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Understanding Magma Ascent through Heat Propagation Buoyancy and Thermal Loss in the Lithosphere: a Thermodynamical Model (MagmaLitho v1)

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Abstract. We introduce a rather simple model for the ascent of the magma through the lithosphere, based on a mix of conduction, rock partial melting, and buoyancy, where we take into account the heat dissipation due to the melting process. Under suitable assumptions, a hyperbolic linear partial differential equation describing the process is derived from a general thermodynamical model whose constitutive equations are compatible with the second principle of thermodynamics. Then, we derive numerically a solution to a physically meaningful initial/boundary condition that can explain many of the experimental observations. More precisely, we suggest a mechanism for the magma emplacement able to explain the different depths of emplacement depending on one of the model parameters. Moreover, this parameter can be tuned to obtain the disappearance of the emplacement, which justifies the occurrence of hot spots (volcanoes forming a plume of molten rocks that rises from deep within the Earth's lithosphere).

1 Introduction

Many mechanisms have been proposed to explain the ascent of magmas through the lithosphere. These can be divided into two main classes (Turcotte, 1982; Burov et al., 2003). One considers low-inertia flows (diapirism) as a result of a viscosity or occasionally a density contrast between magma and wall rocks; so, a magma body moves within its rock matrix because of the difference between the densities of magma and the matrix. This kind of models includes Stokes flows (Cruden, 1988; Schmeling et al., 1988; Cruden, 1990; Mahon et al., 1988; Barnichon et al., 1999; Burov et al., 2003), Rayleigh-Taylor instabilities (Weinberg, 1996; Barnichon et al., 1999; Harig et al., 2010; Seropian et al., 2018), and Rayleigh-Bénard instabilities (Kukowski and Neugebauer, 1990; Hernlund et al., 2008; Yoshida et al., 2017). However, such flow is essentially restricted to the lower crust and the ductile crust. In fact, the mechanism appears to be efficient only at a depth more or less equivalent to the brittle-ductile transition (Taisne and Tait, 2009; Vigneresse and Clemens, 2000).

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Other models consider, besides the buoyancy, also heat diffusion, evaluating the conditions for rock partial melting or for magma segregation in the upper lithosphere (Marsh, 1982; Jackson et al., 2003; Polyansky et al., 2016; Schmeling et al., 2019). However, they are still confined to the upper lithosphere.

Another mechanism responsible for magma ascent is through lithospheric fractures. Again, magma upwelling is induced by the density contrast between the magma and its wall rocks. In this case, magma arises driven by a pressure gradient through crack propagation in the earth lithosphere (Scott et al., 1986; Spence and Turcotte, 1990; Rubin, 1995; Dahm, 2000; Taisne and Tait, 2009). Within this context, Keller et al. (2013) proposed a model coupling the mechanism of magma intrusion with the lithosphere deformation making use of arguments common in viscoelasticity and plasticity. Nevertheless, it has been observed that the crystals content of the magma is incompatible with the density contrast required for magma propagation in fractures (Vigneresse and Clemens, 2000). Consequently, such a mechanism could exist in some rather specific cases, and cannot be a general feature of magma ascent (Vigneresse and Clemens, 2000).

An interesting result of the second kind of models is the development of solitary waves (“magmons”) that can propagate without attenuation in the viscous medium (Scott and Stevenson, 1984; Scott et al., 1986; Dutta and Mandal, 2024). Magmons can also be produced when the magma ascends by capillarity (De Martino et al., 2003): in this case, preexisting fractures are viewed as capillary through which the magma propagates. In any case, the existence of solitary waves contradicts the emplacement of magmons at shallow depth as magma chambers. In fact, once a soliton starts to move, it cannot be stopped (Benjamin, 1972).

Here, we introduce a very different approach in which there does not exist a single mechanism for magma propagation in the lithosphere. So, we suggest a mix of three different physical phenomena: heat propagation, buoyancy, and heat dissipation for rock partial melting. The presence of a heat propagation phenomenon and of a dissipative term in our equation represents the novelty of our approach to magma ascent in the lithosphere.

The only assumption we need is the existence of a magma batch at the boundary between the asthenosphere and the lithosphere. This represents a very reasonable assumption for the existence of a volcano. As an example, a model for the generation of such a batch in the asthenosphere is provided in Godano et al. (2022); Munafò et al. (2024a). Moreover, after proposing a general thermodynamical framework, hereafter we limit ourselves to analyze a linearized model. This involves many approximations, but it considerably simplifies the model and makes the adopted mechanism for magma propagation through the lithosphere more understandable. The plan of the paper is the following. In Section 2 we analyze a general thermodynamical model and exploit its compatibility with the entropy principle. In Section 3, we specialize the model by linearizing the equations, derive a second order hyperbolic partial differential equation for the temperature, and solve numerically a physically meaningful initial/boundary value problem. Finally, in Section 4, some concluding remarks, as well as possible generalizations, are outlined.

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2 The model

The fundamental ingredients of our model are: 1) existence of a magma batch at the boundary between lithosphere and asthenosphere, 2) heat propagation, leading to partial melting of the surrounding rocks, 3) upwelling of the batch in the melted rocks by buoyancy in the newly melted rocks, 4) again partial melting of the surrounding rocks, and 5) repeating of whole mechanism until the temperature of the batch is smaller than the rock melting one.

Here, for simplicity, we assume a circular geometry for the magma batch. The temperature of such a batch T_m is higher than that of the surrounding rocks T_r. Then, due to a simple conductive mechanism, a circular annulus reaches a temperature $T_r < T_i < T_m$ close to the melting temperature of the rocks. As a consequence, the rocks in the circular annulus evolve into a fluid state but with a density $\rho_i > \rho_m$. Even if this condition could not be completely true, it remains a fundamental assumption of our model. This implies that the original magma batch starts to upwell for simple buoyancy reaching the top of the circular annulus. Here, the conductive mechanism resumes to be dominant, and the rocks in contact with the magma body at temperature T_m are melted at temperature T_i. The mechanism continues until the magma temperature becomes lower than the melting temperature of the surrounding rocks, and the magma body stops in the earth crust, forming a magma chamber. A simple sketch of the model is shown in Figure 1.

