

Magnetic fields generated by thermally, chemically, and thermochemically driven dynamos

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Solidification processes at the boundary between the Earth's inner and outer core are the primary source of buoyancy forces that drive convection and Geodynamo in the outer core. This study aims to numerically investigate the impact of different buoyancy sources on convection within the Earth's core and the resulting magnetic field. We analyze three models: thermally driven, chemically driven, and thermochemically driven dynamos. The results confirm that both the magnetic field and the convective velocity field depend on the sources of buoyancy. Additionally, dipole polarity reversals and the super-rotation of the Earth's inner core are influenced by different buoyancy sources. As the Ekman number decreases, this dependence weakens.

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I. INTRODUCTION

The magnetic fields of planets, moons, stars, and galaxies are generated by hydromagnetic dynamos operating within their liquid interiors. The basic principle of a hydromagnetic dynamo is the conversion of the kinetic energy from convective movements in the electrically conductive liquid into magnetic energy. This process not only generates magnetic fields in these celestial objects but also explains their characteristics and time evolution, aligning well with astrophysical and geophysical observations. Buoyancy serves as the primary driver of convection in the liquid parts of their interiors [1–9].

Solidification processes, radiogenic heating, gravitational differentiation, and chemical homogenization in the Earth's core are the primary sources of the buoyancy force that drives convection in the Earth's outer core and sustains the hydromagnetic dynamo [2–4,10]. Thermochemical convection is a complex problem typically addressed through numerical solutions that involve various simplifications. Most notably, instead of modeling a compressible fluid, an anelastic or Boussinesq approximation (incompressible) is often employed.

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In many numerical models, the buoyant force is primarily attributed to the temperature gradient in the Earth's core. Additional sources of buoyancy within the core include concentration gradients, codensity gradients (where thermal and concentration components of buoyancy are combined into a single variable known as codensity), separate temperature and concentration gradients, and heat flow from the inner boundary with subsequent transfer through the outer boundary [3,10–15]. It is worth noting that other models of convection and hydromagnetic dynamos have explored the role of tides or precession as potential energy sources for driving convection and the hydromagnetic dynamo [16–22]. However, in the case of Earth, these resources are only supplementary.

Super-rotation of Earth's inner core refers to its eastward rotation relative to the mantle, with a net rotation rate typically faster than that of the rest of the Earth. Geodynamo simulations by Glatzmaier and Roberts [23] and Roberts and Glatzmaier [6] suggest that the inner core rotates through a mechanism similar to an induction motor. In this model, thermal winds in the outer core drive a circulation pattern with east-west flows near the inner core boundary. Magnetic fields threading both the inner and outer cores produce a magnetic torque, while viscous forces at the boundary help synchronize the rotation rates of the inner core and the surrounding fluid. Evidence for inner core super-rotation also emerges from seismic studies, including observations of temporal changes in seismic wave travel times through the inner core and analyses of the Earth's free oscillations. However, rotation rates inferred from Geodynamo simulations and seismic observations differ. This difference is most likely due to the high viscosities assumed in numerical models and the inclusion or omission of gravitational coupling [3,6,23–26].

Kutzner and Christensen [27] studied three-dimensional, convectively driven dynamos, where the driving convection was either thermal or chemical in nature. For both types, they considered cases with and without internal heating. Regarding boundary conditions, they investigated scenarios with either fixed temperature or concentration, or fixed temperature or concentration fluxes at the boundaries. Their results showed that for low Rayleigh numbers, the generated magnetic field is characterized by a dominant, stable axial dipole. For high Rayleigh numbers, the dipole component weakens, and a multipolar magnetic field dominates. During the transition between these two regimes, the dynamo oscillates between dipole-dominated and multipolar states, with excursions and reversals of the dipole observed. For fixed temperatures or concentrations at the boundaries, this transition is sharp, whereas with imposed fluxes (at one or both boundaries) or with internal sources or sinks, the transition is more gradual.

The diversity of numerical models and the complexity of buoyancy sources have motivated us to examine thermally, chemically, and thermochemically driven hydromagnetic dynamos using a consistent set of parameters. We will focus on how these different sources of buoyancy influence the nature of convection and the characteristics of the generated magnetic fields, including their polarity reversals. In contrast to Kutzner and Christensen [27], we investigate not only thermally and chemically driven dynamos, but also thermochemically driven dynamos. The range of Ekman numbers explored is comparable to Ref. [27]. In addition to examining the characteristics of the velocity and magnetic fields, as well as excursions and reversals, we also analyze how different Ekman numbers and convection types influence the inertial and Lorentz forces, the energy balance, and inner core rotation—factors not addressed by Kutzner and Christensen [27], but essential for understanding the dynamics of the Earth's core.

The structure of the paper is as follows: Section II describes the model and governing equations. Section III details the numerical method and input/output parameters. Our findings are presented in Sec. IV, followed by the discussion and conclusions in Sec. V.

II. MODEL AND BASIC EQUATIONS

Models based on temperature and concentration gradients, and codensity for rotating convection in a spherical shell are being considered, following [13,27–34]. An incompressible, electrically conducting fluid occupies the spherical shell $r_i < r < r_o$ (where r_i and r_o are the inner and outer core radii, respectively), rotating with an angular velocity $\mathbf{\Omega} = \Omega \mathbf{k}$. The standard spherical-coordinate

system (r, θ, φ) is used, where $\mathbf{k}$ denotes the unit vector in the vertical direction (aligned with the rotation axis). Under the Boussinesq approximation, the evolution of the velocity field $\mathbf{U}$, the magnetic field $\mathbf{B}$, and the buoyancy source $\mathcal{A}$ is governed by a system of dimensionless equations

$$E \frac{D\mathbf{U}}{Dt} + 2\mathbf{k} \times \mathbf{U} = -\nabla P + R_a^{\mathcal{A}} \frac{\mathbf{r}}{r_o} \mathcal{A} + E \nabla^2 \mathbf{U} + \frac{1}{P_m} (\nabla \times \mathbf{B}) \times \mathbf{B}, \quad (1)$$

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{U} \times \mathbf{B}) + \frac{1}{P_m} \nabla^2 \mathbf{B}, \quad (2)$$

$$\frac{D\mathcal{A}}{Dt} = \frac{1}{P_r} \nabla^2 \mathcal{A} + S_C, \quad (3)$$

$$\nabla \cdot \mathbf{U} = 0, \quad (4)$$

$$\nabla \cdot \mathbf{B} = 0, \quad (5)$$

where D/Dt denotes the material derivative with respect to the flow velocity $\mathbf{U}$. We investigate three models. The first model (model $\mathcal{T}$) is a thermally driven dynamo, where $\mathcal{A} = T$ and $S_C = 0$, with T representing the temperature. The temperature field is defined as the deviation from the isentropic temperature under the Boussinesq approximation. In comparison to the third model, concentration effects are disabled in this model. The second model (model $\mathcal{CH}$) is a chemically driven dynamo, where $\mathcal{A} = \Xi$ and $S_C = -3/(r_o^3 - r_i^3)$, with Ξ representing the concentration of the light elements. In comparison to the third model, thermal effects are disabled in this model. Finally the third model (model $\mathcal{TCH}$) is a thermochemically driven dynamo, with $\mathcal{A} = C$, where C represents codensity. In this model, both thermal and concentration effects are active.

