



Ice Giants Revisited: Uranus and Neptune as Magma Ocean Worlds

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Abstract

Uranus and Neptune are commonly interpreted as volatile-rich “ice giants,” an assumption that underpins most interior models. Here we show that their observed radii, bulk densities, gravitational harmonics, normalized moments of inertia, intrinsic luminosities, and key features of their atmospheric compositions are consistent with interiors comprising supercritical, hydrogen-rich magma oceans overlain by H_2 -rich envelopes. Our results, based on three fit parameters for each planet, provide a parsimonious explanation for the structures, thermal states, and atmospheric chemistries of Uranus and Neptune. We find that the solar system’s ice giants can be understood as magma-ocean giants, with origins parallel to those of sub-Neptune gas-dwarf planets. A continuum among gas-dwarf planets permits Neptune and Uranus to serve as accessible, data-driven test cases for structure models and material properties used to understand sub-Neptunes.

Unified Astronomy Thesaurus concepts: Neptune (1096); Uranus (1751); Exoplanets (498); Extrasolar ice giants (2024); Exoplanet structure (495); Mini Neptunes (1063)

1. Introduction

Uranus and Neptune were originally classified as “ice giants” on cosmochemical grounds. For a solar-composition protoplanetary disk beyond the snow line, the volatile compounds H_2O , CH_4 , and NH_3 , comprising “ices,” were predicted to outweigh refractory rock-forming material by roughly a factor of 2 (R. T. Reynolds & A. L. Summers 1965; M. Podolak & A. G. W. Cameron 1974; M. Podolak & R. T. Reynolds 1984; M. Podolak et al. 1995), motivating the canonical three-layer interior model of a rocky core, an icy mantle, and a hydrogen–helium envelope (W. B. Hubbard & J. J. MacFarlane 1980; M. S. Marley et al. 1995).

The canonical three-layer models for the interiors of Neptune and Uranus usually include a rocky core, an intermediate “icy” mantle of water, methane, and ammonia, and an outer H – He -rich envelope (W. B. Hubbard & J. J. MacFarlane 1980; J. J. Fortney & N. Nettelmann 2010; N. Nettelmann et al. 2013). However, data neither require the interiors of these ice giants to be fully differentiated nor water dominated in order to fit the observables. Empirical structure models that fit the observed gravitational harmonics without assuming a specific composition reveal that continuous density gradients are equally consistent with the data and that rock-dominated interiors cannot be ruled out (R. Helled et al. 2011, 2020; N. A. Teanby et al. 2020; B. A. Neuenschwander & R. Helled 2022; B. A. Neuenschwander et al. 2024).

The combination of a large number of model parameters and a comparatively smaller number of observables has led to considerable degeneracies among models for the structure and chemistry of the interiors of Uranus and Neptune. We do not attempt to describe all successful models for the interiors of these planets, but we note that new variants that invoke larger fractions of rocky material compared with icy progenitors have been proposed recently (e.g., U. Malamud et al. 2024; L. Morf & R. Helled 2025; R. Tejada Arevalo 2025; V. Ramirez et al.

2026), as well as incorporation of exotic superionic phases (C. Liu et al. 2026).

In this paper we present modeling evidence that the ice giants of the solar system may actually be supercritical magma-ocean giants. The proposition is motivated in part by an evolving understanding of the astrophysical constraints on ices in protoplanetary disks and new findings regarding the physical chemistry of gas-dwarf planet interiors, both of which are in tension with the foundations of the canonical ice-giant models.

