings in another time scale, owing to the thermal diffusion time of a plate with finite thickness. The periodicity of the pulses in case C is reduced relative to case A because of the finite thickness of the lithosphere and the higher $Ra$. The effects of lithospheric shielding on the surface heat flow are clear from comparing Fig. 4 and Fig. 9. This low-pass filtering mechanism of the lithosphere on surface thermal signatures is thus illustrated.

\[ Ra = 2 \times 10^4 \quad n = (5,3,1) \quad A = (10,1,1) \]

Fig. 7. Global snapshots of $\psi$ and $T$ for model C. Model C has a lithosphere which extends to $z = 0.05$. The power-law indices $n(z)$ are 5 for lithosphere, 3 for the upper mantle and 1 for the lower mantle. The box represents zoom-in snapshots that follow later.

Note the hot wide patch to the left being pulled by the large-scale circulation in the upper mantle.
The presence of the lithosphere also helps to maintain better the stationarity of the diapiric signatures on the surface, as can be observed by comparing case A and C (see Figs. 4, 5, 9 and 10). The time histories of the viscosity field, $V_{rms}$, and the surface $Nu$ for case C are shown in Fig. 0.5.

The average viscosity contrast between the upper and lower mantle is about 200. The spiky character of non-Newtonian flows [15,16] is clearly displayed where the time scales are shorter than for the case without the lithosphere (compare with Fig. 6).

$$Ra = 2 \times 10^4 \quad n=(5,3,1) \quad A=(10,1,1)$$

Fig. 8. Sequence of close-up snapshots of $T$ for model C. Time scale for pulsating event is around 0.0004, which is shorter than the overturn time of large-scale background circulation.
3.3 Model H (Newtonian, viscously stratified)

One may well wonder whether a viscosity stratification in a purely Newtonian medium can also produce diapiric structures from the impingement of an ascending plume on an interface with a viscosity contrast of $O(10^2)$. Figure 12 shows the global shot of a Newtonian model with a viscosity jump of 200 and $Ra$ of $2 \times 10^4$. Again, this simulation was started from a steady-state solution perturbed by a thermal perturbation with an amplitude of 0.05 and sinusoidal variations having fundamental wavelengths along both directions.

Comparing Figs. 2, 7 and 12, we see that there are some differences in the flow patterns and temperature fields between a purely Newtonian rheological (case H) and the composite (case A and C) rheological models. The hot tongue so prominent in case C is now absent in the Newtonian model, because the large-scale circulation is not as coherent and strong in drawing out hot anomalies, as in the case with the lithosphere. An important difference with the non-Newtonian models is the large number of spontaneous downwellings from the upper boundary layer.

Figure 13 shows the evolution of the plume as it interacts with the interface. Again the hot material is being extracted from the interface faster than it is supplied, showing that the viscosity decreases is the fundamental reason for the formation of these diapiric structures. Differences exist, in that the plumes developed off the interface are broader and remain in one piece closer to the surface, as the local stress-induced viscosity decrease is absent in these Newtonian models.

From the snapshots of the stream function in Fig. 13 it is clear that the pulsating flow is much less vigorous than in the cases with the non-Newtonian mantle. This point is further emphasized
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by the temporal portrayal of the surface heat flow and the vertical velocity near the top (Fig. 14 and 15). The differences in the fluctuating intensity, periodicity and the plume stationarity between case A and H are striking. The Newtonian signals are much more diffused and tend to drift more. The nature of the diapirc flow is more episodic than periodic. The spontaneous downwellings give rise to the noisy character of $q_s$ and $w$.

Figure 16 summarizes the more sluggish Newtonian behavior by showing the temporal evolution of $V_{rms}$ and the surface $Nu$. Its departure from the spiky time series in Fig. 6 and 11 is quite evident.

4. Discussion and geophysical implications

The principal aim of this study was to investigate the effects of changes in creep laws, as suggested by recent experimental work on perovskite analogues [7], on the dynamics of time-dependent mantle flows. We wished to see what sort of flow structures would result from the interaction of a rising plume and a rheological interface, and whether these secondary plumes penetrate the upper mantle and manifest themselves as hotspots on the Earth's surface.

We found some new features in plume dynamics not encountered previously in time-dependent simulations with a homogeneous rheology throughout the mantle. These include (1) pulsating diapiric structures as a consequence of the viscosity decrease in the upper mantle, (2) fast time scales of the diapirs relative to those associated with the large-scale circulation, (3) the relative stationarity of recognizable surface heat-flow anomalies in models with a lithosphere and (4) the stabilizing influences of plates in modulating the upstream flow of rising plume material beneath the plate [33].