To render the governing equations dimensionless, we use the shell gap thickness $L = r_o - r_i$ as a length scale, the viscous timescale L^2/ν , the magnetic induction scale $(\rho\mu\eta\Omega)^{1/2}$, the temperature scale ΔT , where ΔT is the temperature drop across the spherical shell, the concentration scale $\Delta \Xi$, where $\Delta \Xi$ is the concentration of the light elements drop across the spherical shell, the codensity scale $F/4\pi\nu L$ and the pressure scale $\rho\nu^2/L^2$. Here, ν is the kinematic viscosity, μ is the magnetic permeability, η is the magnetic diffusivity, and F the total codensity flux. The dimensionless parameters appearing in Eqs. (1)–(5) are the Prandtl number, P_r , the magnetic Prandtl number, P_m , the Ekman number, E , and the Rayleigh number, $R_a^{\mathcal{A}}$, defined as

$$P_r = \frac{\nu}{\kappa}, \quad P_m = \frac{\nu}{\eta}, \quad E = \frac{\nu}{\Omega L^2}, \quad \text{and} \quad R_a^{\mathcal{A}} = R_a^{\mathcal{T}}, R_a^{\mathcal{CH}}, R_a^{\mathcal{TCH}}, \quad (6)$$

where κ is the thermal diffusivity and $R_a^{\mathcal{T}}$, $R_a^{\mathcal{CH}}$, and $R_a^{\mathcal{TCH}}$ are the Rayleigh numbers for models $\mathcal{T}$, $\mathcal{CH}$, and $\mathcal{TCH}$, respectively, defined as

$$R_a^{\mathcal{T}} = \frac{\alpha_T g_0 \Delta T L}{\nu \Omega}, \quad R_a^{\mathcal{CH}} = \frac{\alpha_{\Xi} g_0 \Delta \Xi L}{\nu \Omega}, \quad R_a^{\mathcal{TCH}} = \frac{F g_0}{4\pi \rho \nu^2 \Omega}, \quad (7)$$

where α_T is the coefficient of thermal expansion, α_{Ξ} is the coefficient of concentration expansion, ρ is the density, and g_0 is the acceleration due to gravity at $r = r_o$. Note that κ is the same for thermal and concentration diffusion.

In model $\mathcal{T}$, the fixed-temperature boundary conditions are applied, with $T_i = 1$ at the inner boundary and $T_o = 0$ at the outer boundary of the shell. In model $\mathcal{CH}$, a fixed-concentration boundary condition is applied at the inner boundary of the shell with $\Xi_i = 1$, while a fixed-concentration flux boundary condition is applied at the outer boundary with $\partial \Xi_o / \partial r = 0$. Model $\mathcal{TCH}$ is based on codensity. Since Aubert *et al.* [32], Šimkanin *et al.* [33,34] provide a detailed description of the codensity model, only the fundamental principles are briefly outlined here. Codensity combines temperature and composition, assuming equal effective diffusivity for both variables. The boundary conditions for C are specified in terms of fluxes. At the inner-core boundary, the release of latent heat and light elements acts as the chemical sources of convection, generating a positive mass anomaly flux F_i . Conversely, at the outer boundary, the mass anomaly flux imposed by the mantle corresponds

TABLE I. Input parameters used in the simulations.

E	P_r	P_m	R_a^A/Ra_c^A	f_i	N_r	N_θ	N_φ
10^{-2}	1	155	12	0.85	175	113	225
10^{-3}	1	25	4.5	0.85	225	150	300
10^{-4}	1	4.5	4.5	0.85	313	216	432

to the heat flux carried out from the outer core, representing the secular cooling of the outer core. There are no chemical contributions to F_o at the outer boundary, so the light-element mass flux is set to zero there. The values of F_i and F_o are specified as follows:

$$F_i = - \int_{S_i} \kappa \nabla C \cdot d\mathbf{S} \quad \text{and} \quad F_o = - \int_{S_o} \kappa \nabla C \cdot d\mathbf{S}, \quad (8)$$

where S_i and S_o are the inner and outer boundary surfaces, respectively, and $d\mathbf{S}$ is the surface element oriented radially outwards. Note that F_i corresponds to the flux of the codensity field into the outer core at the inner boundary, while F_o corresponds to the flux of the codensity field from the outer core at the outer boundary. The buoyancy ratio associated with the inner boundary is defined as

$$f_i = \frac{F_i}{F}, \quad (9)$$

such that $F = F_i + F_o$. The volumetric source term in Eq. (3) ensures that the averaged codensity does not change over time, see Kutzner and Christensen [27], and is defined as

$$S_C = 3 \frac{1 - 2f_i}{r_o^3 - r_i^3}. \quad (10)$$

Note that the model $\mathcal{T}$ corresponds to $f_i = 0.5$ and the model $\mathcal{CH}$ to $f_i = 1$. Finally, Eqs. (1)–(5) are closed by nonpenetrating boundary conditions for the velocity field at both boundaries, the no-slip condition at the outer boundary and the stress-free condition at the inner boundary. The outer boundary is electrically insulating, ensuring that the magnetic field at this boundary matches an appropriate potential field in the exterior. The inner boundary is electrically conducting, with the electrical conductivity assumed to be the same for both the outer and inner cores.

III. NUMERICAL METHOD AND INPUT/OUTPUT PARAMETERS

The system (1)–(5) is solved numerically using the PARODY dynamo code [12, 35–37]. This semispectral code combines the spherical harmonics decomposition in the lateral directions with a second-order finite-difference scheme in the radial direction, and employs an implicit second-order Crank-Nicolson scheme for time integration. We use adaptive time stepping, and parallelization is achieved through the message-passing interface (MPI). To visualize the results, we utilize MATLAB scripts provided by Julien Aubert [12].

The computations are initiated with a perturbation of a zero velocity field, a strong dipole-dominated magnetic field, and an initial codensity field given by the radially symmetric steady state [33]. The input parameter values used in the computations are listed in Table I. In all simulations, the spatial resolution is set to $(N_r \times N_\theta \times N_\varphi)$, representing the number of grid points in the respective spatial directions. The Prandtl number is set to $P_r = 1$, and the buoyancy ratio is set to $f_i = 0.85$, a value considered close to those relevant to Earth’s core. We set $R_a^{\mathcal{T}}$, $R_a^{\mathcal{CH}}$, and $R_a^{\mathcal{TC\mathcal{H}}}$ so that the R_a^A/Ra_c^A ratio remains fixed for a given E (see Table I). Note that the definitions of the Rayleigh numbers differ for each model [see Eq. (7)], and consequently, their critical values $\text{Ra}_c^{\mathcal{T}}$, $\text{Ra}_c^{\mathcal{CH}}$, $\text{Ra}_c^{\mathcal{TC\mathcal{H}}}$ (corresponding to the dynamo onset), also vary, as described in more detail by Zhang and Busse [38] and Busse and Simitev [5]. Our decision to maintain the same R_a^A/Ra_c^A ratio ($R_a^{\mathcal{T}}/\text{Ra}_c^{\mathcal{T}} = R_a^{\mathcal{CH}}/\text{Ra}_c^{\mathcal{CH}} = R_a^{\mathcal{TC\mathcal{H}}}/\text{Ra}_c^{\mathcal{TC\mathcal{H}}}$) for a given E across all three models was driven

by the goal of isolating the effect of changing the type of buoyancy source (switching among models $\mathcal{T}$, $\mathcal{CH}$, and $\mathcal{TCH}$), while keeping all other parameters constant. Higher values of R_a^A/Ra_c^A and P_m are preferred for numerical modeling, as Kutzner and Christensen [27] found that dynamos capable of generating dipole-dominant magnetic fields with polarity reversals tend to occur at higher values of R_a^A and P_m .

The outputs of the numerical simulations are presented in terms of global characteristics. The influence of the inertial and Lorentz forces on the hydromagnetic dynamo is assessed using the modified Rossby number, Ro_ℓ , and the modified Elsasser number, Λ_ℓ , defined as

$$\text{Ro}_\ell = R_o \frac{\ell_U}{\pi}, \quad \Lambda_\ell = \frac{\Lambda L}{R_m \ell_B}, \quad (11)$$

where $R_o = E(2E_k)^{1/2}$ is the Rossby number, ℓ_U is the characteristic convective length scale, $\Lambda = 2P_m E E_m$ is the Elsasser number, ℓ_B is the characteristic magnetic dissipation length scale, $R_m = P_m(2E_k)^{1/2}$ is the magnetic Reynolds number, E_k is the kinetic energy density, and E_m is the magnetic energy density [28]. ℓ_U is calculated from the time-averaged kinetic energy spectrum, and ℓ_B is obtained by integrating over the entire volume of the spherical shell

$$\ell_U = \frac{1}{2E_k} \sum_l l \langle \mathbf{U}_l \cdot \mathbf{U}_l \rangle, \quad \ell_B^2 = \frac{\int \mathbf{B}^2 d^3\mathbf{x}}{\int (\nabla \times \mathbf{B})^2 d^3\mathbf{x}}, \quad (12)$$

where $\mathbf{U}_l$ is the velocity component at harmonic degree l and the angle brackets denote the average over the spherical shell and over a sufficient time interval. The definition of ℓ_U assumes that the radial and horizontal length scales are approximately similar. Comprehensive explanations for calculating ℓ_U and ℓ_B can be found in Christensen and Aubert [13], Dormy [39], Dormy *et al.* [8], and Šimkanin and Kyselica [31]. The Rossby number measures the ratio of inertial to Coriolis forces; however, Christensen and Aubert [13] introduced a modified Rossby number that depends on the characteristic convective length scale rather than the shell thickness. This modification provides a more accurate measure of the balance between inertial and Coriolis forces. In the strong-field dynamo regime, Λ_ℓ is of order unity, whereas in the weak-field dynamo regime, $\Lambda_\ell \ll 1$. The modified Elsasser number offers a more appropriate measure of the balance between Lorentz and Coriolis forces [8,39].

