

On ‘steady-state’ heat flow and the rheology of oceanic mantle

Jun Korenaga

Department of Earth and Planetary Science, University of California, Berkeley, California, USA

Thomas H. Jordan

Department of Earth Sciences, University of Southern California, Los Angeles, California, USA

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[1] A constraint on the rheology of oceanic mantle is derived on the basis of a recently established scaling law for the onset of convection with temperature-dependent viscosity. Our approach is free from the assumption of statistically steady-state convection, which has commonly been employed in previous studies. The estimated range of asthenospheric viscosity is $1\text{--}4 \times 10^{19}$ Pa s if the activation energy is 300 kJ mol^{-1} , and $4 \times 10^{19}\text{--}10^{20}$ Pa s if the activation energy is 100 kJ mol^{-1} . The former range is consistent with laboratory data as well as geodynamic inference based on the geoid. The latter is consistent with geodynamic inference based on seamount loading history. Different activation energies predict different temperature contrasts in convecting mantle. Seismic tomography has the potential to discriminate between these possibilities. **INDEX TERMS:** 8120 Tectonophysics: Dynamics of lithosphere and mantle—general; 8130 Tectonophysics: Evolution of the Earth: Heat generation and transport; 8162 Tectonophysics: Evolution of the Earth: Rheology—mantle; 8180 Tectonophysics: Evolution of the Earth: Tomography. **Citation:** Korenaga, J., and T. H. Jordan, On ‘steady-state’ heat flow and the rheology of oceanic mantle, *Geophys. Res. Lett.*, 29(22), 2056, doi:10.1029/2002GL016085, 2002.

1. Introduction

[2] The dynamic state of oceanic asthenosphere is poorly known. Assuming that it is convecting (Figure 1a), however, we can relate surface observables to mantle rheology using scaling laws for thermal convection. Inferred rheology can then be compared with other types of constraints to verify the initial assumption. Surface heat flow is relatively constant at old ocean basins ($40\text{--}50 \text{ mW m}^{-2}$) [e.g., *Sclater et al.*, 1980; *Stein and Stein*, 1992], and this so-called steady-state heat flow has been used as a constraint on the viscosity of asthenospheric mantle. Heat flux is sensitive to the strength of convection, which in turn is a function of viscosity. Thus, if asthenospheric mantle is indeed convecting, oceanic heat flow can be used to infer its viscosity.

[3] Most of previous attempts have used a heat flux scaling law derived for steady-state convection (in a time-averaged sense), and their estimates are consistently within the range of $10^{18}\text{--}10^{19}$ Pa s [*Davaille and Jaupart*, 1994; *Doin et al.*, 1997; *Dumoulin et al.*, 1999; *Solomatov and Moresi*, 2000] (the scaling law derived by Davaille and

Jaupart is based on data during a quasi-equilibrium state achieved in a system only cooled from above). Even if asthenospheric mantle is convecting, however, whether or not the convection is at a statistically steady state is another question, which depends on the onset time of convection as well as the size of a convection system. The lifetime of an oceanic plate is typically less than 200 Myr. There may not be enough time for sublithospheric convection to reach a steady state. Given that the spatial extent of sublithospheric convection may not be limited to the upper mantle, arriving at a steady state before the plate subducts seems even more unlikely. A recent study indicates that a viscosity jump and an endothermic phase transition expected at 670-km depth can sustain upper-mantle-scale sublithospheric convection only for limited duration [*Korenaga and Jordan*, 2002].

[4] In this paper, we present an alternative, potentially more robust approach based on a scaling law derived for the onset of convection. This approach is similar to that of *Jaupart and Parsons* [1985], though time-independent viscosity structure was used in their analysis. Our study employs a newly established scaling law for temperature-dependent viscosity [*Korenaga and Jordan*, 2002], and it is dynamically consistent with the thermal and rheological evolution of oceanic mantle. Our estimate is $\sim 10^{19}\text{--}10^{20}$ Pa s, which is one order of magnitude higher than the previous estimates based on steady-state scaling laws. We will discuss its significance in comparison with relevant laboratory and geophysical studies.