In addition to these points we have also shown the importance of a non-Newtonian upper mantle and lithosphere in maintaining plate-like behavior in time-dependent convection, which helps to generate coherent large-scale circulation in the upper mantle. This flow, in turn, can entrain hot anomalies away from a rising plume. Such shallow hot patches have recently been observed beneath the East Pacific by surface wave tomography [34]. The correlation of the horizontal extent of the negative seismic velocity contour with spreading rate [34] argues strongly for the dynamic role played by plate motions in pulling out hot patches from the ridge. On the other hand, Newtonian viscously stratified models do not produce pronounced broad hot patches in the upper mantle.

The question arises as to whether this type of diapiric upwelling can explain the periodicity in volcanism along hotspot chains. In Fig. 17 we show a comparison of the surface heatflow $q_s$ for the three models, in which the parameters are dimensionalized. The periodicity of the upwellings varies between 100 Myr (case A) to 30 Myr (case C), which is more than an order of magnitude larger than the period of volcanism that is, for example, observed in the spacing between the Hawaiian islands (1–2 Myr). From the speed with which the thermal signals travel
along the top boundary, we can estimate a plate velocity of approximately 1 cm/yr, which is at the lower end of the range of observed present-day plate velocities.

An important parameter in controlling the vigor of both the large-scale convection and the small-scale upwellings is the Rayleigh number. Increasing $Ra$ will increase the frequency of the pulsations (compare model A with model C) and the magnitude of the plate velocity.

The magnitude of the viscosity jump mainly influences the rate at which the hot material is extracted from the interface, especially in the cases of non-Newtonian viscosity, where the extraction rate is much faster than the supply rate. Decreasing the average viscosity in the upper mantle (or increasing the stress-induced drop in local viscosity by increasing the power-law index $n$) will tend to speed up the diapirs in the upper mantle; the period between the diapirs is however determined by factors governing the supply rate (the viscosity of the lower mantle and the vigor of convection), defined by the Rayleigh number.

Our choice of $Ra$ is conservative. It would be more realistic to choose a Rayleigh number that is a factor of 10 higher. Exploring the behavior of these layered models with higher $Ra$ is however strongly limited by computational time. Smaller spatial and temporal discretization must be used to resolve these small-scale features, and these 2D calculations challenge the limits of present-day computers.

Many interesting new consequences arising

\[
Ra = 2 \times 10^4 \quad n=(1,1) \quad A=(0.005,1)
\]

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![Streamfunction and Temperature](image)

**Fig. 12.** Global snapshot of $\psi$ and $T$ for the Newtonian model (H). The lower mantle is 200 times more viscous than the upper mantle. Note that the hot tongue from the plume to the left is small because of the more sluggish flow there.
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from vertical variations in creep laws have been discussed here. Important features for inclusion in future studies would be temperature- and pressure-dependent viscosity and depth-dependent thermal expansivity, as depth-dependent properties in Newtonian convection have been shown to

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Ra = 2 \times 10^4 \quad n=(1,1) \quad A=(0.005,1)
\]

Fig. 13. Sequence of snapshots depicting \( T \) and \( \psi \) of the Newtonian model (H).
stabilize the upwellings [35,36]. Phase transitions can also influence the dynamics of the diapirc structures [17]. The development of plume structures are shown to be strongly affected by the change in mantle rheological laws across the lower mantle-upper mantle boundary. Taken together with effects from phase transitions [17], this type of rheological layering, non-Newtonian on top of Newtonian, may greatly help in promoting diapiric flows in the upper mantle with fast time scales.

The seismological discontinuities in the upper mantle can be related to phase changes. Their

combined thermodynamic and rheological effects can be envisaged to modulate the behavior of upwelling diapirs and to bring shorter time scales into play as a consequence of the shorter distance to the surface.

Acknowledgements

We thank Ulli Hausen and Shun Karato for interesting discussion. Comments by an anonymous reviewer helped to improve the manuscript. Support for this research came from both the

Fig. 14. Same as for Fig. 4, but for Newtonian rheology.

Fig. 15. Same as for Fig. 5, but for Newtonian model.
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Dutch Academy For Scientific Research (NWO), the American National Science Foundation and the U.S. Army High Performance Computing Research Center. Computing was performed at both Minnesota Supercomputing Center and the SARA Supercomputing Center at Amsterdam. Wim Spakman is gratefully acknowledged for making his graphic software available to us. We thank M. Lundgren for help in preparing this manuscript.

References


Fig. 16. Same as for Fig. 6, but for Newtonian model.

Fig. 17. Comparison of the dimensionalized surface heat-flow characteristics for the three cases. The scheme for dimensionalization is given in Fig. 2.


