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The clustering of rising diapirs and plume heads

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Abstract. Buoyant diapirs or plume heads in the Earth's mantle are typically modelled as a solitary bodies, ascending vertically and uninfluenced by neighboring diapirs. However, laboratory experiments and basic theory suggest that two or more rising diapirs can become attracted to each other, thereby influencing their ascent trajectory such that they cluster and possibly coalesce. Clustering is controlled by the horizontal and vertical separation between diapirs, as well as their relative sizes. This phenomena has implications for the origins of large igneous provinces as well as the Earth's long-wavelength hotspot distribution.

Introduction

Mantle plumes have been invoked to explain a wide variety of geological events such as hot spot volcanism, flood basalts, and continental rifting [see White and McKenzie, 1995]. Mantle plumes possibly initiate as large diapirs or plume heads in the D" layer above the core-mantle boundary. After the diapirs separate from D" it is assumed that they continue to rise along a vertical line as essentially solitary bodies. Any deflection from vertical ascent is presumed to arise from large-scale mantle flow associated with plate tectonic motions [e.g., Skilbeck and Whitehead, 1978; Olson and Singer, 1985; Richards and Griffiths, 1988]. However in this paper we present laboratory experiments and basic theory which suggest that rising diapirs can interact, attracting each other to form diapir clusters, and possibly merge into larger diapirs. The clustering and coalescence of diapirs naturally has implications for many geological phenomena, most notably for the origin of flood basalts, and the global distribution of hotspots.

Laboratory Experiments

The phenomena of diapir attraction occurs naturally in simple Rayleigh-Taylor experiments wherein a heavy viscous fluid overlies a thinner, less viscous and less dense fluid. To isolate and examine this phenomena we also performed experiments where the distance between the plumes could be controlled.

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General Rayleigh-Taylor Experiment

We conducted several basic Rayleigh-Taylor experiments in a cylindrical tank with a 30.75cm outer diameter, 29.8cm inner diameter, and 30.5cm height. The tank was filled to a height of 28.8cm with LSI Specialty Products Liquidose No. 444 corn syrup with density of 1460kg/m^3 and a viscosity of 607,329.8cP (where 1cP=10^{-3}Pa s). Over the Liquidose syrup was added a 5mm layer of Karo light corn syrup (dyed purple to enhance its visibility) with a density of 1320kg/m^3 and a viscosity of 3839cP. The tank was equilibrated for 12 hours, sealed and turned upside down. Approximately 11 min thereafter, the Karo syrup layer developed dome-shaped undulations 10–25mm in width. After another 15 min the undulations separated from the source area and began to rise as distinct diapirs. Approximately 10 min later, the attraction between two of the diapirs became apparent (Figure 1). One of the diapirs was slightly larger and rose faster than the other, eventually becoming the leading diapir. The trailing diapir became more oblate as its ascent path was deflected toward the leading diapir. In time the lower diapir moved to a trajectory almost directly underneath the leading diapir which also became distorted; after another 3 min the two diapirs merged.

Controlled Experiments

Controlled experiments were conducted in a rectangular tank (53cm long by 40.5cm high by 21.5cm wide) filled to a height of 33 cm with Karo syrup. Mounted above the tank were two dispensing burets beneath which were fixed 42cm long glass tubes with ends curved into a J-shape; the inner diameter of the tubes was 0.8cm and the upward-curved end of each tube was 7cm long. The burets were partially filled with a mixture of 40% (by volume) water and 60% Karo syrup with a density of 1080kg/m^3 and a viscosity of 103.2cP; fluid flux was controlled by stopcocks on the buret-nozzles. In each experiment, the diapirs were formed by completely filling the J-tubes with the plume mixture and then stopping the source feed. The diapirs inflated at the mouth of each J-tube to about 1.3 cm diameter after which they separated from their tube and proceeded to rise with a total ascent time of approximately 10 seconds. Multiple experiments were run for a range of separation distances between the J-tubes. No significant interaction between diapirs was noticed until the J-tubes were 3.4cm apart or less. Four runs were conducted with the 3.4cm separation. In two runs which