2. Onset Time of Convection

[5] For a fluid with temperature-dependent viscosity, which has initially uniform temperature and whose surface temperature is lowered by ΔT instantaneously at $t = 0$, the onset time of convection may be expressed as [*Korenaga and Jordan*, 2002]

$$t_c^* = \left(\frac{4}{\pi^2 F(\mu^*, T^*)} \frac{Ra_c}{Ra} \right)^{\frac{2}{3}}. \quad (1)$$

The Rayleigh number, Ra , is defined as $\alpha \rho_0 g \Delta T D^3 / (\kappa \mu_0)$, where α is the coefficient of thermal expansion, g is gravitational acceleration, κ is thermal diffusivity, D is system height, and ρ_0 and μ_0 are reference density and viscosity, respectively, at the initial interior temperature. Time is normalized by diffusion time scale, D^2/κ . The critical Rayleigh number, Ra_c , is $\sim 2 \times 10^3$ for this transient cooling problem. The nonlinearity of an evolving tempera-

ture profile, T^* , and the temperature dependency of viscosity, μ^* , are all contained in the functional F , the definition of which is given by *Korenaga and Jordan* [2002]. We note that the onset time is independent of the system height D . The onset time is only sensitive to the rheological structure of shallow upper mantle.

[6] We use the Arrhenius law of temperature dependency, $\mu(T) = \exp(E/RT)$, where E is activation energy, and R is the universal gas constant. Figure 1b illustrates an example of predicted onset time for a range of activation energy. The previous scaling laws [*Davaille and Jaupart*, 1994, Equation 19; *Choblet and Sotin*, 2000, Equation 21] are defined in terms of the temperature derivative of logarithmic viscosity with internal temperature. Thus, with the Frank-Kamenetskii approximation, they can be applied to the Arrhenius law of temperature-dependent viscosity. Significant discrepancy from the previously known scaling laws can be seen, and readers are referred to *Korenaga and Jordan* [2002] for detailed discussion. The new scaling law

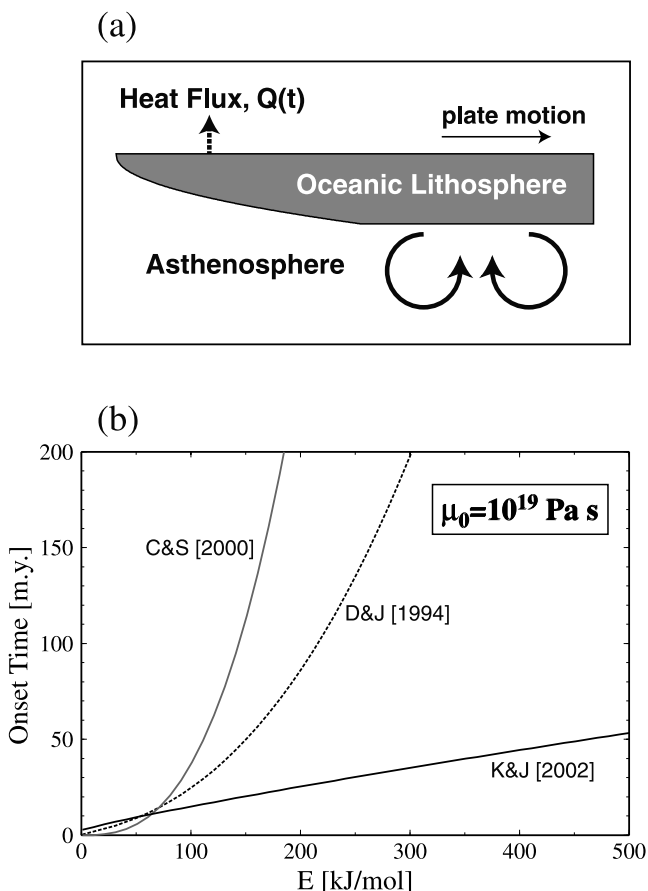


Figure 1. (a) Schematic drawing for the possible small-scale convection in oceanic asthenosphere. (b) Predicted onset time of convection based on the scaling law of *Korenaga and Jordan* [2002] is shown for reference asthenospheric viscosity of 10^{19} Pa s and the critical Rayleigh number of 2000. Predictions based on previous scaling laws [*Davaille and Jaupart*, 1994; *Choblet and Sotin*, 2000] are also shown. The new scaling law implies a much increased likelihood of sublithospheric convection in the oceanic mantle.