To determine the energy balance, we calculate the power generated by buoyancy forces (buoyant power) as

$$\text{PB} = \frac{R_a^A}{r_o} \int \mathcal{A} \mathbf{r} \cdot \mathbf{U} d^3\mathbf{x}, \quad (13)$$

and the Ohmic and viscous dissipations, respectively, as

$$Q_\eta = \frac{1}{P_m^2} \int |\nabla \times \mathbf{B}|^2 d^3\mathbf{x} \quad \text{and} \quad Q_\nu = E \int |\nabla \times \mathbf{U}|^2 d^3\mathbf{x}. \quad (14)$$

PB represents the buoyant power input required to sustain the dynamo. In a statistically steady state, this must be balanced by the total dissipation $Q = Q_\nu + Q_\eta$, meaning $\text{PB} = Q$ [6,10,32], as confirmed by all the present simulations. In addition to PB, we present plots of Q_ν/PB and Q_η/PB , which represent the fractions of the total dissipation attributed to Ohmic and viscous dissipation, respectively.

The inner core is free to rotate due to the electromagnetic, viscous, and gravitational coupling, which may result in a difference between the rotation rate of the inner core and the modeled mantle [6,23,40]. The angular momentum balance is described by the angular momentum equation, which in its dimensional form is expressed as

$$I \frac{d\omega_i}{dt} = \Gamma_\eta + \Gamma_\nu + \Gamma_g, \quad (15)$$

where I is the moment of inertia of the inner core, ω_i is the inner core's rotation rate, Γ_η is the magnetic torque, Γ_v is the viscous torque, and Γ_g is the gravitational torque. Aubert and Dumberry [41] incorporated this equation into their model in the following form:

$$\frac{8}{15}\rho\pi r_i^5 \frac{d\omega_i}{dt} = \rho\nu \int_{\text{ICB}} s^2 \frac{\partial U_\varphi/s}{\partial r} dS + \frac{r_i}{\mu} \int_{\text{ICB}} \mathbf{B}_r \mathbf{B}_\varphi \sin \theta dS - \Gamma_v \omega_i, \quad (16)$$

where τ_v represents the viscous relaxation time, and s denotes the cylindrical radius. The PARODY dynamo code solves Eq. (16) to determine the evolution of ω_i . Seismic observations estimated ω_i to range between 0.25° and 0.48° per year, consistent with super-rotation [3,6,23]. However, numerical dynamo simulations often yield ω_i values ranging from 2° to 3° per year [3,6,23]. This discrepancy may result from the use of unrealistically high Ekman numbers in numerical models, as shown by Glatzmaier and Roberts [23], and Roberts and Glatzmaier [6]. In addition, Buffett and Glatzmaier [24] demonstrated that steady inner core rotation can be efficiently inhibited by gravitational coupling, yielding values of $\omega_i \sim 0.02^\circ$ per year. These findings were later corroborated by further Geodynamo simulations conducted by Aubert and Dumberry [41]. Finally, based on seismic data analyses, Witze [25] and Yang and Song [26] reported that the inner core rotation rate varies over time (see Sec. V).

IV. NUMERICAL RESULTS

Numerical simulations were conducted using the parameters listed in Table I, with runs of 2000 time units for $E = 10^{-2}$ and 19 time units for $E = 10^{-3}$ and 10^{-4} . Figure 1 presents the temporal evolution of kinetic energy density, E_k , and magnetic energy density, E_m , as defined by Christensen *et al.* [28]. For $E = 10^{-2}$, thermochemical convection strongly dominates over both chemical and thermal convection [see Fig. 1(a)]. In contrast, for $E = 10^{-3}$ and 10^{-4} , thermal convection becomes the dominant mechanism, surpassing both chemical and thermochemical convection [see Figs. 1(b) and 1(c)]. The values of E_m for thermally driven dynamos are consistently higher than those for chemically and thermochemically driven dynamos across all investigated values of E . For $E = 10^{-3}$ and 10^{-4} , the E_m values for models $\mathcal{T}$ and $\mathcal{TCH}$ are comparable [see Figs. 1(d)–1(f)]. We do not present the toroidal and poloidal components of kinetic and magnetic energy here. However, in general, the toroidal component dominates over the poloidal component for both kinetic and magnetic energy, except in one case: for model $\mathcal{T}$ at $E = 10^{-2}$, where the poloidal kinetic energy is the dominant component.

The spatial distributions of the radial magnetic field component B_r are shown in Fig. 2. The generated magnetic field exhibits a dipolar character in most cases, except for the $\mathcal{TCH}$ model at $E = 10^{-2}$, where the magnetic field is dipolar but concentrated primarily in the polar regions (a phenomenon related to inertial convection), similar to the findings in Šimkanin *et al.* [33]. However, the power spectra of the magnetic field, presented in Figs. 3(a)–3(c), reveal that for $E = 10^{-2}$, the dipole-dominant dynamo occurs only in the $\mathcal{TCH}$ model, while the dynamos in the $\mathcal{T}$ and $\mathcal{CH}$ models generate a quadrupole-dominant magnetic field [see Fig. 3(a)]. At $E = 10^{-3}$ and $E = 10^{-4}$, all three models exhibit a dipole-dominant magnetic field [see Figs. 3(b) and 3(c)].

The velocity field is primarily dominated by the $l = 2$ and $l = 3$ ($\mathcal{T}$ model) modes at $E = 10^{-2}$, by the $l = 3$ and $l = 4$ ($\mathcal{CH}$ model) modes at $E = 10^{-3}$, and by the $l = 3$ ($\mathcal{T}$ model), $l = 4$ ($\mathcal{CH}$ model) and $l = 6$ ($\mathcal{TCH}$ model) modes at $E = 10^{-4}$ [see Figs. 3(d)–3(f)]. This is consistent with the expectation, consistent with many other studies, that, as the Ekman number decreases, smaller scales begin to dominate.