First, it is not clear that Uranus and Neptune formed from progenitors containing as much ice as is often cited. A budget of planet-forming materials based on solar elemental abundances shows that the commonly assumed ice-to-rock mass ratio of approximately 2:1 requires the condensation and incorporation of CO, and therefore formation beyond the CO snow line. Assigning the carbon in a solar-composition gas to CO, the oxygen remaining after CO formation first to silicates and then to water, and Fe to metal yields a total inventory composed of approximately 54% CO, 23% silicate, 12% water, and 11% Fe by mass. If the full CO inventory were incorporated into planetesimals, the condensed material beyond the CO snow line would therefore be approximately 66% ice by mass, consisting of 54% CO and 12% water. The dominance of CO ice persists even if one assumes that one-half of the carbon resides in organic solids. Between the water and CO snow lines, however, CO remains in the gas phase and does not contribute to the condensed planet-forming material. The remaining solids are then composed of approximately 50% silicate, 27% water, and 23% Fe by mass. Thus, water by itself comprises only roughly one-quarter of the planet-forming material outside the water snow line but inside the CO snow line, excluding H_2 and He, considerably less than the 2:1 water-to-rock ratios assumed in many ice-giant models. This calculation also has the merit of offering a simple explanation for the approximately 2:1 ratio of silicate to Fe metal in Earth. Timing of accretion, incomplete condensation, and the efficiency with which CO ice was incorporated into planetesimals affect the specific proportions, but an ice-dominated progenitor composition requires substantial accretion of CO-rich material from beyond the CO snow line.



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The position of the CO snow line in the solar protoplanetary disk at any given time is poorly known, and observations of disks around low-mass pre-main-sequence stars allow for a wide range of possible locations. In TW Hya, often treated as a solar-nebula analog, N_2H^+ emission places the CO snow line near ~ 30 au, similar to the current position of Neptune but beyond that of Uranus, while rare CO isotopologue observations place the midplane CO snow line closer to ~ 20 au (C. Qi et al. 2013; K. R. Schwarz et al. 2016; K. Zhang et al. 2017), similar to the present-day semimajor axis of Uranus. Other T Tauri disks exhibit a similarly broad range of inferred CO snow line locations. CO isotopologue modeling yields midplane locations of $\sim 12 \pm 5$ to $\sim 30 \pm 5$ au for AS 209, IM Lup, and GM Aur (K. Zhang et al. 2021), whereas N_2H^+ -based analyses suggest more distal CO and N_2 snow line positions. C. Qi et al. (2019) report CO snow line radii of 58, 48, and 75 au for LkCa 15, GM Aur, and DM Tau, respectively, with corresponding N_2 snow line radii of 88, 78, and 145 au. These analogs indicate that, in solar-type protoplanetary disks during the T Tauri phase, the CO and N_2 snow lines may have resided beyond the present-day semimajor axes of Uranus and Neptune. Consequently, it cannot simply be assumed that the progenitors of either planet formed in a region where CO was available as a major condensed component.

The inference that the progenitors of the ice giants could have formed inside of the CO snow line is supported by the properties of the vestiges of planet formation in the outer solar system. Kuiper Belt objects provide, arguably, the best direct evidence remaining of the composition of the outer planet-forming regions of the solar system. These bodies are approximately 70% silicate (C. Bierson & F. Nimmo 2019). The ices comprising Pluto are dominantly water, with CO, CH_4 , and N_2 ices representing only a relatively thin veneer on the surface (S. A. Stern et al. 2015; W. B. McKinnon et al. 2017; B. Schmitt et al. 2017). Jupiter-family comets also sample volatile-rich planetesimals formed in the outer solar system. Their compositions are consistent with those of Kuiper Belt objects in suggesting dust-to-ice ratios of about three, corresponding to ice fractions of $\sim 25\%$ (M. Pätzold et al. 2016; M. Fulle et al. 2019).

The solar-abundance budget, astronomical constraints on snow line locations, and the compositions of surviving outer solar system bodies therefore provide no compelling evidence that the progenitors of Uranus and Neptune must have been dominated by ice. Although condensation of CO could produce an ice-dominated inventory at sufficiently large heliocentric distances, observational evidence for bodies dominated by ices other than water remains limited. For these reasons, it is not at all clear that the progenitors of the interiors of planets formed in the outer solar system contained the large ice fractions commonly assumed in ice-giant models.