These general considerations are taken into account for building a general thermo–mechanical model within the framework of Extended Irreversible Thermodynamics (EIT) (Truesdell, 1969; Jou et al., 2008, 2010). Then, by a linearization procedure, we shall discuss the solutions of a linear equation embedding the mechanisms above described.

2.1 General thermodynamical model

In recent decades, many experiments have shown that classical transport laws, such as Fourier's law for heat conduction, Fick's law for diffusion, and Newton's law for viscosity, fail to accurately describe phenomena occurring under extreme conditions. In such situations, processes take place very rapidly and the assumption of local equilibrium is no longer valid. Classical laws lead to parabolic equations, implying infinite propagation speeds of perturbations, which contradict both experiments and fundamental theoretical considerations. To address this issue, it is necessary to include relaxation times of dissipative fluxes (heat, mass, momentum). However, classical thermodynamic frameworks cannot easily incorporate such evolution equations, since in some cases they would imply negative entropy production, violating the second law of thermodynamics. The need to describe nonequilibrium processes has led to the development of extended theories, such as Rational Extended Thermodynamics, Extended Irreversible Thermodynamics, thermodynamics with internal variables. These approaches treat dissipative fluxes as independent variables and introduce a generalized nonequilibrium entropy that remains compatible with the second law. These theories are also crucial for engineering applications, where relaxation times are not negligible. Most engineering models of heat conduction rely on Fourier's law, which relates the heat flux $\mathbf{q}$ to the gradient of the temperature T:

$$\mathbf{q} = -\lambda \nabla T.$$

Although highly successful in many contexts, Fourier's law combined with the energy balance equation leads to a parabolic heat equation with infinite propagation speed of thermal disturbances. This is physically unrealistic and violates causality.

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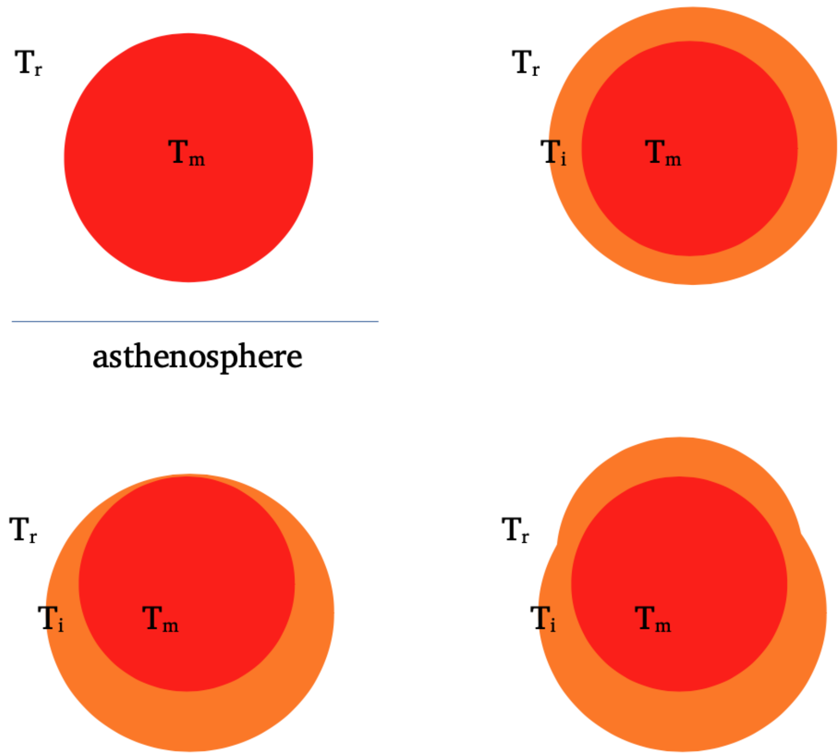

Figure 1. Sketch of model in which a magma body propagates through the lithosphere.

Experiments throughout the 20th century showed phenomena that cannot be explained within the Fourier framework. The most notable example is second sound, where heat propagates as a wave. These observations motivated the development of generalized heat conduction models. The first major extension is the Maxwell-Cattaneo-Vernotte (MCV) equation, which introduces the relaxation time τ for the heat flux and leads to a hyperbolic heat equation:

$$\tau(\mathbf{q}_t + \mathbf{v}\cdot\nabla\mathbf{q}) + \mathbf{q} = -\lambda\nabla T.$$

This model accounts for memory effects and restores finite propagation speed. However, experimental evidence shows that material parameters such as thermal conductivity, specific heat, relaxation time, and sometimes even mass density can vary significantly with temperature. Introducing these nonlinearities strongly affects heat propagation.

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In the framework of Extended Irreversible Thermodynamics (EIT) for heat conducting continua (Jou et al., 2008, 2010), let us consider the one–dimensional system of balance laws (z denotes the ascending vertical direction)

$$\begin{aligned}
&\mathcal{E}_1 \equiv \rho_t + v\rho_z + \rho v_z = 0,\\
&\mathcal{E}_2 \equiv \rho(v_t + v v_z) - \mathbb{T}_z = -\rho g,\\
&\mathcal{E}_3 \equiv \rho(\varepsilon_t + v\varepsilon_z) + q_z - \mathbb{T} v_z = r_1,\\
&\mathcal{E}_4 \equiv \tau(q_t + v q_z) + \phi_z = r_2,
\end{aligned} \tag{1}$$

where ρ is the mass density, v the velocity, $\mathbb{T}$ the Cauchy stress tensor, g the gravity acceleration, ε the internal energy per unit volume, q and ϕ the heat flux and the flux of heat flux, respectively, r_1 and r_2 the productions of ε and q, respectively, whereas the subscripts $_t$ and $_z$ denote the time and spatial derivatives, respectively. The balance equation for heat flux is a generalization of Cattaneo equation, and τ is a relaxation time (Cattaneo, 1948; Munafò et al., 2024b).

In continuum thermodynamics, the entropy principle constitutes a basic tool in developing constitutive theories. It requires the constitutive relations to be such that all thermodynamical processes, i.e., all solutions of the field equations, satisfy the entropy inequality. Along any arbitrary thermodynamic process the entropy production must be non-negative, so that, in local form, the entropy production writes as

$$\sigma_s = s_t + v s_z + J_z \geq 0 \tag{2}$$

(Clausius–Duhem inequality), where s is the entropy per unit volume and J is the local entropy flux.