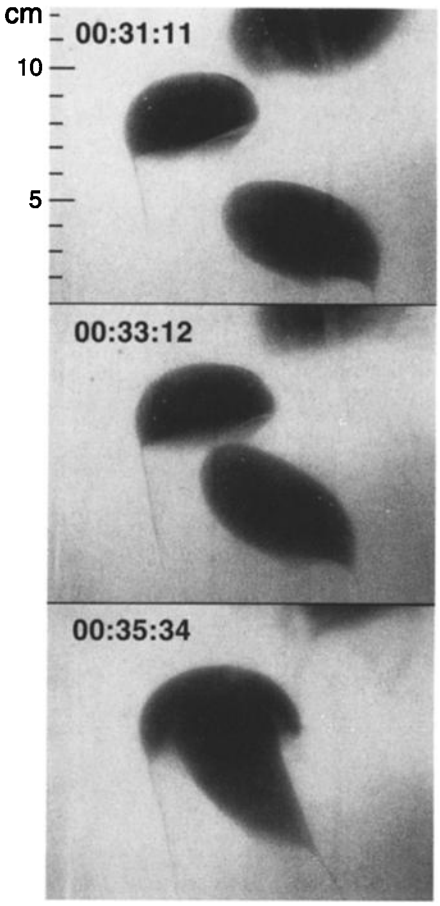

Figure 1. Sequential photographs of two clustering diapirs in a basic Rayleigh-Taylor experiment. Scale is shown on the left of the top frame; times (since start of experiment) are indicated in the upper left corner of each frame. See text for complete discussion of experiments.

showed strong diapir attraction, one diapir always lead the other by a significant distance (Figure 2). In one run, the two diapirs coalesced before reaching the surface of the tank. In the two runs which showed no noticeable interaction, neither diapir lead the other as the diapirs formed simultaneously and ascended at nearly the same rate.

Discussion of Experiments

Our experiments suggest that the vertical separation between diapirs is one of the primary criteria for clustering. For the diapirs to show any attraction, one of the two needed to be slightly higher than the other, by at least 1 cm. Total separation between diapirs also controlled attraction; diapirs displayed little interaction when separated by more than 3.4 cm. In the Rayleigh-Taylor experiments, attraction was apparent only when the leading diapir formed and, approximatey 5 minutes later, a second trailing diapir formed nearby.

Theory

In this section we theoretically examine diapir clustering by considering the interaction between two ascending low-viscosity spheres. The majority of work on the interaction between rising diapirs or bubbles is mostly numerical [e.g., Wijngaarden, 1993; Manga and Stone, 1993]. Here we devise a simple approximate theory of how two diapirs interact through the dipole flow and pressure fields that they establish in the surrounding viscous medium [see also Manga and Stone, 1993].