Second, chemistry at the conditions appropriate for the interiors of planets is equally important, if not more important, in determining the nature of their interiors than their radial position of accretion. This point is underscored by recent models for sub-Neptunes (H. E. Schlichting & E. D. Young 2022; E. D. Young et al. 2024, 2025). Copious amounts of water can be produced by reactions between primary hydrogen atmospheres and hot cores beneath (E. Kite & L. Schaefer 2021; H. E. Schlichting & E. D. Young 2022; E. D. Young et al. 2023; F. Miozzi et al. 2025; A. Werlen et al. 2025a, 2026). Therefore, the amount of H_2O in a planet is not directly related to where it

accreted. Indeed, while more water may accrete to growing bodies in the outer protoplanetary disks, subsequent chemical evolution can lead to similar water concentrations for planets formed interior to the water ice line (A. Werlen et al. 2025a). What is more, the chemical memory of accreted water ices is, to first order, most likely erased in the interiors of bodies due to miscibility with silicates. To see this, consider a planet assembled from planetesimals that were 25% by mass water ice, 50% rock (MgSiO_3), and 25% Fe (Earth-like rock and metal). Density functional theory (DFT) molecular dynamics (MD) calculations by T. Kovavevic et al. (2022) indicate that molten MgSiO_3 and H_2O can become fully miscible at high pressure and temperature, similar to the behavior in the $\text{MgSiO}_3\text{--H}_2$ system. The initial 25% H_2O by mass contributes 22.2% oxygen and 2.8% hydrogen to the bulk planet. The oxygen will comprise part of the dense oxide-like component in the supercritical miscible phase, whereas the hydrogen increases the low-density component. Water-derived oxygen may be regarded as providing a means to oxidize Fe to FeO, increasing the rock-to-metal ratio, but this is an intraphase redox redistribution and is not readily discernible in the fully miscible melt. The first-order effect of accreting water to an $\text{MgSiO}_3 + \text{Fe} + \text{H}_2$ mixture is therefore an increase in the effective hydrogen inventory, while the oxygen augments the dense interior. This water-derived H_2 is in addition to the free molecular H_2 accreted by the growing planet. Our models should largely capture this density effects in the form of total hydrogen.

Finally, the suggestion that the ice giants may be more rocky than previously thought (N. A. Teanby et al. 2020) requires an analysis of the physical chemistry of “rock” when intermingled with a large reservoir of H_2 at temperatures of thousands of degrees and pressures of tens to hundreds of GPa. Recent work indicates that silicate, hydrogen, and iron are miscible at the conditions that prevailed in the interiors of ice giants, necessitating a reexamination of what is meant by “rocky” in this context (E. D. Young et al. 2024, 2025; T. Gilmore & L. Stixrude 2026).

This paper is organized as follows. In Section 2 we describe the observational constraints on the compositions of the ice giants. In Section 3 we describe the supercritical magma-ocean model, and in Section 4 we describe our methods for deriving fits to the planet observables using this model. Results are described in Section 5, and various implications, predictions, and limitations are described in Section 6. We provide our conclusions in Section 7.

2. Ice-giant Observational Constraints

The primary observational constraints used in interior structure models for these planets are the 1 bar equatorial radius, bulk density, and the gravitational harmonics J_2 and J_4 . The 1 bar radius and the first determinations of J_2 and J_4 were established by the Voyager 2 flyby (G. L. Tyler et al. 1989; G. F. Lindal 1992; Tables 1 and 2). No spacecraft has revisited either planet since. The J_2 and J_4 values used in modern interior models are not the raw Voyager 2 numbers, however. Rather, they have been substantially refined through long-baseline astrometry of satellite orbital motions (R. A. Jacobson 2009; M. Brozović et al. 2020; B. Wang et al. 2023). In addition, the gravitational parameters are composed of a static component due solely to oblateness and a dynamic component arising from winds (Y. Kaspi et al. 2013).