The field equations (1) and the entropy inequality (2), once the variables entering the state space have been assigned, must be supplemented by the constitutive equations for the stress tensor $\mathbb{T}$, the flux of the heat flux ϕ, the productions r_1 and r_2, the specific entropy s, and entropy flux J. We use the framework of EIT, a modern thermodynamic theory in which both the dissipative fluxes and their gradients are regarded as non-equilibrium state variables. In this way, let us assume the state space to be spanned by

$$\mathbb{Z} = \{\rho, \varepsilon, q, \rho_z, v_z, \varepsilon_z, q_z\}. \tag{3}$$

To find a set of conditions which are at least sufficient to guarantee the inequality (4) to be satisfied, we can use either the classical Coleman-Noll (Coleman and Noll, 1963) procedure or the Liu (Liu, 1972) one. However, when dealing with non-local constitutive relations, an extended approach reveals successfully to ensure the compatibility of non-local constitutive equations with the second law of thermodynamics. In this paper, we apply an extended Liu procedure (Liu, 1972), incorporating new restrictions consistent with higher-order non-local constitutive theories (Cimmelli et al., 2011; Gorgone et al., 2020).

According to this procedure, in the entropy inequality, we have to use as constraints the field equations, and their gradient extensions too, up to the order of the gradients entering the state space. In our case, due to the choice of the state space (3), we have to use as constraints the first order gradients of mass, linear momentum, energy and heat equations. More in detail, this technique consists of subtracting from the Clausius–Duhem inequality a linear combination of the field equations and of the spatial gradients of the latter, and the coefficients of this linear combination are the so–called *Lagrange multipliers* which

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may depend on the field and state space variables. Therefore, the number of independent constraints to be taken into account is always equal to that of the unknown fields and the gradients entering the constitutive equations as independent state variables.

Then, the entropy production reads

$$\sigma_s = (s_t + v s_z) + J_z - \lambda_1 \mathcal{E}_1 - \lambda_2 \mathcal{E}_2 - \lambda_3 \mathcal{E}_3 - \lambda_4 \mathcal{E}_4 \\ - \Lambda_1 (\mathcal{E}_1)_z - \Lambda_2 (\mathcal{E}_2)_z - \Lambda_3 (\mathcal{E}_3)_z - \Lambda_4 (\mathcal{E}_4)_z \geq 0, \tag{4}$$

where λ_i and Λ_i ($i = 1, \ldots, 4$) are suitable Lagrange multipliers to be further determined.

There is a simple rationale for using the gradients of the governing equations as additional constraints in the entropy inequality when dealing with non-local constitutive equations. Thermodynamical processes are the solutions of the field equations, and, if these solutions are smooth enough, are trivially solutions of their differential consequences. Since the entropy inequality (4) has to be satisfied in arbitrary smooth processes, then, from a mathematical viewpoint, we have to use also the differential consequences of the equations governing those processes as constraints for such an inequality (Rogolino and Cimmelli, 2019). On the contrary, if we limit ourselves to consider as constraints only the field equations, we are led straightforwardly to a specific entropy and Lagrange multipliers which are independent of the gradients of the variables entering the state space, and, as a consequence, a Cauchy stress tensor depending on first order gradients would become incompatible with second law of thermodynamics.

By expanding derivatives with the chain rule, the inequality (4) can be written as a polynomial in some derivatives of field variables with coefficients depending on field and state space variables, in which we can distinguish:

- *highest derivatives*, that are both the time derivatives of the field and state space variables, which cannot be expressed through the governing equations as functions of the thermodynamical variables, and spatial derivatives with highest order;

- *higher derivatives*, that are the spatial derivatives whose order is not maximal but higher than that of the gradients entering the state space.

As a consequence of the choice of the state space (3), the components of the highest and higher derivatives are

$$\mathbf{X} = \{\rho_t, e_t, q_t, v_t, \rho_{xt}, e_{xt}, q_{xt}, v_{xt}, \rho_{xxx}, \varepsilon_{xxx}, q_{xxx}, v_{xxx}\},$$
$$\mathbf{Y} = \{\rho_{xx}, \varepsilon_{xx}, q_{xx}, v_{xx}\},$$

respectively; using these identifications, the entropy inequality (4) can be rewritten in compact form as

$$\mathbf{A} \cdot \mathbf{X} + \mathbf{Y}^T \mathbf{B} \mathbf{Y} + \mathbf{C} \cdot \mathbf{Y} + D \geq 0, \tag{5}$$

where $\mathbf{B}$ is a symmetric matrix of order m, $\mathbf{A}$ is a vector with n components, $\mathbf{C}$ is a vector with m components, and D is a scalar; $\mathbf{A}$, $\mathbf{B}$, $\mathbf{C}$, and D depend at most on the field variables and the gradients entering the state space; moreover, n is the number of the highest derivatives, and m the number of the higher derivatives. The left-hand side of the inequality (5) is a polynomial which is linear in the highest derivatives and quadratic in the higher ones. Since the highest and higher derivatives

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may assume arbitrary values, in order to fulfill the second law of thermodynamics, their coefficients must vanish; otherwise, the entropy inequality could be easily violated. Since in principle nothing prevents the possibility of a thermodynamical process where $D = 0$, in order for the inequality (5) to be satisfied for every thermodynamical process, the conditions

$$
\begin{aligned}
&\mathbf{A} = \mathbf{0},\\
&\mathbf{C} = \mathbf{0},\\
&\mathbf{B} \quad \text{is a positive semidefinite matrix},\\
&D \geq 0,
\end{aligned}
\tag{6}
$$

are sufficient for the fulfillment of the entropy inequality. We note that the terms involving highest derivatives provide us the Lagrange multipliers and some constraints on the constitutive equations. On the contrary, the terms involving higher derivatives provide us further restrictions on the constitutive equations.

We also observe that the thermodynamic restrictions are obtained by the exploitation of a rigorous mathematical procedure; due to their length, we do not try to attribute to them a detailed physical meaning. Nevertheless, the effectiveness of the procedure can be judged on the basis of the physical admissibility of the results it produces. By using some algebra, we are able to solve all the restrictions (6), and explicit constitutive equations for the stress tensor, the flux of the heat flux, the productions, the specific entropy, and the entropy flux are determined.