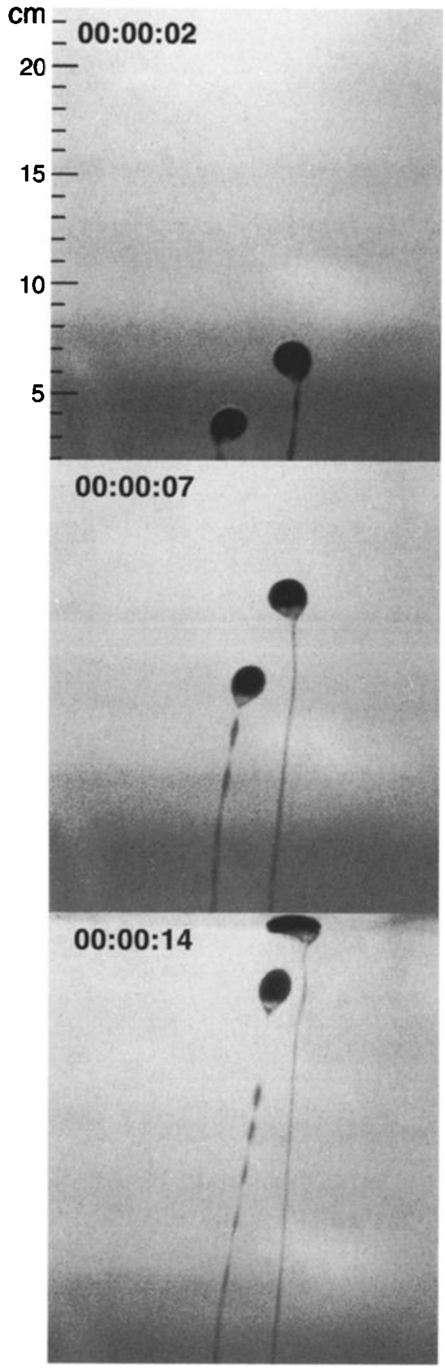

Figure 2. Sequential photographs of two clustering diapirs in a controlled experiment. The diapirs are generated by feeding fluid through two separate tubes, with adjustable separation, into the base of a tank of more viscous fluid. Scale is shown on the left of the top frame; times (since start of experiment) are indicated in the upper left corner of each frame. See text for complete discussion of experiments.

Our two diapirs have buoyant density anomalies (relative to the surrounding medium) of $\Delta\rho$, and have radii of a_1 and a_2. The surrounding medium is infinite, incompressible and has a viscosity of μ, which is much larger than the viscosities of the two diapirs. We define the x-z plane to contain the direction of gravity along the z axis and the line connecting the two diapirs. We assume that the diapirs maintain their spherical symmetry, thus there are no forces to pull them out of the x-z plane. The positions and velocities of each diapir are therefore $\mathbf{r}_1 = (x_1, 0, z_1)$, $\mathbf{r}_2 = (x_2, 0, z_2)$ and $\mathbf{v}_1 = (u_1, 0, w_1)$, $\mathbf{v}_2 = (u_2, 0, w_2)$, respectively. We assume that each diapir induces flow and pressure fields in the viscous medium which differ little from the dipole fields that would occur in solitary spheres moving through an infinite medium; this implicitly assumes that $|\mathbf{r}_1 - \mathbf{r}_2| \gg a_1, a_2$. Assuming creeping flow, the net force balance on diapir 1 is

$$0 = \frac{4}{3}\pi a_1^3 \Delta\rho g \hat{z} - 4\pi a_1 \mu \mathbf{v}_1 + \mathbf{F}_2 \tag{1}$$

where the first term on the right represents the buoyancy force of diapir 1, the second term is the drag force on diapir 1 due to its own motion, and $\mathbf{F}_2$ is the force due to the flow and pressure fields in the viscous medium induced by the motion of diapir 2. The flow and pressure fields caused by diapir 2 are governed by

$$0 = -\nabla P - \mu \nabla \times \omega \tag{2}$$

where P is the nonhydrostatic pressure, and ω is the vorticity vector in the viscous fluid [Batchelor, 1967]. Integration of (2) around a volume of the viscous fluid yields the sum of forces acting on that volume by the surrounding medium:

$$-\int\int_A (P\hat{n} + \mu\hat{n} \times \omega) dS = 0 \tag{3}$$

where A is the surface area of the volume, $\hat{n}$ is the unit normal to this surface; the above forces are in balance according to (2). However, if this volume were occupied by diapir 1 then the force balance of (3) would not hold. Much of the imbalance arises because the surface of diapir 1 is essentially shear-stress free and thus all tractions tangential to its surface, i.e., perpendicular to $\hat{n}$, would be negligible. Elimination of tractions tangent to the surface of diapir 1, i.e., $\mu\hat{n} \times \omega$, in (3) yields an approximate force on diapir 1 due to flow induced by diapir 2, i.e.,