Equatorial sections of T , Ξ and C are shown in Fig. 4, while Fig. 5 presents equatorial sections of the radial velocity field V_r . Convection occurs throughout the entire volume and generally exhibits a large-scale columnar structure, except in the $\mathcal{TCH}$ model at $E = 10^{-2}$, where the convection corresponds to a transition regime between the inertial and large-scale columnar convection. This transition, along with inertial convection, has been explored in greater detail by Zhang and Busse [38] and Šimkanin *et al.* [33]. A threefold symmetry in large-scale columnar convection is typical at

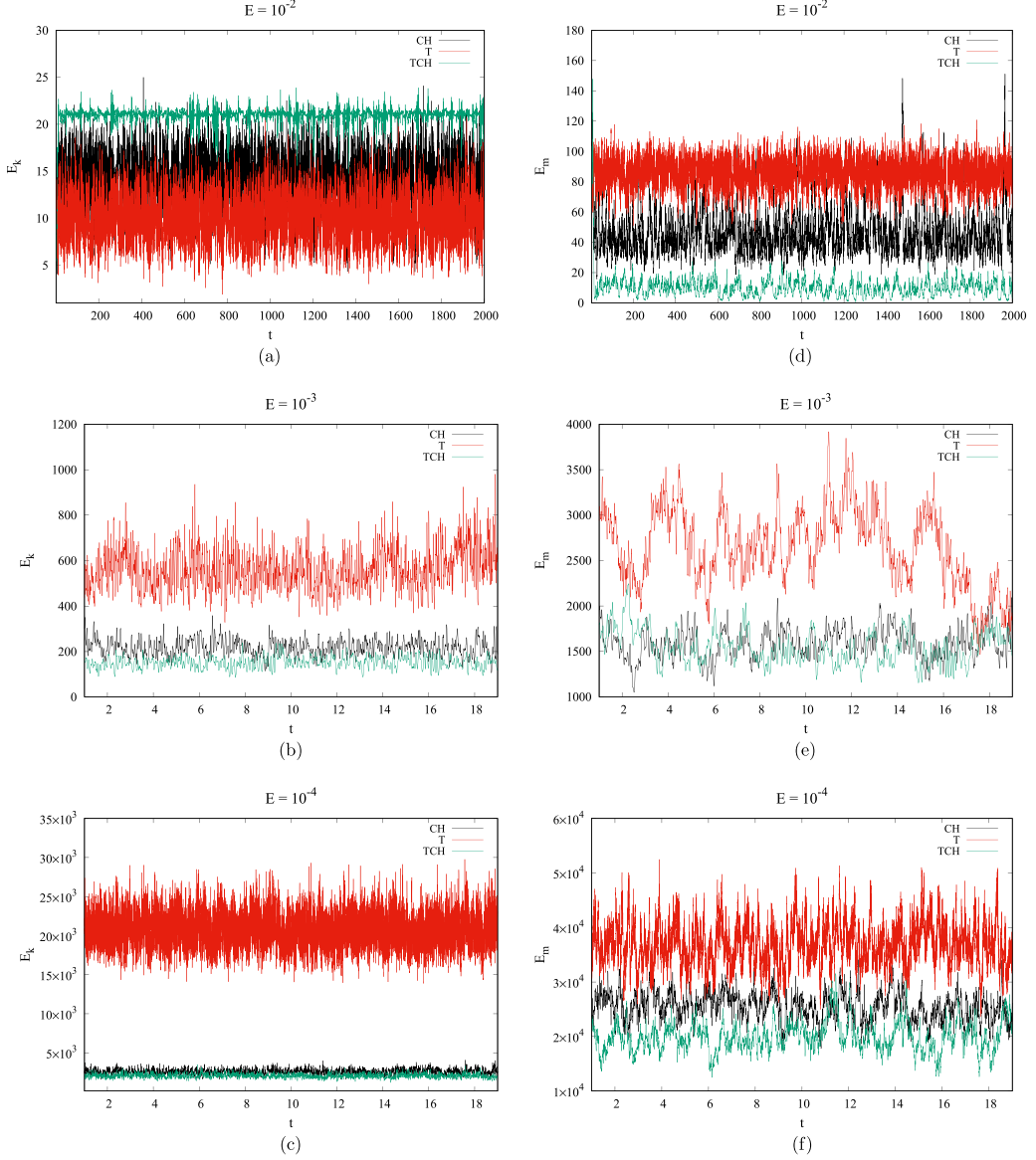


FIG. 1. Temporal evolution of kinetic energy density, E_k (left column), and magnetic energy density, E_m (right column), as a function of the model type and the Ekman number, E .

$E = 10^{-2}$ (except in the previously mentioned $\mathcal{TCH}$ model), while fourfold symmetry is characteristic at $E = 10^{-3}$. Finally, at $E = 10^{-4}$, a fivefold symmetry emerges, with smaller scales becoming much more prominent compared to higher E values (see Figs. 4 and 5).

The time evolution of the dipole tilt is shown in Fig. 6. The dipole tilt refers to the angle between the magnetic dipole axis and the Earth's rotation axis. At $E = 10^{-2}$, the dipole component of the generated magnetic field is quite unstable for the $\mathcal{CH}$ and $\mathcal{TCH}$ models, resulting in frequent polarity reversals [see Fig. 6(a)]. In contrast, the dipole component in the $\mathcal{T}$ model is more stable, with no polarity reversals observed, only frequent excursions, as shown in Fig. 6(a). At $E = 10^{-3}$ and $E = 10^{-4}$, the dipole component of the generated magnetic field remains stable, and no polarity

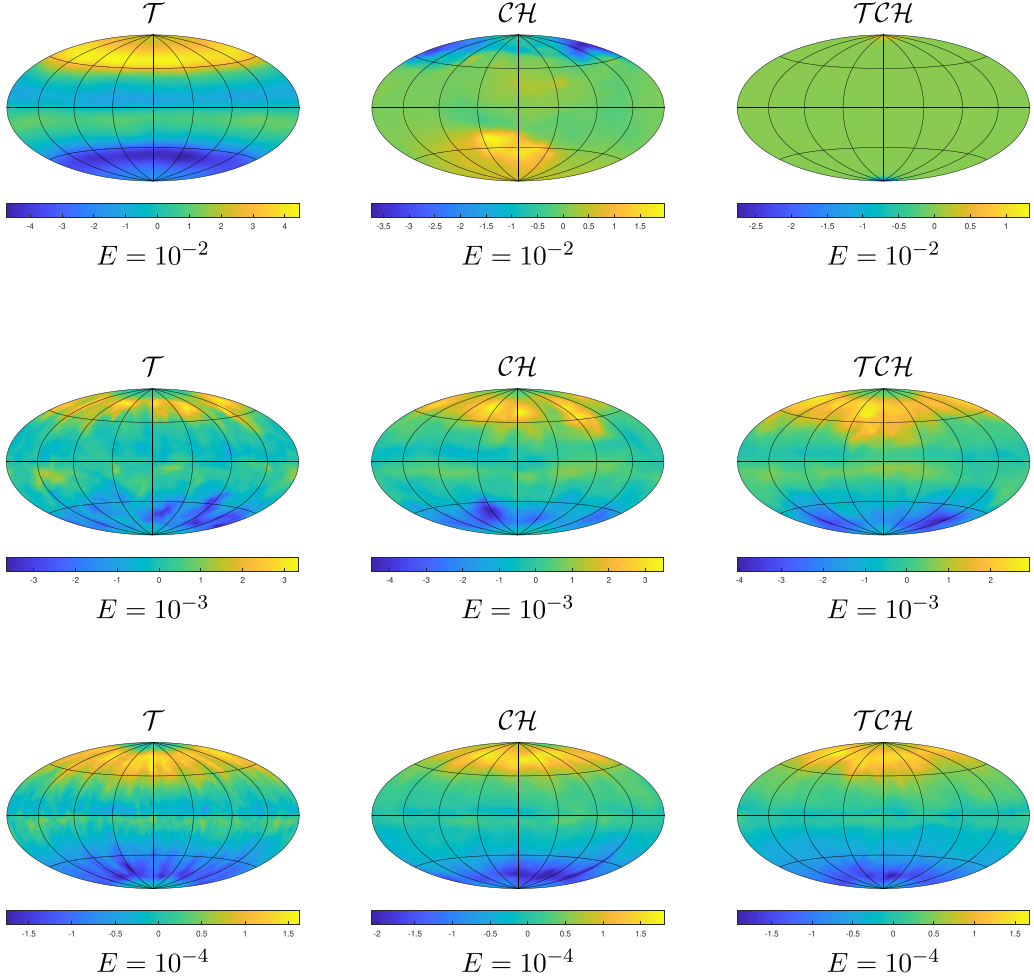


FIG. 2. Spatial distribution (Hammer projection) of radial magnetic field component, B_r , at $r = r_o$, as a function of the model type and the Ekman number, E . The orange (blue) colors indicates positive (negative) values. The snapshots are taken at $t = 2000$ for $E = 10^{-2}$, and $t = 19$ for $E = 10^{-3}$ and 10^{-4} .

reversals are observed [see Figs. 6(b) and 6(c)]. At $E = 10^{-3}$, excursions of the dipole component are observed, but these excursions are weak at $E = 10^{-4}$ compared to the $E = 10^{-3}$ case. Note that polarity reversals at $E = 10^{-3}$ and $E = 10^{-4}$ might be observable over much longer time series than those used in this analysis, and possibly under different parameters than those considered here.