Table 1
Observational Constraints Used in the Neptune Interior MCMC

Quantity	Value	Uncertainty	Source
Mass (M , kg)	1.02413×10^{26}	fixed	G. L. Tyler et al. (1989)
1 bar equatorial radius (R_{eq} , km)	24,766	± 15	G. F. Lindal (1992)
Rotation period (P_{rot} , s)	58,000	fixed ^a	G. L. Tyler et al. (1989)
Equilibrium temperature (T_{eq} , K)	46.6	± 1.1	J. C. Pearl & B. J. Conrath (1991)
Thermal constraints			
1 bar temperature ($T_{1 \text{ bar}}$, K)	72	± 2	G. F. Lindal (1992)
Intrinsic luminosity (L_{int} , W)	3.3×10^{15}	$\pm 0.35 \times 10^{15}$	J. C. Pearl & B. J. Conrath (1991)
Gravitational harmonics at $R_{\text{ref}} = 25,225$ km			
J_2	3401.655×10^{-6}	$\pm 3.994 \times 10^{-6}$	B. Wang et al. (2023)
J_4	-33.294×10^{-6}	$\pm 10.000 \times 10^{-6}$	B. Wang et al. (2023)
Renormalized to $R_{\text{eq}} = 24,766$ km (this work)			
J_2	3528.912×10^{-6}	$\pm 4.143 \times 10^{-6}$	B. Wang et al. (2023), renorm.
J_4	-35.832×10^{-6}	$\pm 10.762 \times 10^{-6}$	B. Wang et al. (2023), renorm.
Model consistency target			
Norm. moment of inertia (C/MR^2)	0.2410	± 0.0048 (2%)	N. Nettelmann et al. (2013)

Notes. Gravitational harmonics from B. Wang et al. (2023) are renormalized from $R_{\text{ref}} = 25,225$ km to the 1 bar equatorial radius $R_{\text{eq}} = 24,766$ km via $J_n(R_{\text{new}}) = J_n(R_{\text{old}})(R_{\text{old}}/R_{\text{new}})^n$, increasing J_2 by $\sim 3.7\%$. The C/MR^2 uncertainty reflects the assumed model-to-observation systematic.

^a Voyager 2 magnetic field periodicity; known to be inconsistent with Neptune’s observed shape (N. Nettelmann et al. 2013). See text.

Published harmonics for Neptune are tabulated (B. Wang et al. 2023) at a reference radius $R_{\text{ref}} = 25,225$ km (B. Wang et al. 2023), which exceeds the physical 1 bar equatorial radius $R_{1 \text{ bar}}$ of 24,766 km (R. G. French et al. 2024) adopted as the reference radius for our interior models. We therefore renormalize the J_{2n} values to $R_{1 \text{ bar}}$:

$$J_{2n}(R_{\text{new}}) = J_{2n}(R_{\text{ref}}) \left(\frac{R_{\text{ref}}}{R_{\text{new}}} \right)^{2n}, \quad (1)$$

yielding $J_2 = 3528.91 \times 10^{-6}$ and $J_4 = -35.83 \times 10^{-6}$ at the 1 bar level, derived from the B. Wang et al. (2023) values.

Neptune’s rotation period is fixed to the Voyager 2 magnetic field periodicity of 58,000 s (G. L. Tyler et al. 1989), which is known to be inconsistent with the planet’s observed oblateness (R. Helled et al. 2010; N. Nettelmann et al. 2013). This introduces a model systematic distinct from the purely astrometric uncertainties in Table 1 (B. Wang et al. 2023). The normalized moment of inertia C/MR^2 has not been directly measured for either planet, but model values exist (N. Nettelmann et al. 2013; E. Bailey & D. J. Stevenson 2021; B. A. Neuenschwander & R. Helled 2022) that can serve as useful arbiters for structure models.