$$\mathbf{F}_2 = -\int\int_A P\hat{n} dS \quad . \tag{4}$$

However, as we only use the dipole flow field induced by diapir 2 as if it were unaffected by the presence of diapir 1, we neglect the viscous normal stresses due to this flow impinging on and being deflected by diapir 1. (These normal stress effects can be approximated for rigid spheres using the method of reflections [Happel and Brenner, 1965; Manga and Stone, 1993]; there is, however, no equivalent method for inviscid diapirs.) Thus, our approximate force $\mathbf{F}_2$ is only valid as long as the viscous normal stresses on diapir 1 due to its own ascent are much larger than those due to diapir 2's induced dipole flow. Therefore, not only must diapir 1 be very far away from diapir 2 (so that diapir 2 maintains a dipole field), but diapir 1 must not be so small compared to diapir 2 that it moves too slowly. In brief, the theory assumes that the diapirs are not only widely separated but also are not of greatly disparate sizes.

Since the pressure associated with diapir 2 is assumed the same as if the diapir were solitary, then

$$P = \mu a_2 \frac{\mathbf{v}_2 \cdot (\mathbf{r} - \mathbf{r}_2)}{|\mathbf{r} - \mathbf{r}_2|^3} \tag{5}$$

[Batchelor, 1967] where $\mathbf{r}$ is the position vector of any point in space. The integral in (4) is around the surface of a sphere centered on $\mathbf{r}_1$; i.e.,

$$\mathbf{F}_2 = -\mu a_2 \int_0^{2\pi} \int_0^{\pi} \frac{\mathbf{v}_2 \cdot (\mathbf{r} - \mathbf{r}_2)}{|\mathbf{r} - \mathbf{r}_2|^3} \; \frac{\mathbf{r} - \mathbf{r}_1}{a_1} \; a_1^2 \sin\theta d\theta d\phi \tag{6}$$

where θ is colatitude and ϕ longitude of a spherical coordinate system centered at $\mathbf{r}_1$. Thus (6) is evaluated at $\mathbf{r} = \mathbf{r}_1 + a_1(\cos\phi \sin\theta, \sin\phi \sin\theta, \cos\theta)$. Defining $\Delta x = x_1 - x_2$, $\Delta y = y_1 - y_2$ and $\Delta r = \sqrt{\Delta x^2 + \Delta z^2}$, and expanding $|\mathbf{r} - \mathbf{r}_2|^{-3}$ to first order in $a_1/\Delta r$, (6) becomes

$$\mathbf{F}_2 = 4\pi a_2 \mu \frac{a_1^3}{\Delta r^3}\left(\alpha u_2 + \gamma w_2, 0, \gamma u_2 + \beta w_2\right) \tag{7}$$

where

$$\alpha = \frac{\Delta x^2}{\Delta r^2} - \frac{1}{3}, \qquad \beta = \frac{\Delta z^2}{\Delta r^2} - \frac{1}{3}, \qquad \gamma = \frac{\Delta x \Delta z}{\Delta r^2}. \tag{8}$$

Substitution of (7) into (1) yields two equations for u_1 and w_1; by symmetry, an interchange of subscripts 1 and 2 in these equations yields two more equations for u_2 and w_2. Combining all four equations, while neglecting terms of order $(a_1^m a_2^{3-m})^2/\Delta r^6$ ($m = 1$ or 2), yields

$$u_i = \frac{a_j a_i^2}{\Delta r^3} \frac{\Delta x \Delta z}{\Delta r^2} W_j \tag{9}$$

$$w_i = \frac{a_j a_i^2}{\Delta r^3}\left(\frac{\Delta z^2}{\Delta r^2} - \frac{1}{3}\right) W_j + W_i \tag{10}$$