Figure 7(a) shows the dependence of the time-averaged modified Rossby number (Ro_ℓ) on the model type. The Ro_ℓ values increase with rising E across all investigated models, indicating that inertia becomes more significant at higher Ekman numbers [see Fig. 7(a)]. Additionally, the dependence of Ro_ℓ on the model type is notable [Fig. 7(a)].

Figure 7(b) illustrates the dependence of the time-averaged modified Elsasser number (Λ_ℓ) on the model type. Λ_ℓ values generally increase with E for all models except for the TCH model at $E = 10^{-2}$ [see Fig. 7(b)]. The dependence of Λ_ℓ on the model type is apparent, particularly at higher Ekman numbers, while it is weaker at lower values. Since $\Lambda_\ell > 1$, the dynamos in our analysis are classified as strong-field dynamos, with one exception: the TCH model at $E = 10^{-2}$, where a weak-field dynamo ($\Lambda_\ell < 1$) is observed [see Fig. 7(b)].

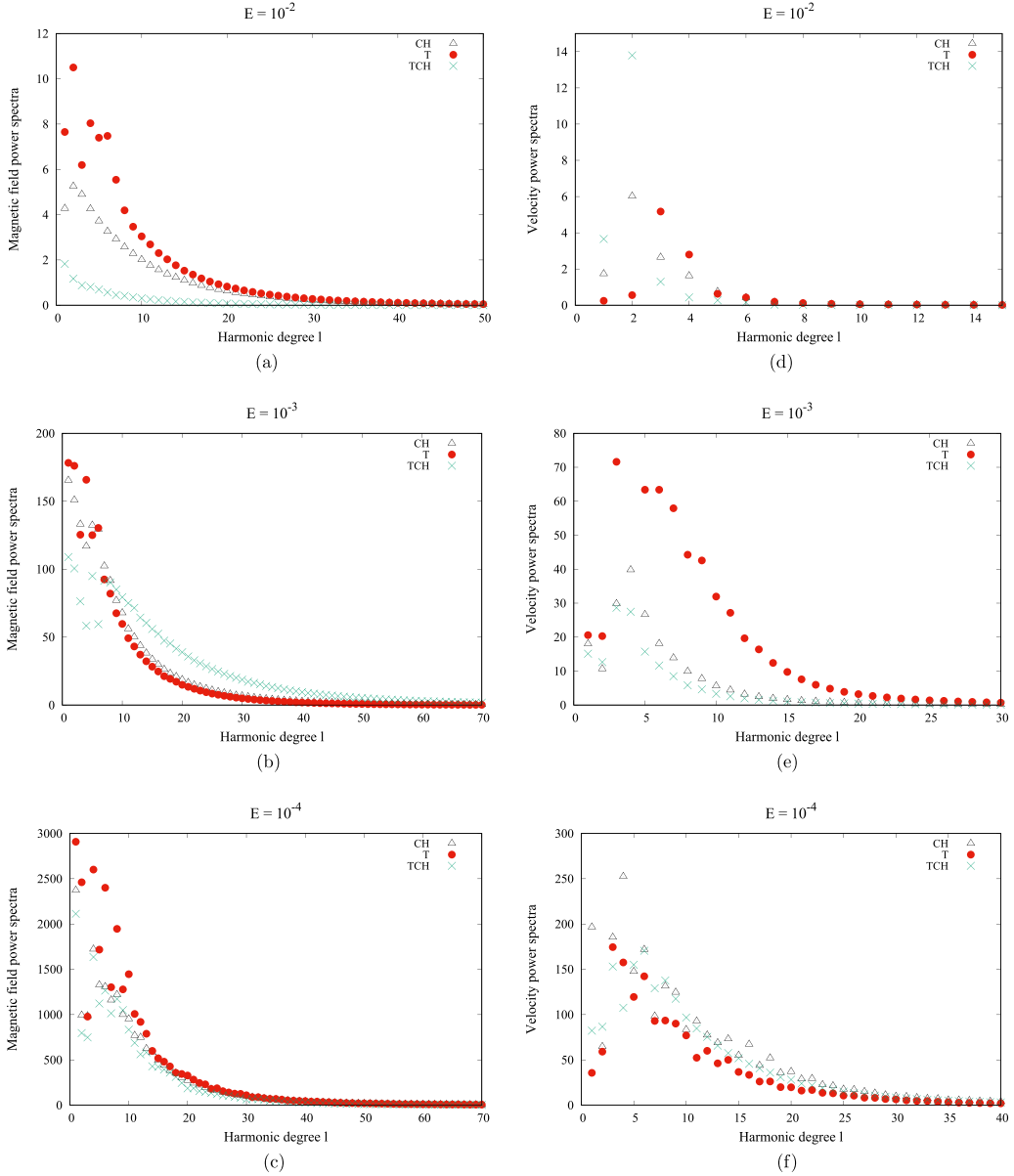


FIG. 3. Power spectra of the magnetic field (left column) at $r = r_o$ and the velocity field (right column) $r = 0.85$ as functions of spherical harmonic degree l , the model type, and the Ekman number, E . For each l , the sum over all values of m is shown. The axis for the magnetic field power spectra is in the log scale.

The power generated by buoyancy forces (PB) is shown in Fig. 8(a). PB increases as E decreases, with a noticeable dependence on the model type across all considered values of E . The ratios of ohmic and viscous dissipation to buoyancy power (Q_η/PB and Q_v/PB) are presented separately in Figs. 8(b) and 8(c). Q_η/PB decreases with increasing E , while Q_v/PB increases with increasing E . The dependence of both ratios on the model type is clearly evident [see Figs. 8(b) and 8(c)].

Seismic observations indicated that Earth's inner core is in super-rotation—meaning it rotates faster than the mantle (in our context, the upper boundary). This phenomenon was also reproduced