The intrinsic luminosities of Uranus and Neptune provide additional constraints on their interiors. Intrinsic luminosities are inferred from the planetary energy balance, requiring both the effective temperature (from thermal emission) and the Bond albedo (from reflected sunlight), the latter being the primary source of uncertainty. The intrinsic luminosities of Uranus and Neptune were first determined from Voyager 2 IRIS data (J. C. Pearl et al. 1990; J. C. Pearl & B. J. Conrath 1991). These data indicate energy balances of $\mathcal{E} = 1.06 \pm 0.08$ and 2.61 ± 0.28 respectively, where $\mathcal{E}$ is the ratio of emitted to absorbed power related to the intrinsic luminosity by $L_{\text{int}} = (\mathcal{E} - 1)(1 - A)\pi R^2 \mathcal{F}_{\odot}$. Here A is the Bond albedo, R is the planet radius, and $\mathcal{F}_{\odot}$ is the incident solar flux at the planet’s orbital distance. The measured values yield intrinsic luminosities of $L_{\text{int}} \lesssim 3.5 \times 10^{14}$ and 3.3×10^{15} W for Uranus and Neptune, respectively. Neptune’s substantial

internal heat source is well established, with albedo-driven uncertainties of order 20%, while Uranus’ near-zero intrinsic luminosity has been revised upward recently to $L_{\text{int}} \approx 6.3 \times 10^{14}$ W (P. G. J. Irwin et al. 2025; W. Wang et al. 2025), driven primarily by a reanalysis of the Bond albedo. Compositional stratification and limitations on convection have been invoked to match the observed thermal states (A. Vazan & R. Helled 2020; B. A. Neuenschwander et al. 2024). Sources of stratification proposed previously include phase separation between partially miscible H_2 and H_2O (E. Bailey & D. J. Stevenson 2021; M. Cano Amoros et al. 2024) and $\text{H}_2\text{O}-\text{CH}_4-\text{NH}_3$ exsolution to distinct water and C–H–N-rich layers (B. Militzer 2024).

The 1 bar temperatures, $T_{1 \text{ bar}}$, provide additional constraints for interior models for both ice giants. Because the equilibrium temperatures are low, the values for $T_{1 \text{ bar}}$ are sensitive to the deep structure that sets L_{int} as well as to the details of the opacities.

The atmospheres of both planets are dominated by H_2 and He, with CH_4 being the third-most abundant constituent. In Neptune, the deep tropospheric CH_4 mole fraction, retrieved below the ~ 1.5 bar condensation level and taken to represent the well-mixed bulk abundance, varies from 6% to 7% at equatorial latitudes to $\sim 3\%$ near the poles (E. Karkoschka & M. G. Tomasko 2011; P. G. J. Irwin et al. 2019a, 2021, 2023). Uranus exhibits roughly one-half this abundance, with 1% to 4% at low latitudes and $\sim 1\%$ near the poles below its ~ 1.2 bar condensation level (L. A. Sromovsky et al. 2014, 2019). In both cases the subcondensation mole fraction is the closest available proxy for the bulk atmospheric carbon abundance, since deeper levels are inaccessible to remote sensing. These concentrations of methane imply carbon enrichments of $\sim 80 \times$ solar for Neptune’s atmosphere and ~ 40 to $50 \times$ solar for Uranus’ atmosphere (T. Cavalie et al. 2020). H_2S has also been detected near 1 to 4 bar in both planets with mole fractions of 0.4 to 0.8 ppm for Uranus and ~ 1 to 3 ppm for Neptune. NH_3 is undetected at any level in both atmospheres (I. de Pater et al. 1991), although its presence has been inferred on the basis of latitudinal variations in microwave opacities