Here, we exhibit a solution to all the thermodynamic restrictions imposed by the entropy inequality (5):

$$
\begin{aligned}
&\phi = \phi_0 + \frac{1}{3}\frac{d(\rho s_1 \phi_1)}{d\rho}\left(\frac{d(\rho s_1)}{d\rho}\right)^{-1} q^2,\\
&r_1 = r_{10} + r_{11}q^2 + r_{12}q\rho_z + r_{13}q\varepsilon_z + r_{14}q_z + r_{15}v_z,\\
&r_2 = r_{20}q + r_{21}\rho_z + r_{22}\varepsilon_z + \left(\phi_1 - \frac{1}{3}\frac{d(\rho s_1 \phi_1)}{d\rho}\left(\frac{d(\rho s_1)}{d\rho}\right)^{-1}\right) q q_z\\
&\quad + r_{23}q v_z + r_{24}\rho_z v_z + r_{25}\varepsilon_z v_z + r_{26}q^3,\\
&s = s_0 + s_1 q^2,\\
&J = \frac{\partial s_0}{\partial \varepsilon}(1 - r_{14})q + \frac{2}{3}\frac{d\phi_1}{d\rho}\left(\frac{d(\rho s_1)}{d\rho}\right)^{-1}\rho^2 s_1^2 q^3,\\
&\mathbb{T} = \rho^2 \frac{\partial s_0}{\partial \rho}\left(\frac{\partial s_0}{\partial \varepsilon}\right)^{-1} - r_{15} + \rho\left(\frac{ds_1}{d\rho}\rho - 2r_{23}s_1\right)\left(\frac{\partial s_0}{\partial \varepsilon}\right)^{-1} q^2\\
&\quad - 2r_{24}\rho s_1\left(\frac{\partial s_0}{\partial \varepsilon}\right)^{-1} q\rho_z - 2r_{25}\rho s_1\left(\frac{\partial s_0}{\partial \varepsilon}\right)^{-1} q\varepsilon_z + \mu v_z,
\end{aligned}
\tag{7}
$$

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where $\phi_0 \equiv \phi_0(\rho,\varepsilon)$, $\phi_1 \equiv \phi_1(\rho)$, $r_{1i} \equiv r_{1i}(\rho,\varepsilon)$ $(i = 1,\ldots,5)$, $r_{2j} \equiv r_{2j}(\rho,\varepsilon)$ $(j = 1,\ldots,6)$, $s_0 \equiv s_0(\rho,\varepsilon)$, $s_1 \equiv s_1(\rho)$, $\mu \equiv \mu(\rho,\varepsilon)$ are arbitrary functions of the indicated arguments, along with the constraints

$$\begin{aligned}
&r_{11}\frac{\partial s_0}{\partial \varepsilon} + 2r_{20}\rho s_1 \geq 0, \qquad s_1 \leq 0, \qquad \mu \geq 0,\\
&2\frac{\partial \phi_0}{\partial \rho}\rho s_1 + \frac{\partial}{\partial \rho}\left(r_{14}\frac{\partial s_0}{\partial \varepsilon}\right) - \frac{\partial^2 s_0}{\partial \rho \partial \varepsilon} - r_{12}\frac{\partial s_0}{\partial \varepsilon} - 2r_{21}\rho s_1 = 0,\\
&2\frac{\partial \phi_0}{\partial \varepsilon}\rho s_1 + \frac{\partial}{\partial \varepsilon}\left(r_{14}\frac{\partial s_0}{\partial \varepsilon}\right) - \frac{\partial^2 s_0}{\partial \varepsilon^2} - r_{13}\frac{\partial s_0}{\partial \varepsilon} - 2r_{22}\rho s_1 = 0.
\end{aligned} \tag{8}$$

Some comments about the thermodynamic restrictions (7) and (8) are in order. By defining the absolute temperature T with the classical thermodynamical relation $\dfrac{1}{T} = \dfrac{\partial s_0}{\partial \varepsilon}$, in (7) we note that the entropy flux is no longer given by the constitutive equation postulated in rational thermodynamics, namely, $J = \dfrac{q}{T}$, and an entropy extra-flux is obtained. Also, as expected, the Cauchy stress tensor $\mathbb{T}$ and the productions r_1 and r_2 involve terms in the gradients entering the state space. Moreover, in the specific entropy s we recognize two contributions: the first one, s_0, representing the equilibrium entropy defined for homogeneous states, and the second one, quadratic in q, whose coefficient must satisfy the principle of maximum entropy at the equilibrium ($s_1 \leq 0$). Finally, we notice that the viscosity coefficient μ must be non-negative, in order to be physically meaningful.

3 Linearized model

In this Section, after specializing the constitutive functions according to (7) and (8) so that the most relevant mechanisms are captured, we linearize the model equations (1), and derive a second order partial differential equation ruling the evolution of the temperature T.

Let us consider the assignments

$$\begin{aligned}
&\phi_0 = \frac{\lambda}{c_v}\varepsilon, \qquad \rho^2\frac{\partial s_0}{\partial \rho}\left(\frac{\partial s_0}{\partial \varepsilon}\right)^{-1} = -p\\
&r_{10} = \kappa\rho\varepsilon, \qquad r_{20} = -1, \qquad r_{21} = 0,
\end{aligned} \tag{9}$$

where $p \equiv p(\rho,\varepsilon)$ represents the pressure, τ is a positive constant, κ and λ are arbitrary constants, c_v is the specific heat, along with the functions r_{22} and μ that are assumed to be constant. The assumptions in (9) can be easily explained by noting that we want to obtain a linear equation. Other coefficients are fixed at specific values to ensure that the terms have the correct dimensions.