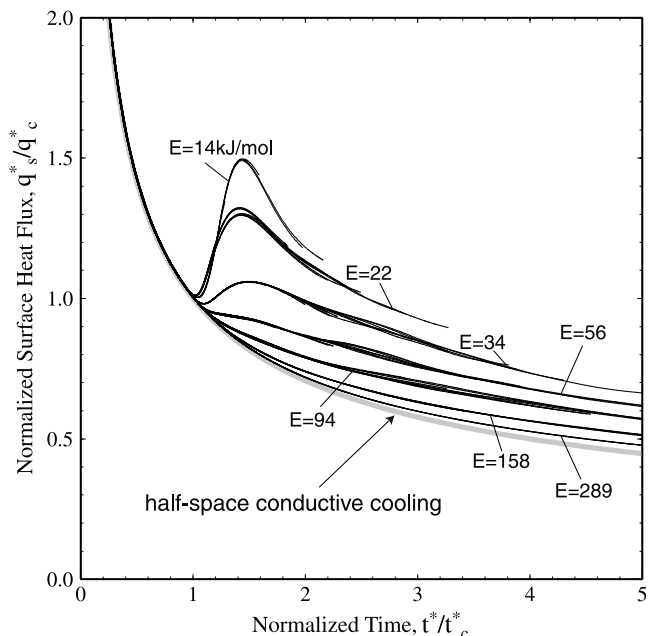


Figure 2. Summary of surface heat flux measurements from 2-D whole-mantle transient cooling models, with various combinations of the Rayleigh number, viscosity stratification between the upper and lower mantle, and an endothermic phase boundary at the base of the upper mantle. Modeling details are given by *Korenaga and Jordan* [2002]. The activation energy of the Arrhenius law is controlled so that all model runs have the same rigid lid thickness of 100 km at the onset of convection. That is, model runs with higher activation energy have higher Rayleigh number (lower reference viscosity). Time is normalized by onset time, and surface heat flux is normalized by the heat flux at the onset of convection.

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3. Heat Flow Constraint

[7] For realistic values of activation energy (i.e., a few hundreds kJ mol^{-1}), the initiation of sublithospheric convection does not significantly affect surface heat flux with respect to purely conductive cooling (Figure 2). Observed oceanic heat flux is indeed reasonably well explained by the half-space conductive cooling model. Slowly-varying heat flow during a limited time window (<200 Myr) probably appears to be relatively constant over old ocean basins. Using the half-space cooling model, therefore, this quasi-constant heat flux, q_{QC} (which is assumed to be 45 mW m^{-2} as derived by *Davaille and Jaupart* [1994]), may be expressed as

$$q_{QC} = \frac{k\Delta T}{D\sqrt{\pi\chi t_c^*}} \quad (2)$$

where k is thermal conductivity, and χ denotes the ambiguity of the onset time. For $\chi = 1$, convection takes place when surface heat flux reaches this quasi-constant value (i.e., at ~ 100 Myr). For $\chi = 2$, convection takes place

at as early as ~ 50 Myr. By combining equations (1) and (2), reference viscosity for asthenospheric mantle may be estimated as

$$\mu_0 = \frac{2F}{3Ra_c} \frac{\alpha \rho_0 g \Delta T^4 k^3}{\kappa \chi^{3/2} q_{0C}^3}. \quad (3)$$