Table 2
Observational Constraints Used in the Uranus Interior MCMC

Quantity	Value	Uncertainty	Source
Mass (M , kg)	8.68099×10^{25}	fixed	R. A. Jacobson & R. S. Park (2025)
1 bar equatorial radius (R_{eq} , km)	25,559	± 4	G. F. Lindal et al. (1987)
Rotation period (P_{rot} , s)	62,064	fixed	M. D. Desch et al. (1986)
Equilibrium temperature (T_{eq} , K)	58.1	± 1.0	J. C. Pearl et al. (1990)
Thermal constraints			
1 bar temperature ($T_{1 \text{ bar}}$, K)	76	± 2	G. F. Lindal et al. (1987)
Intrinsic luminosity (L_{int} , W)	6.3×10^{14}	$\pm 1.5 \times 10^{14}$	W. Wang et al. (2025)
Gravitational harmonics at $R_{\text{ref}} = R_{\text{eq}} = 25,559$ km			
J_2	3509.291×10^{-6}	$\pm 0.412 \times 10^{-6}$	R. G. French et al. (2024)
J_4	-35.522×10^{-6}	$\pm 0.466 \times 10^{-6}$	R. G. French et al. (2024)
Model consistency target			
Norm. moment of inertia (C/MR^2)	0.2250	± 0.0045 (2%)	N. Nettelmann et al. (2013), M. Podolak & R. Helled (2012)

Note. Gravitational harmonics from R. G. French et al. (2024) are reported at $R_{\text{ref}} = R_{\text{eq}} = 25,559$ km. The C/MR^2 uncertainty reflects the assumed model-to-observation systematic.

(M. D. Hofstadter & B. J. Butler 2003). The detection of gaseous H_2S and the nondetection of NH_3 implies a sulfur to nitrogen ratio substantially greater than solar in the accessible regions of the atmospheres for both planets (I. de Pater et al. 1991; P. G. J. Irwin et al. 2018, 2019b). Stratospheric photochemistry driven by CH_4 photolysis produces C_2H_6 and C_2H_2 , detected in both atmospheres, together with trace CO and HCN (J. I. Moses et al. 2020).

3. Model

A growing understanding of the physical chemistry of the silicate–hydrogen–iron system is elucidating the nature of gas-dwarf sub-Neptune planets. We use the term gas dwarf here to refer to planets thought to consist of magma oceans¹ overlain by H_2 -rich envelopes but small enough to have avoided runaway growth. We have learned that at the temperatures and pressures imposed by accretion and self-gravity, the interiors of these planets are likely to be a single, supercritical mixture of silicate, iron, and hydrogen, modeled as miscible MgSiO_3 , Fe, and H_2 based on ab initio DFT-MD simulations (E. D. Young et al. 2024, 2025; J. G. Rogers et al. 2025; T. Gilmore & L. Stixrude 2026). These calculations are supported by recent experimental findings (F. Miozzi et al. 2025). The boundary between this supercritical magma ocean and an overlying H_2 -rich envelope is a first-order phase transition with coexisting compositions in the MgSiO_3 – H_2 system defined by a pressure-dependent binodal in temperature (T) versus composition (x_{H_2}) space. The same DFT-MD simulations that inform us about the phase equilibria also provide constraints on the equation of state (EoS) for the supercritical magma oceans.

The picture of these planets that emerges from the new findings is one composed of four main layers (Figure 1). The deepest layer is a magma ocean composed of a supercritical mixture of silicate, iron, and hydrogen, similar to mixtures of MgSiO_3 , Fe, and H_2 . For weight percent fractions of total hydrogen, the density of the overlying hydrogen-rich envelope is sufficient to preserve a molten magma ocean for several Gyr timescales (S. Ginzburg et al. 2016). At large radial distances from the center, at the conditions specified by the binodal for

the first-order phase transition, the magma ocean gives way to an H_2 -rich envelope phase. At the base of the envelope is a boundary layer in which convection is inhibited (G. Ahlers et al. 2009), augmented by molecular-weight gradients predicted by the phase equilibria (S. Markham et al. 2022; W. Misener et al. 2023; E. D. Young et al. 2024). Estimates for the Rayleigh number in the supercritical melt are sufficiently high that the supercritical magma ocean should be actively convecting (E. D. Young et al. 2024). Above the basal boundary layer, the envelope is convective up to altitudes where radiative diffusion dominates over convective heat transport.