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Then, considering the velocity field as a constant, *i.e.*, $v = u$, $\rho = \rho_0 + \tilde{\rho}$, where $\tilde{\rho}$ represents the perturbation of the constant mass density ρ_0, and neglecting the higher order terms, the model equations (1) become

$$\tilde{\rho}_t + u\tilde{\rho}_z = 0, \tag{10a}$$

$$p_z = -(\rho_0 + \tilde{\rho})g, \tag{10b}$$

$$\rho_0(\varepsilon_t + u\varepsilon_z) + q_z = \kappa\rho_0\varepsilon, \tag{10c}$$

$$\tau(q_t + uq_z) = -q + \left(r_{22} - \frac{\lambda}{c_v}\right)\varepsilon_z, \tag{10d}$$

By cross differentiating Equations (10c) and (10d), using the relation $\varepsilon = c_v T$, and choosing $r_{22} = -\tau\rho_0 u^2$, we derive the following second order differential equation for the temperature T:

$$\tau\frac{1}{1-\kappa\tau}T_{tt} + T_t - d\frac{1}{1-\kappa\tau}T_{zz} + uT_z + 2\tau u\frac{1}{1-\kappa\tau}T_{tz} - \frac{\kappa}{1-\kappa\tau}T = 0, \tag{11}$$

where $d = \dfrac{\lambda}{\rho_0 c_v}$ is the thermal diffusivity. Under the hypothesis that $1 - \kappa\tau \approx 1$ (this will be the case), we can write Equation (11) as:

$$\tau T_{tt} + T_t - dT_{zz} + uT_z + 2\tau u T_{tz} - \kappa T = 0, \tag{12}$$

where the subscripts $_t$ and $_z$ denote the time and space partial derivatives, respectively.

Equation (12) represents the described model. It is a linear hyperbolic equation; the first three terms corresponds to the Cattaneo equation (Cattaneo, 1958; Chester, 1963), the fourth one takes into account the buoyancy, the fifth term (12) describes the heat dissipated in the rock partial melting process, and the last term comes from an energy source (here supposed to be the heat coming from the asthenosphere).

As far as the parameters in (12) are concerned, τ is the viscoelastic relaxation time of the lithosphere, d the thermal diffusivity, u the buoyancy velocity, $r = 2\tau u$ the radius of the melted circular annulus, and κ a constant taking into account the contribution of the asthenosphere heat. Here we would like to emphasize that the introduction of the first term in equation (12) represents the main novelty of our model. The second-order time derivative of the temperature, first introduced by Cattaneo (1948, 1958), resolves the paradox of infinite propagation speed inherent in Fourier's equation, transforming the problem from a diffusive one into a propagative one. Moreover, we introduce a new dissipative term (the fifth one), which, to the best of our knowledge, has never been used before.

In obtaining Equation (12), we assumed constant viscosity, velocity and density. This could be an over-simplification for our batch of magma. However, in modeling magma ascent through the lithosphere, it is common to assume constant density, viscosity, and buoyant velocity in order to isolate the primary transport mechanisms and retain mathematical tractability. This simplification is well established in geophysical fluid dynamics. For instance, the Boussinesq approximation treats density as constant except in the buoyancy term, enabling simplified yet physically meaningful descriptions of thermally driven flows (Turcotte and Schubert, 2014). Similarly, many analytical and semi-analytical models of magma migration and dike propagation assume constant viscosity to avoid strong nonlinear coupling between temperature, composition, and rheology (Marsh,

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1982, 1989; Head and Wilson, 1992). In addition, constant ascent velocity is often used as a first order parameterization in advection–diffusion models of heat and mass transfer in magmatic systems, particularly when focusing on steady or quasi-steady regimes (Spiegelman, 1993), or in reactive melt transport and advection models (Dominguez et al., 2024). Similarly, recent fluid mechanical models of magma ascent, such as those describing the rise of magma batches, bubbles, or xenolith-bearing melts, typically rely on constant density contrasts and simplified rheological parameters to derive ascent velocities and scaling laws (Berbineau et al., 2026; Dragoni, 2022). Overall, these assumptions represent a standard and justified compromise that allows one to capture first order physics in order to isolate the dominant transport mechanisms and maintain analytical or numerical tractability, before introducing additional complexities. The obtained solution is simple and realistic, so justifying our assumptions.

Of course, density, viscosity and buoyancy velocity are temperature dependent. The analysis of a second order nonlinear hyperbolic equation for temperature where density, viscosity and buoyancy velocity depend on temperature is currently under investigation, and will be the content of a forthcoming paper.

Equation (12) describes the evolution of temperature in time and along the vertical direction. This approximation comes from the observation that a magma batch propagates mainly in one direction. In fact, we checked that, even in three space dimensions, when the propagation is allowed only in one direction, our results do not change substantially.

Let us define dimensionless time and space variables coherent with the phenomenon at hand:

$$\bar{t} = \frac{t}{\alpha}, \qquad \bar{z} = \frac{z}{\beta},$$

where $\alpha = 10^8\,s$ (roughly corresponding to three years), and $\beta = 2 \cdot 10^4\,m$ (the lithosphere thickness). Therefore, Equation (12) (dropping the bars to simplify the notation) reads as

$$\frac{\tau}{\alpha} T_{tt} + T_t - \frac{d\alpha}{\beta^2} T_{zz} + \frac{u\alpha}{\beta} T_z + \frac{r}{\beta} T_{tz} - \kappa\alpha T = 0, \tag{13}$$

to be studied in the spatial domain $z \in [0,1]$ for $t \geq 0$. Taking as relaxation time $\tau = 10^5\,s$ (a little bit more than one day) (Weber et al., 2025), and $\kappa = -2 \cdot 10^{-9} s^{-1}$, the solutions of Equation (13) will be affected by the values of u and d. In the following, we will consider the following four cases

- $u = 10^{-4}\,m\,s^{-1}$, $d = 10^{-6} m^2 s^{-1}$
- $u = 1.5 \cdot 10^{-4}\,m\,s^{-1}$, $d = 10^{-6} m^2 s^{-1}$
- $u = 10^{-4}\,m\,s^{-1}$, $d = 2 \cdot 10^{-6} m^2 s^{-1}$
- $u = 1.5 \cdot 10^{-4}\,m\,s^{-1}$, $d = 2 \cdot 10^{-6} m^2 s^{-1}$

Let us consider the following initial (see Figure 2) and boundary conditions:

$$\begin{aligned} &T(z,0) = 1230 \exp\left(-\frac{(z-0.1)^2}{0.05}\right), \\ &T(0,t) = 1007.04. \end{aligned} \tag{14}$$

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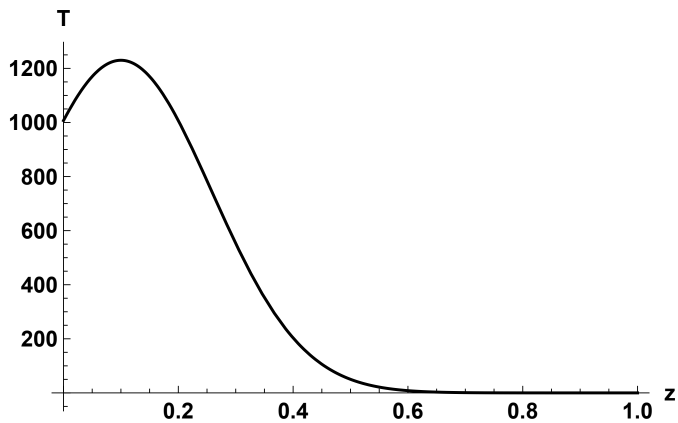

Figure 2. Initial condition for temperature.