[8] The possible range of asthenospheric viscosity is plotted as a function of activation energy in Figure 3, for several different combinations of χ , ΔT , and Ra_c . The rest of parameters are fixed as $\alpha = 3 \times 10^{-5} \text{ K}^{-1}$, $\rho_0 = 3300 \text{ kg m}^{-3}$, $k = 3.2 \text{ W m}^{-1} \text{ K}^{-1}$, $g = 9.8 \text{ m s}^{-2}$, and $\kappa = 10^{-6} \text{ m}^2 \text{ s}^{-1}$. The thickness of a rigid lid, which is the portion of the thermal boundary layer undisturbed by convection, is plotted in Figure 3b. Corresponding temperature difference driving sublithospheric convection is also shown in Figure 3c. The temperature difference is calculated on the basis of the lid thickness and the conducting geotherm at the onset of convection. Diffusion and dislocation creeps are two dominant deformation mechanisms for the Earth's mantle, with different activation energies ($240\text{--}540 \text{ kJ mol}^{-1}$) [Karato and Wu, 1993]. Because our analysis is based on the scaling law for the onset of convection, we can reasonably restrict ourselves to diffusion creep. The amplitude of convection is very small at the onset (by definition), so the effective viscosity of dislocation creep is virtually infinity, leaving diffusion creep as the only plausible deformation mechanism. Though dislocation creep associated with plate motion may be operating, it is probably unrelated to the onset of sublithospheric convection, whose planform tends to align perpendicular to plate motion to minimize convective interference [e.g., Richter, 1973]. The activation energy of diffusion creep ranges from 240 kJ mol^{-1} (wet) to 300 kJ mol^{-1} (dry), and for this range, our estimate is $1\text{--}4 \times 10^{19} \text{ Pa s}$. For temperature of $1350\text{--}1450^\circ\text{C}$ at the depth of 100 km , diffusion creep for $\sim 50\%$ water-saturated olivine [Hirth and Kohlstedt, 1996] predicts a similar range of viscosity. Hager [1991] estimated asthenospheric viscosity to be $2 \times 10^{19} \text{ Pa s}$ on the basis of the geoid, plate velocities and advected heat flux.

[9] One possible complication in the rheological structure of oceanic mantle is the effect of dehydration during mantle melting beneath mid-ocean ridges. Hirth and Kohlstedt [1996] proposed that the viscosity of residual mantle increases by two orders of magnitude. Gaherty et al. [1999] presented seismic evidence in support of the shallow depth of the dehydration boundary in the central Pacific. This introduces additional (time-independent) depth-dependency in mantle rheology. Under normal circumstances, however, this viscosity increase by dehydration becomes prominent only above $50\text{--}60 \text{ km}$ depth (see Figure 5 of Hirth and Kohlstedt [1996]). This depth level is within the rigid lid predicted by our analysis (Figure 3b), so the initiation of sublithospheric convection is unlikely to be influenced by this melting-related phenomenon at the present time.

[10] Recently, the geodynamic study of seamount loading history suggested asthenospheric viscosity of 10^{20} Pa s and activation energy of 120 kJ mol^{-1} , with the assumption of Newtonian rheology (diffusion creep) [Watts and Zhong, 2000]. This activation energy is too low to be consistent with currently available laboratory studies. It is interesting,