Details of the planet model, including material properties, EoSs, phase equilibria, and governing equations are given in Appendices A through D. Summarizing the salient features briefly here, the supercritical magma ocean has a density much lower than that of pure MgSiO_3 . Differences in the pressure dependence of ρ from pure MgSiO_3 is constrained from DFT-MD simulations at several hydrogen concentrations (see Appendix D). Hydrogen alters the compression behavior of the liquid. While the reference bulk modulus, K_T , is comparable in the EoSs of pure MgSiO_3 and the supercritical mixture of MgSiO_3 and 4% H_2 by mass, the pressure derivative of the bulk modulus decreases from 5.71 for dry MgSiO_3 liquid to 4.32 for the H-bearing mixture (S. P. Marcum et al. 2026). This lower K' implies that the supercritical mixture stiffens less rapidly with increasing pressure. Thus, hydrogen results in a lower density and makes the phase more compressible at high pressure. We use a linear fit of K' with hydrogen concentration to extrapolate the EoS to higher concentrations of H_2 (Appendix D).

The fluid comprising the magma-ocean phase is expected to have a melting temperature of order 1000°C lower than that of pure MgSiO_3 due to the melting-point depression effect of mixing (Appendix D). Iron that would otherwise form an iron metal core is entirely miscible in the MgSiO_3 – H_2 melt, based on phase equilibria in the ternary MgSiO_3 –Fe– H_2 system (E. D. Young et al. 2025), and has a limited effect on the density (Appendix D). The physical chemistry thus ensures that the interior of the planets is fully molten and composed of a single convecting phase. Immediately above the magma ocean, the binodal specifies that there is a heavy load of Mg, Si, and O present in the hydrogen-dominated phase. We consider that the silicate species will condense moving outward as temperature decreases. We account for condensation as a Rayleigh-like distillation as silicates and oxides rain

¹ Because the supercritical mix is so hydrogen rich, the use of the term “magma” may no longer be entirely appropriate. We use it to distinguish it from the even more H_2 -rich envelope.

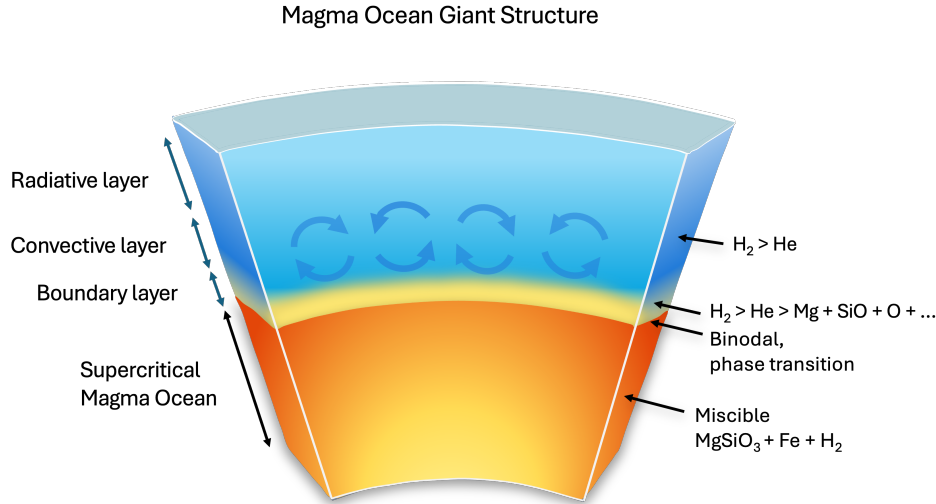


Figure 1. Schematic showing the layered structure of Neptune and Uranus as proposed here (not to scale).

out back to the surface. This leads to a significant molecular-weight gradient that stabilizes the base of the envelope against convection. The Rayleigh number relative to the critical Rayleigh number, Ra/Ra_{crit} , is $\geq 10^9$, with the determination of precise values hampered by uncertainties associated with various scaling laws. Nonetheless, once the molecular-weight gradient subsides, the envelope should be actively convecting. We employ a moist adiabat to account for the enthalpy effect of condensation. The convective layer transitions smoothly into a radiative-diffusion regime above, and eventually to a radiative layer, which we model using a two-stream solution (Appendix A).