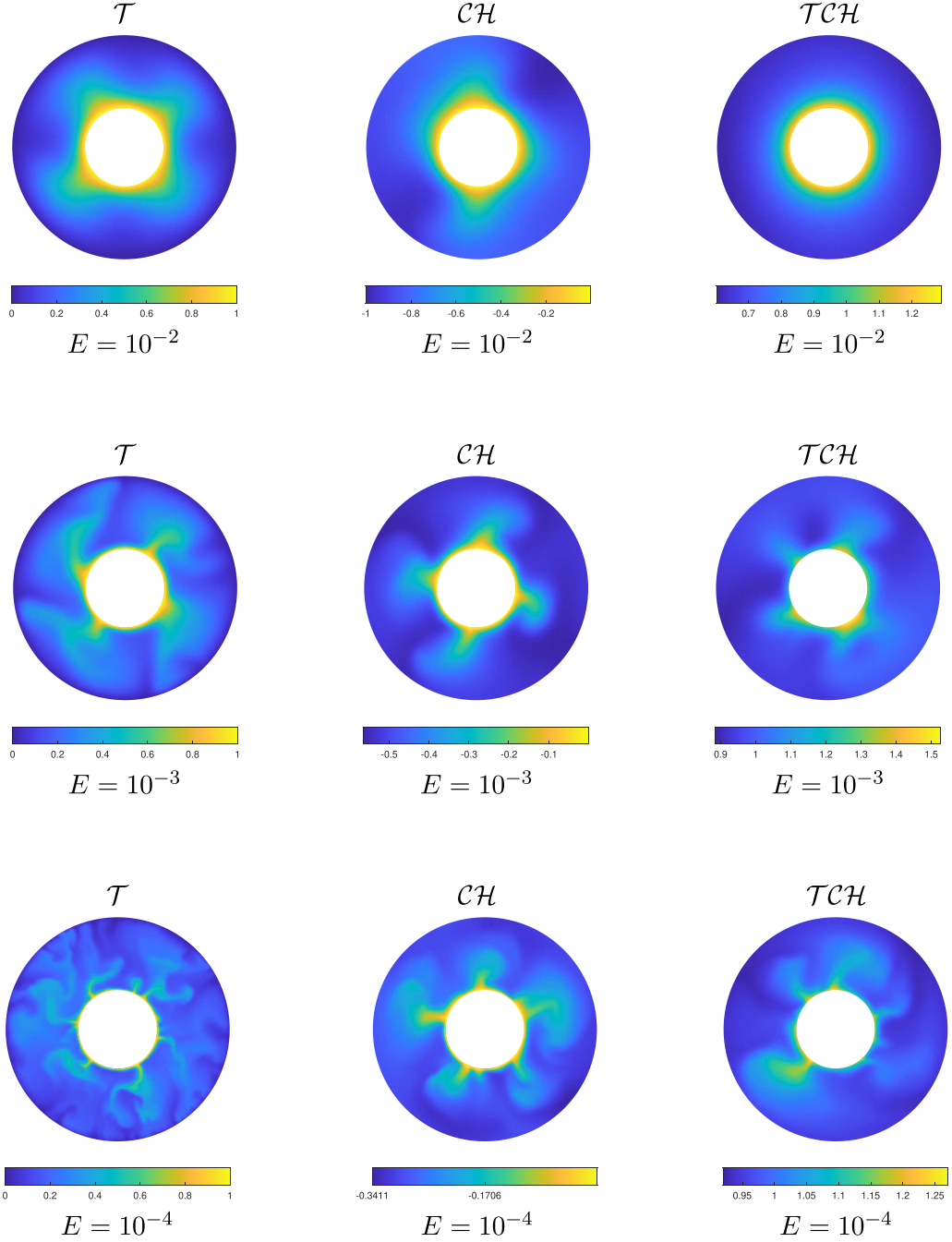


FIG. 4. Equatorial sections of temperature, T (left column), concentration of the light elements, Ξ (middle column), and codensity, C (right column), as functions of the Ekman number, E . The snapshots are taken at $t = 2000$ for $E = 10^{-2}$, and $t = 19$ for $E = 10^{-3}$ and 10^{-4} .

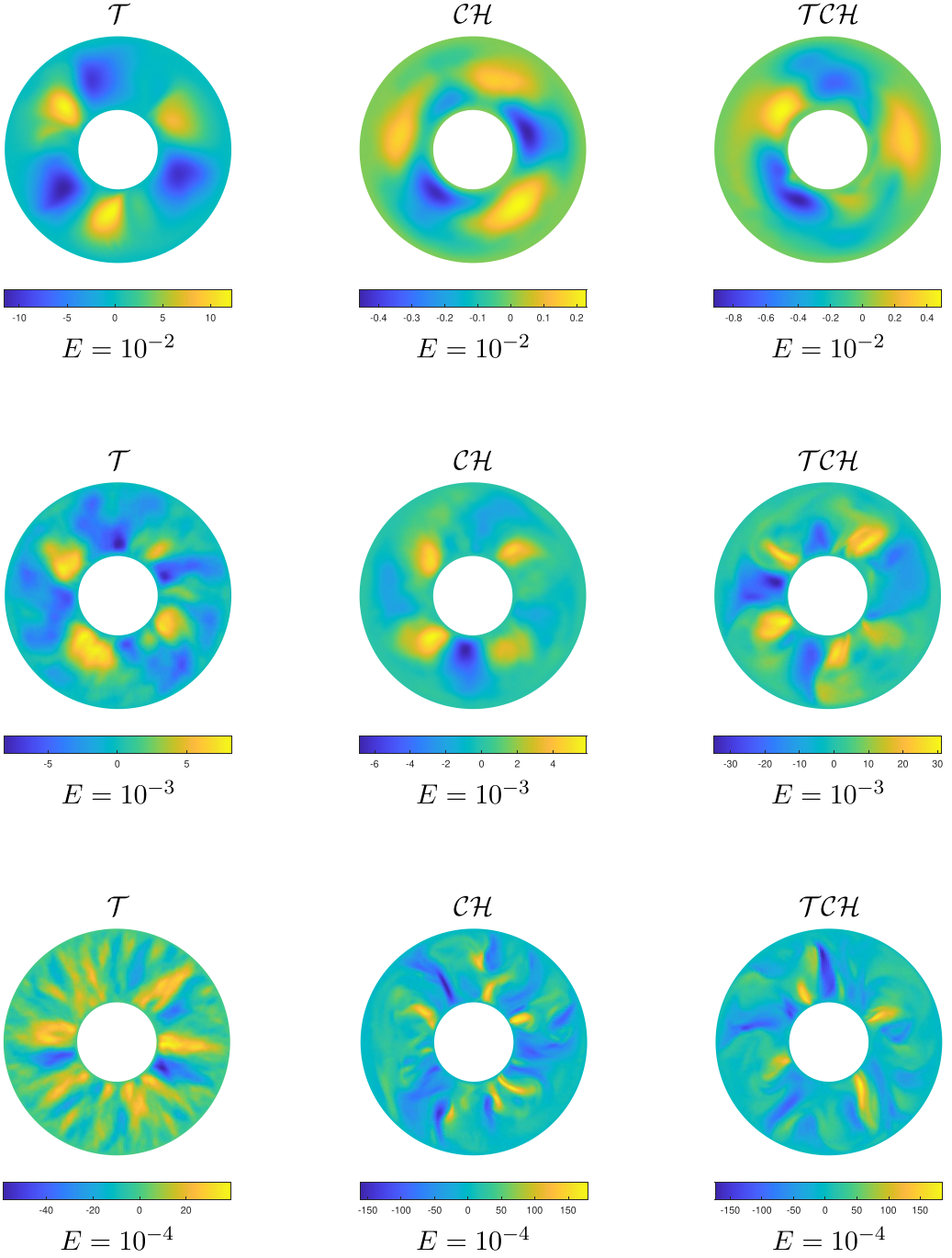


FIG. 5. Equatorial sections of velocity fields, U_r , as a function of the model type and the Ekman number, E . The orange (blue) colors indicates positive (negative) values. The snapshots are taken at $t = 2000$ for $E = 10^{-2}$, and $t = 19$ for $E = 10^{-3}$ and 10^{-4} .

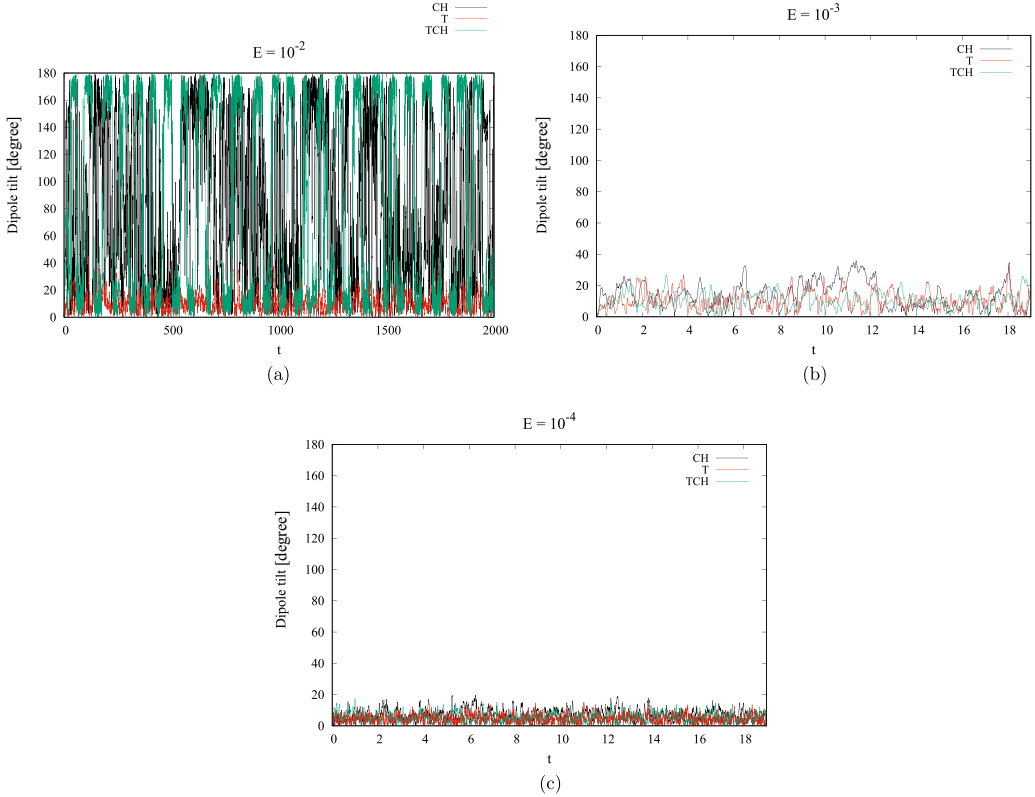


FIG. 6. Temporal evolution of dipole tilt as a function of the model type and the Ekman number, E .