The numerical integration (Wolfram Research, Inc., 2024) of Equation (13) with conditions (14) yields the solution shown in Figure 3.

Moreover, in Figures 4 and 5, the solutions profiles at different times are displayed.

Although Equation (13) does not represent a complete description of the physics of magma melting, crystallization, and differentiation during its lifetime, it provides a general and simple approach for the investigation of magma rising in the lithosphere that finally forms the magma chamber.

The initial and boundary conditions (14) represent the temperature profile of the superposition of the magma batch and the lithospheric temperature profile. In fact, the temperature profile without the magma batch in the lithosphere is approximately logarithmic (Ranalli and Rybach, 2005; Cloetingh et al., 2010). This means that, when we start from the bottom of the lithosphere, the profile decays exponentially. Adding a gaussian profile for the magma batch temperature, we get the profile (14) because the exponential decay becomes negligible with respect to the perturbation represented by the magma batch. Moreover, (14) helps us to show how the magma batch propagates through the lithosphere. This implies that the batch dimension is ≈ 4 km assuming that the lithosphere temperature decays exponentially from $\approx$ 1000°C, at the boundary with the asthenosphere, to 0°C at the Earth surface. Equation (13) describes the propagation of the magma batch through the lithosphere (see Figure 3). The magma propagates until its temperature becomes lower than the rock melting one, here assumed as 1000°C. This occurs at $t = 0.9$ (Figures 4 and 5) implying that the propagation duration is $\approx$2.7 years. Of course, the magma emplacement depends on the value of u, whereas it appears to be independent of d (Figures 4 and 5). The parameter u can be tuned in order to obtain different emplacement depths up to the limit of the Earth surface (in this limit we can consider it an hot spot, namely a direct alimentation from the asthenosphere). Moreover, the different values of u can explain the different magma compositions. In fact, in the extreme case of a small value for u, we should have a greater contamination of the original magma with lithospheric materials. Conversely, when u assumes larger values, buoyancy is more efficient and the magma remains more similar to its original composition (Reiners et al., 1995). Of course, any intermediate value of the parameters can provide an explanation of the various magma spectra from mafic to felsic.

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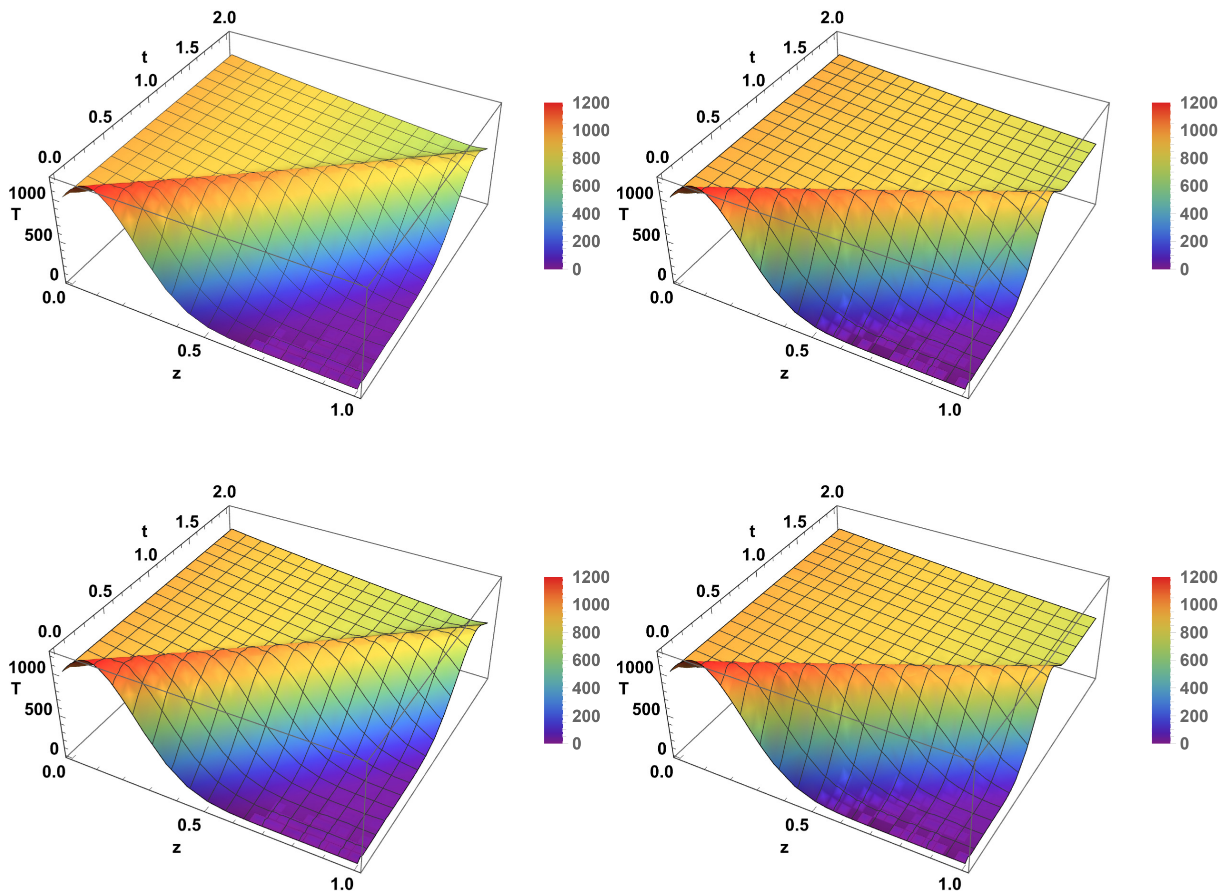

Figure 3. Numerical solution of Equation (13) with conditions (14): $d = 10^{-6}m^2s^{-1}$, $u = 10^{-4}ms^{-1}$ (top left); $d = 10^{-6}m^2s^{-1}$, $u = 1.5 \cdot 10^{-4}ms^{-1}$ (top right); $d = 2 \cdot 10^{-6}m^2s^{-1}$, $u = 10^{-4}$ (bottom left); $d = 2 \cdot 10^{-6}m^2s^{-1}$, $u = 1.5 \cdot 10^{-4}ms^{-1}$ (bottom right).