where $(i,j) = (1,2)$ or $(2,1)$ and $W_i = \Delta\rho g a_i^2/(3\mu)$. Equations (9) and (10) show that the horizontal motion of one diapir is coupled to the vertical motion of the other diapir. If we nondimensionalize Δx, Δz and Δr by $\sqrt{a_1 a_2}$, and velocities $\mathbf{v}_1$ and $\mathbf{v}_2$ by $\Delta\rho g a_1 a_2/(3\mu)$, then the relative velocity of diapir 1 with respect to diapir 2 has components

$$u_1 - u_2 = \frac{d\Delta x}{dt} = \frac{1-\eta}{\sqrt{\eta}} \frac{\Delta x \Delta z}{\Delta r^5} \tag{11}$$

$$w_1 - w_2 = \frac{d\Delta z}{dt} = \frac{1-\eta}{\sqrt{\eta}}\left(\frac{3\Delta z^2 - \Delta r^2}{3\Delta r^5} - \frac{\eta+1}{\sqrt{\eta}}\right) \tag{12}$$

where $\eta = a_1/a_2$. Therefore there is no relative motion between the diapirs if they are of the same size (i.e., $\eta = 1$). The diapirs only converge on each other vertically if $d\Delta z^2/dt < 0$ which, by (12), occurs only if $(1-\eta)\Delta z > 0$ (given that $\Delta r \gg 1$), i.e., when the larger diapir is beneath the smaller diapir. The diapirs converge horizontally if $d\Delta x^2/dt < 0$ which, by (11), occurs only if $(1-\eta)\Delta z < 0$, i.e., if the smaller diapir is beneath the larger one. The theory therefore predicts two scenarios:

1. If the larger diapir begins above the smaller one, the vertical separation will grow, but the two diapirs will be drawn toward the same horizontal position and thus will cluster and rise along the same vertical track.
2. If the larger diapir begins beneath the smaller one, the vertical separation will decrease while the horizontal separation increases (as the larger diapir catches up to and pushes aside the smaller one) Once the larger diapir passes the smaller one, the system is described by Scenario 1 above. (See also Wijngaarden [1993] and Manga and Stone [1993].)

Although this model is only accurate for widely separated diapirs, it is in keeping with observations of the laboratory experiments, in particular that 1) there is little or no clustering when the diapirs are at the same level ($\Delta z = 0$); 2) that the interaction forces fall off rapidly with separation distance between diapirs; and 3) clustering largely occurs when the smaller trailing diapir moves toward the same horizontal position as the leading diapir.

Implications and Conclusions

Flood basalts Continental flood basalts and oceanic plateau basalts possibly originate from large plume heads which undergo partial melting upon arriving at the base of the lithosphere [Richards et al., 1989; c.f. White and McKenzie, 1989, 1995]. However, if rising plume heads can cluster and even coalesce, then the large plume heads which arrive at the lithosphere may result from the clustering and even merging of smaller diapirs. Since the final plume heads would have spent a portion of their ascent as smaller diapirs, their net rise time would be longer and their entrainment of surrounding mantle more profound. Moreover, the apparent occurence of double flood basalt events [Bercovici and Mahoney, 1994] may also be explained by the clustering effect (wherein a plume head draws a trailing plume head into its vertical ascent path).

Global hotspot distribution The global hotspot distribution was found by Ribe and de Valpine [1994] to have a very long wavelength signature, dominated by spherical harmonic degree $\ell = 2$ with secondary power at $\ell = 1$. This quadrupolar signature reflects two regions with a high density of hotspots in the South Pacific and Africa. Ribe and de Valpine inferred that if this distribution resulted from an infinitesimal Rayleigh-Taylor instability in the D" with dominant mode between $\ell = 1$ and 2, then the mantle in the D" would have to be approximately 10^6 times less viscous than the adjacent mantle. However, this assumes that the hotspot distribution at the surface reflects the wavelength of a linear instability at the core-mantle boundary. The clustering phenomenon discussed here shows that nonlinear interactions between fully developed plume heads can strongly influence their final distribution. That is, if plume conduits are also deflected toward each other as they follow their clustering plume heads, then the resulting hotspots would conceivably be clustered as well.

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