in the first numerical Geodynamo simulations [6,23]. Figure 9 shows the dependence of the time-averaged inner core rotation rate (ω_i) on the Ekman number. The rotation rate is dimensional and expressed in degrees per year ($[\text{deg}/\text{year}]$). In all cases investigated here, only super-rotation is observed; subrotation or intervals with no inner core rotation are not seen (see Fig. 9). The values

TABLE II. Inner core rotation rate ω_i [degrees/year] versus Ekman number, with and without gravitational torque (Γ_g), from dynamo simulations.

Authors	E	Γ_g	ω_i in [degrees/year]
This study	10^{-2}	Yes	22–429
This study	10^{-3}	Yes	6–21
This study	10^{-4}	Yes	1.6–2.2
Aubert and Dumberry [41]	3×10^{-5}	No	0.67–1.36
Aubert and Dumberry [41]	3×10^{-5}	Yes	0.0037–0.54
Pichon <i>et al.</i> [42]	3×10^{-5}	Yes	0.1–0.75
Aubert and Dumberry [41]	10^{-5}	No	0.85–0.97
Aubert and Dumberry [41]	3×10^{-6}	No	0.84–0.93
Aubert and Dumberry [41]	3×10^{-6}	Yes	0.011–0.87
Glatzmaier and Roberts [23]	10^{-6}	No	2–3
Buffett and Glatzmaier [24]	10^{-6}	Yes	0.02
Šimkanin and Kyselica [43]	10^{-7}	No	1.5

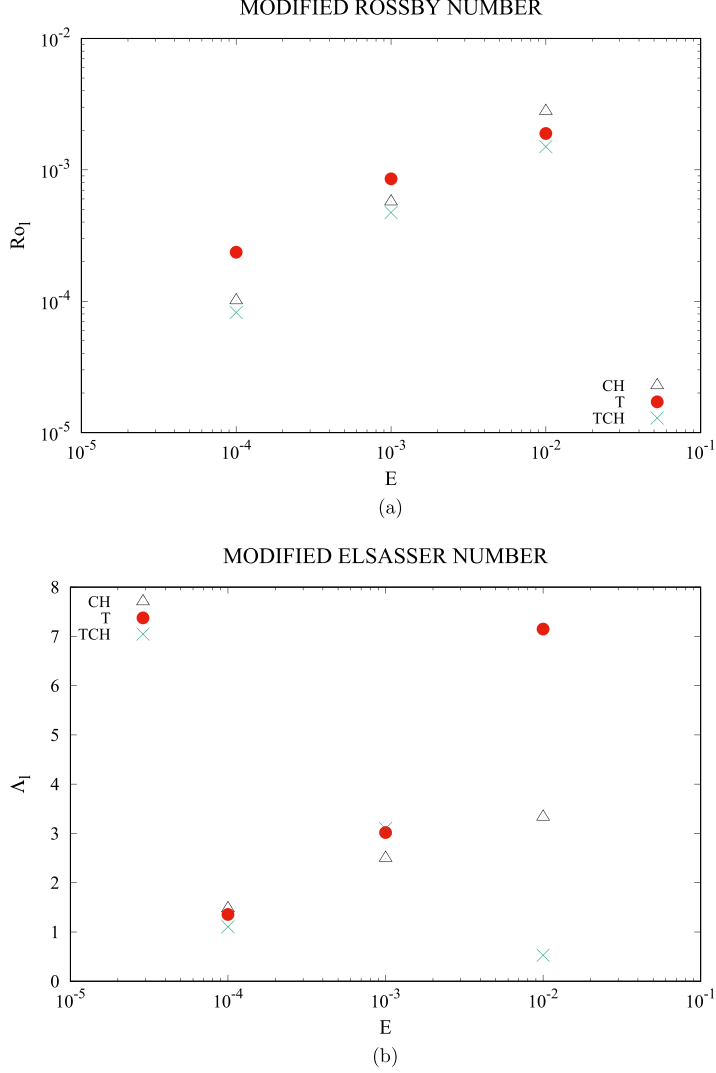


FIG. 7. Time-averaged modified Rossby numbers, Ro_ℓ , (top) and the time-averaged modified Elsasser numbers, Λ_ℓ , (bottom) as a function of the model type and the Ekman number, E . The axes are in log scales except the vertical axis for Λ_ℓ .

of ω_i decrease as the Ekman number decreases. At $E = 10^{-2}$, the ω_i values are unrealistically high due to the relatively large Ekman number used. In contrast, at $E = 10^{-4}$, the values ω_i are much lower. The influence of the model type ($\mathcal{T}$, $\mathcal{CH}$, or $\mathcal{TCH}$) on ω_i is more pronounced at higher Ekman numbers and becomes less significant as E decreases (see Fig. 9).

Table II compares our inner core rotation rates with results from other numerical Geodynamo simulations [6,23,24,41–43]. This comparison suggests that while ω_i decreases with decreasing E , the inclusion of gravitational coupling in Geodynamo models has a more pronounced effect. Note that in our model, Γ_g is incorporated. As shown by Buffett and Glatzmaier [24] and Aubert and Dumberry [41], accounting for Γ_g reduces ω_i more effectively than lowering E alone. Notably, Aubert and Dumberry [41] obtained ω_i values even smaller than those inferred from seismic

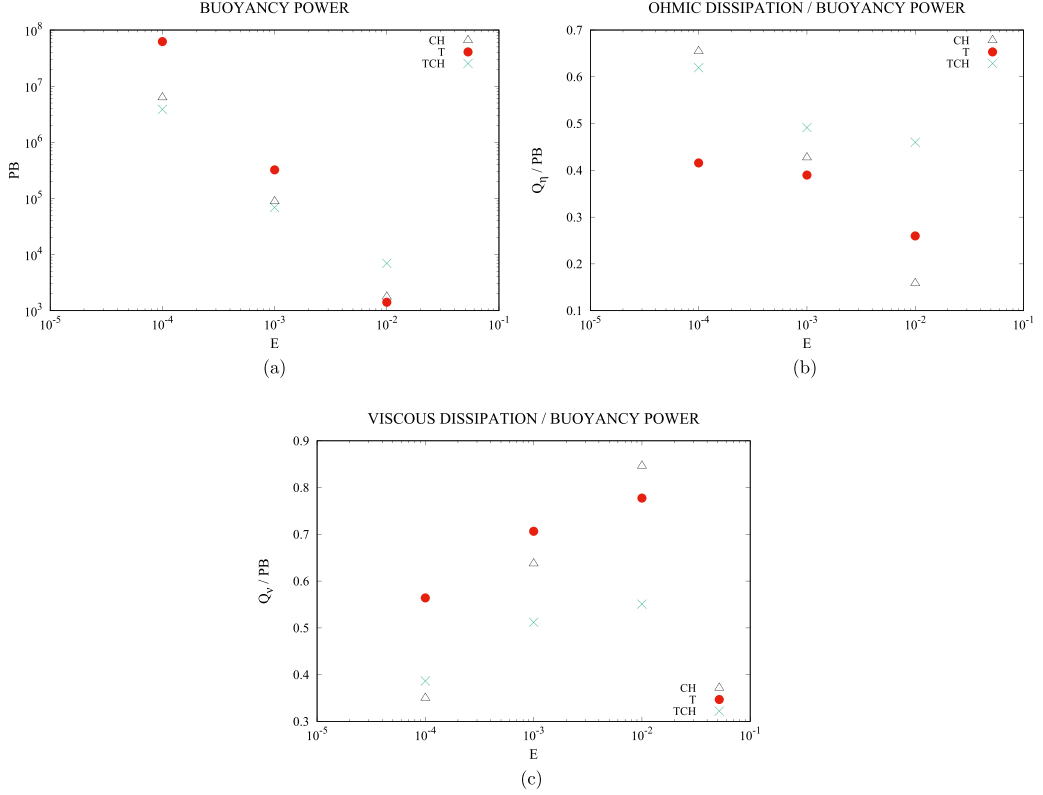


FIG. 8. Power generated by buoyancy forces, PB, which in the statistically steady dynamo state is equal to the total energy dissipation, Q , together with the ratios of the ohmic and viscous dissipations, Q_η and Q_v , respectively, to the total dissipation, as a function of the model type and the Ekman number, E . The plot of PB has both axes in log scales, while the rest of the plots has only the horizontal axes in log scales.