4 Conclusions

We presented a model for the upwelling of a magma batch through the lithosphere. Conceptually, the model views the process as a mix of heat propagation and buoyancy, where also heat dissipation is considered. More precisely, a batch of magma rises through the lithosphere, melting the surrounding rocks within a circular annulus of radius R, then it rises through the melted rocks by buoyancy. The process continues up to a depth at which the temperature becomes smaller than the melting temperature of the surrounding rocks. Here, the batch ends its rising and magma emplaces in a magma chamber. The model is well represented by the Equation (12) which has been derived from general balance laws (see 2.1). A numerical solution has been presented. It can model different depth of emplacement of the magma batch explaining the formation of magma chambers, whose vertical extension depends on heat propagation, buoyancy, and heat dissipation for rock partial melting. The physical

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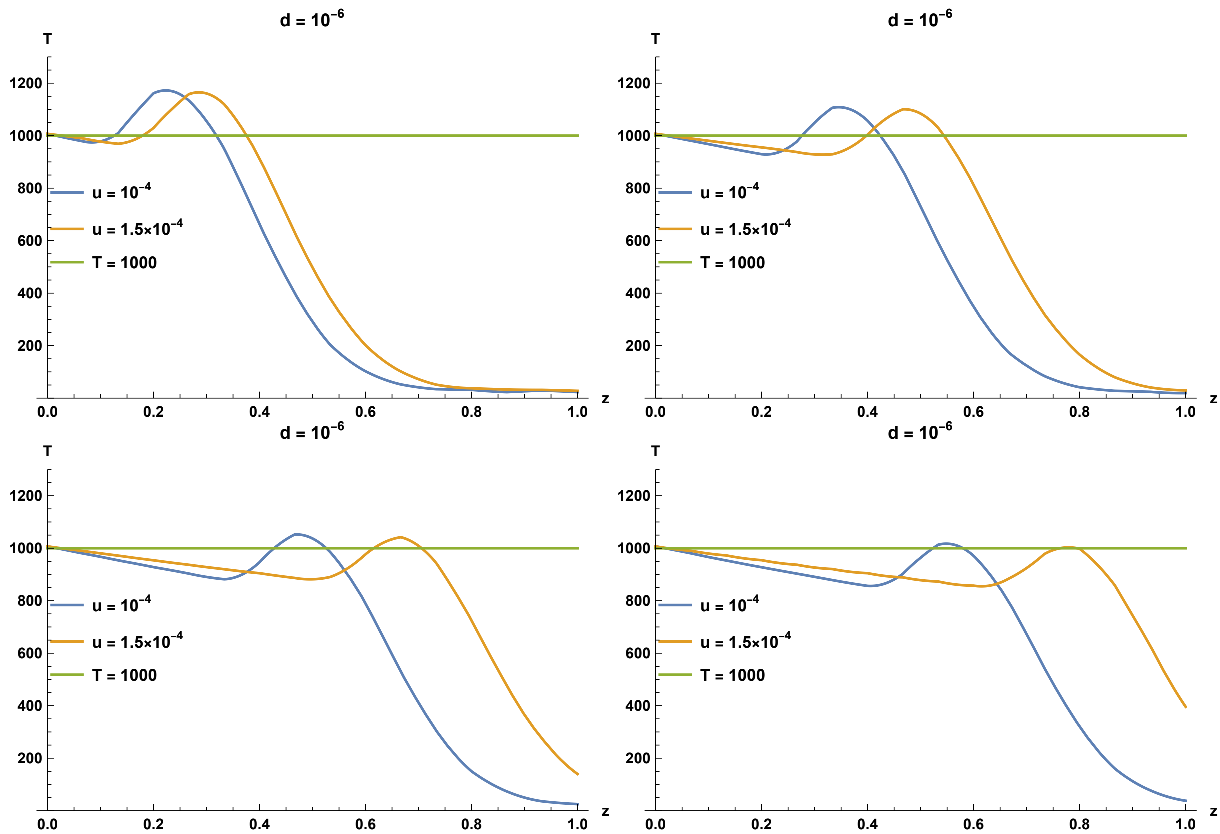

Figure 4. Profiles of the numerical solution at different times for $d = 10^{-6}$.

model can also explain the vertical compositional stratification of the crust, from the mafic lower part to the more evolved granite-dominated upper crust (Rudnick, 1995). In our model, (even if not included in Equation (12)), magma differentiation possibly occurs by assimilation (magma can incorporate and partially melt surrounding crustal rocks, altering its chemical composition) and degassing, during the rising phase of a parental magma batch. When the magma reaches the accumulation zone, where the temperature becomes smaller than the melting temperature of the surrounding rocks, the differentiation of the magma continues mainly through the process of fractional crystallization, further degassing, and magma mixing. These processes can lead to the over-pressurization of the magma chamber and to a further ascent of magma in the form of dykes within the shallower and more brittle portion of the crust. Beyond this point, magma ascent is no longer dominated by buoyancy forces. Instead, the driving forces include magma overpressure, gas and water content, and the regional or local tectonic stress field. This mechanism is dominant within the brittle crust. Magma may also migrate through pre-existing fractures, faults, and joints without necessarily forming new dykes, and thus its ascent depends on the crustal geological structure. This mechanism likewise occurs in brittle zones, particularly in highly tectonically active regions (Rubin & Pollard, 1987). Although the model does not include the composition and the chemical properties of the magmas, it represents a very general approach able to