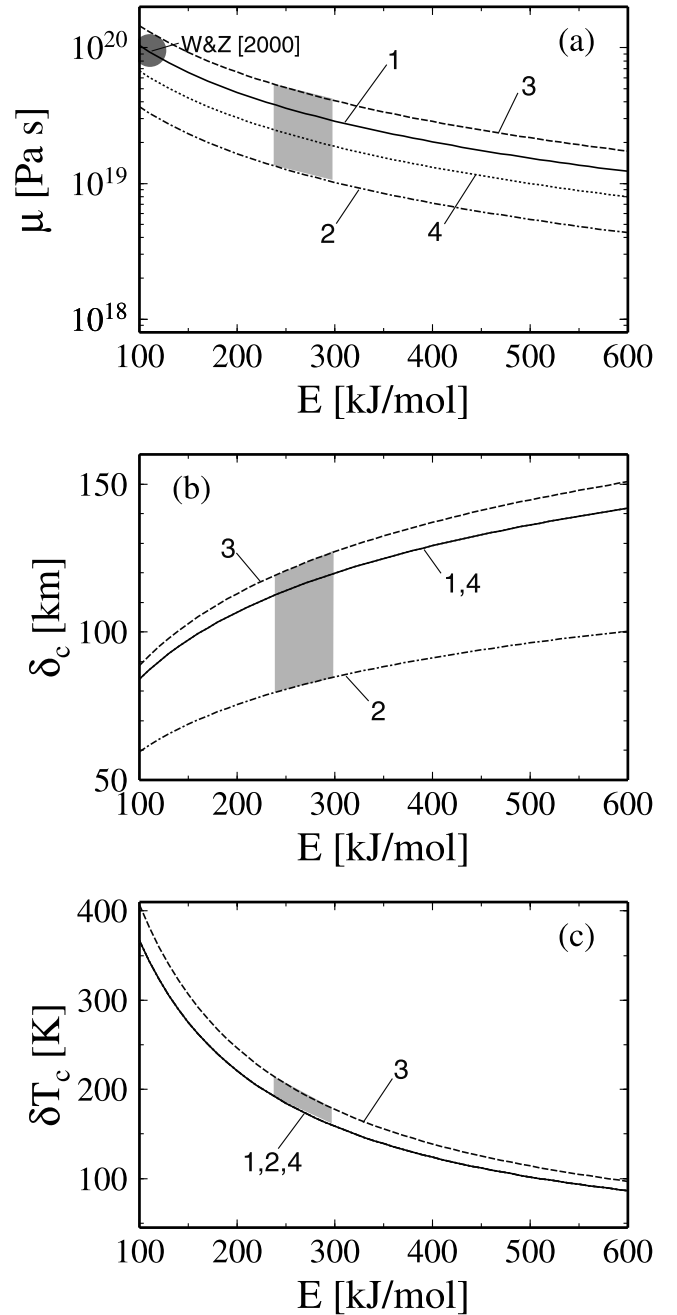


Figure 3. Heat flow constraint on mantle rheology and sublithospheric convection. (a) Asthenospheric viscosity, (b) rigid-lid thickness, and (c) temperature difference driving sublithospheric convection are plotted as a function of activation energy. Four cases are shown: 1. $\chi = 1$, $\Delta T = 1350^\circ\text{C}$, and $Ra_c = 1300$ (solid), 2. $\chi = 2$, $\Delta T = 1350^\circ\text{C}$, and $Ra_c = 1300$ (dot-dashed), 3. $\chi = 1$, $\Delta T = 1450^\circ\text{C}$, and $Ra_c = 1300$ (dashed), and 4. $\chi = 1$, $\Delta T = 1350^\circ\text{C}$, and $Ra_c = 2000$ (dotted). Onset time of convection is 93.0 Ma (for cases 1 and 4), 46.5 Ma (for case 2), and 107.3 Ma (for case 3). Shading indicates the range corresponding to the activation energy of diffusion creep ($240\text{--}300 \text{ kJ mol}^{-1}$). Estimate by Watts and Zhong [2000] is also shown in (a).

therefore, that our results can also be interpreted to support their inference (Figure 3a). Our approach can constrain only the pair of viscosity and activation energy; neither of them cannot be uniquely estimated. Different activation energies, however, predict significantly different temperature contrasts in convecting asthenosphere (Figure 3c). In the recent high-resolution tomography of the Pacific upper mantle [Katzman *et al.*, 1998], the peak-to-trough amplitude of shear-wave velocity anomalies is $\sim 3\text{--}4\%$, which can be interpreted as the temperature variation of $\sim 250\text{--}350\text{ K}$. If these velocity anomalies originate in sublithospheric convection, and if diffusion creep is an adequate deformation mechanism for such convection, the low activation energy suggested by Watts and Zhong [2000] may be preferable.

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J. Korenaga, Department of Earth and Planetary Science, University of California, 377 McCone Hall, Berkeley, CA 94720-4767, USA. (korenaga@seismo.berkeley.edu)

T. H. Jordan, SCI 103, Department of Earth Sciences, University of Southern California, Los Angeles, CA 90089-0740, USA. (tjordan@usc.edu)