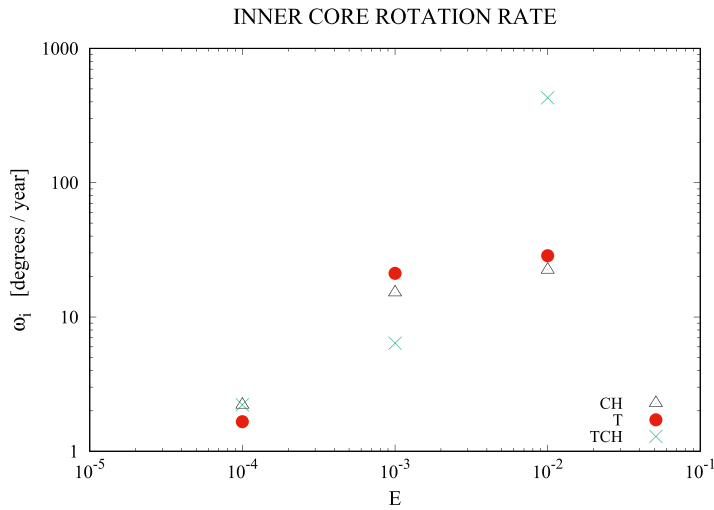


FIG. 9. The time-averaged inner core rotation rate, ω_i , in dimensional units [degrees/year] as a function of the model type and the Ekman number, E . The plot has both axes in log scales.

observations. Moreover, their results suggest that seismic observations cannot represent a steady rotation.

V. CONCLUSIONS

Previous analyses of convection in the Earth's core and the geodynamo have provided insights into the characteristics of the magnetic and velocity fields, temperature, concentration, codensity, and heat flow, depending on various parameters such as the Ekman, Prandtl, magnetic Prandtl, and Rayleigh numbers. These studies have also considered different sources of convection, boundary conditions, and their impacts [3–7,12–14,27,29–33,38,44].

In this study, we examine the characteristics of the generated magnetic and velocity fields, temperature, concentration, codensity, polarity reversals, energy balance, and inner core rotation rate, as influenced by different sources of buoyancy. This comprehensive approach allows us to compare all three types of convection and their respective impacts on the dynamo process across a range of Ekman numbers. All analyses are conducted using the fixed set of parameters listed Table I.

Our results demonstrate that the source of buoyancy significantly affects the generated magnetic field and convective velocity, as shown in Figs. 1–5. These findings support previous studies, particularly those of Kutzner and Christensen [27]. Their work showed that at lower values of R_a and P_m , a stable, dipole-dominant, nonreversing magnetic field is favored. As R_a and P_m increase, the dipole-dominant field persists but exhibits excursions and reversals. With a further increase in R_a and P_m , they observed a transition to a nondipole dynamo in which the dipole component undergoes more frequent excursions and reversals. In this study, we selected higher values of R_a and P_m to access the regime of reversing dynamos, as described by Ref. [27]. The characteristics of the velocity and magnetic fields in our results are also consistent with their findings. Computational constraints—primarily CPU time on the supercomputer—prevented us from extending the simulations long enough to capture longer time series at lower Ekman numbers. As a result, we could not observe the dipole polarity reversals at these lower E values. Under these conditions, the dipole remains strong and stable and does not reverse within the investigated time intervals, consistent with Kutzner and Christensen [27]. Thus, for $E = 10^{-3}$ and 10^{-4} , the selected R_a and P_m values were appropriate, but the simulation durations were insufficient to capture the first reversals.

Our study extends the findings of Kutzner and Christensen [27] by incorporating codensity-driven convection and demonstrating how different convection sources and Ekman numbers influence Ro_ℓ , A_ℓ , PB, Q_η , Q_v and ω_i —parameters not presented in Ref. [27]. We show that the dependence on the source of convection is also reflected in the Elsasser numbers and the inner core rotation rate. However, as the Ekman number decreases, this dependence diminishes and eventually becomes negligible (Figs. 7 and 9). In contrast, the dependence of the modified Rossby numbers, buoyancy power, and the ratios of ohmic and viscous dissipation to buoyancy power remains significant across all Ekman numbers considered (Figs. 7 and 8). For thermally driven dynamos, we did not observe polarity reversals, which is consistent with the previously described sharp transition from reversing to nonreversing dynamos (see Kutzner and Christensen [27]). For higher Ekman numbers, polarity reversals are influenced by the buoyancy source (Fig. 6). In general, we find that the influence of different buoyancy sources is stronger at higher E values and weaker at lower ones.

It is important to highlight that the characteristics of polarity reversals at high Ekman numbers differ significantly from those at lower values. Wicht and Meduri [45] concluded that polarity is amplified during dipole recovery, and that reversals and grand excursions—defined as excursions during which the dipole temporarily assumes reverse polarity—are equally probable. While the overall reversal behavior appears Earth-like, a more detailed comparison with paleomagnetic observations suggests that the simulated reversal events are excessively prolonged, and that grand excursions occur too infrequently. At particularly large Ekman numbers, they identified a second, less Earth-like type of reversal, in which the total magnetic field decays completely and subsequently recovers after a certain period.

Furthermore, our analysis supports the hypothesis proposed by Glatzmaier and Roberts [23], and Roberts and Glatzmaier [6], indicating that lower ω_i values are obtained as the Ekman number decreases. However, because our model incorporates gravitational coupling, the resulting ω_i values are further reduced. A direct comparison with the results of Šimkanin and Kyselica [43] shows comparable ω_i values, despite their use of a much lower Ekman number but without accounting for gravitational coupling (see Table II). When studying the Earth's inner core rotation in numerical Geodynamo models, accounting for gravitational coupling is likely more important than to focus on viscosity (see Refs. [24,41,42,46]).

Aubert and Dumberry [41] suggested that seismic observations cannot reflect a steady inner core rotation. This aligns with the findings of Tkalčić *et al.* [47], who reported quasioscillatory inner core rotation rates with a period of about 20 years between 1961 and 2007. Livermore *et al.* [46] proposed that decadal variations in the geomagnetic field may drive these fluctuations, making them consistent with the observed quasioscillatory behavior. Similarly, Witze [25] and Yang and Song [26] also demonstrated that the inner core rotation rate varies over time. Yang and Song [26], for example, analysed seismic waves from the early 1990s and found that while seismic wave propagation previously showed significant temporal variations, it has exhibited minimal change over the past decade—suggesting that the inner core's super-rotation may have recently paused. By comparing seismic records dating back to 1964, they identified a gradual change in inner core rotation as part of an approximately 70-year oscillation. Additionally, their analysis of seismic waves traveling through the inner core since the 1960s revealed a multidecadal oscillation of the inner core's rotation, which correlates with variations in the length of day and changes in the Earth's magnetic field. This leads to the conclusion that, in addition to periods of super-rotation, the inner core can also experience subrotation and intervals during which it does not rotate at all. All three modes of Earth's inner core rotation will be investigated in future Geodynamo simulations.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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