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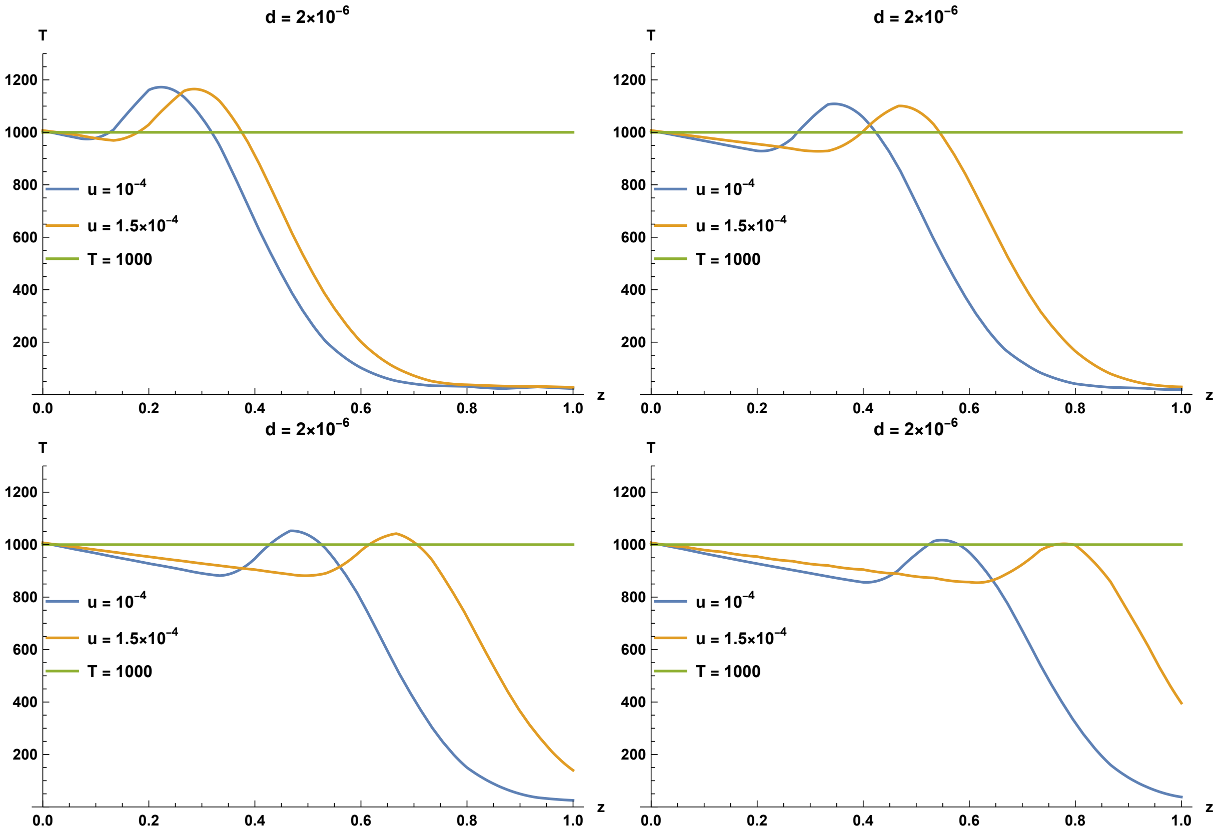

Figure 5. Profiles of the numerical solution at different times for $d = 2 \cdot 10^{-6}$.

explain the magma ascent through the lithosphere, and introduces a new perspective for the mechanism of magma ascent through the rigid lithosphere. More precisely, we introduced some oversimplifications (constant velocity and heat conductivity, $T_{tz} \gg T_{tx}$ and $T_{tz} \gg T_{ty}$) implying that we are focusing on the physical mechanism for transporting a batch of magma through the lithosphere neglecting the effects of the medium. Indeed, we are able to describe the propagation of a magma batch from the border between the lithosphere and asthenosphere up to the emplacement of a magma chamber. A more general analysis, with velocity and heat conductivity depending on temperature will be matter of a paper in preparation.

Code and data availability. The current version of Mathematica code used for the numerical integration of equation (13) is available from the project website https://zenodo.org/records/21271908 (Godano, 2026) and it does not need any license. The exact version of the model used to produce the results used in this paper is archived on repository under DOI 10.5281/zenodo.21271907 (Godano, 2026), no data and scripts to run the model and produce the plots for all the simulations presented in this paper are necessary.
No data are used in our paper.

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Appendix A: Model parameters

A1 The upwelling velocity u

The buoyancy velocity can be estimated by the following relation (Turcotte and Schubert, 2014)

$$u = \frac{2}{9}\frac{\Delta\rho}{\mu} g \mathcal{R}^2 \tag{A1}$$

where $\Delta\rho$ is the density difference between the batch of magma and the melted rocks, μ is the viscosity of the melted rocks, g is the gravity acceleration, and $\mathcal{R}$ is the radius of the magma batch. Assuming $\Delta\rho = 300$ kg m^{-3} (Turcotte and Schubert, 2014), $\mathcal{R} \in [50, 500]\ m$, roughly estimated from the distribution of the erupted volumes (Papale et al., 2021) and assuming that the erupted volumes are $\approx$ 10% of the total volume of the magma chamber (Bower and Woods, 1998), and $\mu \in [10^{-1}, 10^{14}]$ Pa (Giordano et al., 2008), we get $u \in [10^{-7}, 3 \cdot 10^{2}]$ m s^{-1}.

A2 The diffusivity d

It has been observed (Whittington et al., 2009) that the diffusivity coefficient is in the range $[5 \cdot 10^{-7}, 2 \cdot 10^{-6}]$ ms^{-1}. However, these values are obtained in laboratory experiments. As a consequence, the value of d in the real lithosphere could be very different from the measured ones in the laboratory.

A3 The relaxation time τ

For evolving magmatic systems, at depth larger than about 10 km, the lithosphere viscosity μ can be estimated to be of the order of $10^{14} - 10^{15}$ Pa (Weber et al., 2025). The viscoelastic relaxation time τ is equal to μ/G where G is the shear modulus in Pa. Thus, assuming $G = 0.5$ GPa (Weber et al., 2025), we get $\tau \in [10^5, 10^6]$ s.

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