4. Methods

We have applied this picture of gas-dwarf planets to Uranus and Neptune. We modeled the interior structures of Uranus and Neptune with `Planet_LAB`,² a forward code that constructs a 1D spherical planet realization from three parameters. The parameters are the pressure of the binodal separating the magma ocean from the overlying envelope, P_{binodal} , for the planet, the total H_2 mass fraction for the planet, x_{H_2} , and the ratio of the Rayleigh number for the atmosphere to the critical Rayleigh number that specifies the width of the boundary layer at the base of the atmosphere, Ra/Ra_{crit} (Appendices A–D). For each proposed realization, `Planet_LAB` integrates hydrostatic equilibrium on spherical shells outward, crossing through a binodal-defined magma ocean–envelope boundary. We use trapezoidal quadrature over 1.2×10^5 radial steps, returning the density profile $\rho(r)$ and the radius $R_{1\text{ bar}}$ at which the column reaches 1 bar. Because the integration is spherical, $R_{1\text{ bar}}$ is a volumetric mean radius. The mass and (spherical) moment of inertia are

$$M = \int_0^{R_{1\text{ bar}}} 4\pi r^2 \rho(r) dr, \quad (2)$$

$$I = \int_0^{R_{1\text{ bar}}} \frac{8\pi}{3} \rho(r) r^4 dr, \quad (3)$$

with the normalized moment of inertia $C/MR_{1\text{ bar}}^2 = I/(MR_{1\text{ bar}}^2)$.

The gravitational moments and rotational figures are computed a posteriori from the spherical density profiles using the concentric Maclaurin spheroid (CMS) method (W. B. Hubbard 2013), implemented in the Python translation of the original Matlab code (N. Movshovitz et al. 2019). We pass $\rho(r)$, $R_{1\text{ bar}}$, and the published rotation period to the CMS module, which iterates the equipotential surfaces of the rotating fluid figure and returns the dimensionless multipole coefficients J_2 and J_4 . Fully self-consistent oblate methods—the Theory of Figures (N. Nettelmann et al. 2013; R. Helled & J. J. Fortney 2020) and the natively coupled CMS treatments of S. M. Wahl et al. (2017) and N. Movshovitz et al. (2019)—iterate the density profile and the equilibrium shape simultaneously on level surfaces, allowing the density distribution to feed back into the rotating figure as it relaxes. The CMS as implemented here solves for the shape and gravity moments self-consistently at the published rotation rate but treats the input $\rho(r)$ as fixed. Allowing the density profile to relax in response to rotational deformation would modify J_n at the $\mathcal{O}(q_{\text{rot}}^2)$ level where $q_{\text{rot}} = \Omega^2 R^3 / GM \approx 0.026$ for both planets (G. Schubert et al. 2004); the shifts are less than the observational uncertainties and certainly less than systematic uncertainties related to the dynamic contributions. The use of a 1D structure code is motivated by the miscibility–immiscibility phase boundary. The radial position of the H_2 –silicate binodal is determined by local (P, T) conditions, and locating it self-consistently across a Markov Chain Monte Carlo (MCMC) ensemble requires evaluating that boundary at every realization. A 1D spherical model permits mapping the boundary as a single radius at each step. Because exsolution creates compositionally stratified layers in which the isothermal level-surface assumption fails, a self-consistent oblate method would require coupling the figure iteration to a 2D heat-transport treatment, which is computationally expensive and beyond the present scope.

In the CMS calculation, lengths are normalized internally by an outer spheroid radius a_0 . Because CMS as implemented here computes the gravity field of a prescribed density distribution rather than iterating the shape and structure

² On GitHub at https://github.com/eyoung UCLA/planet_Lab; with a copy of the most recent version deposited to Zenodo at doi: [10.5281/zenodo.22240097](https://doi.org/10.5281/zenodo.22240097).

jointly, the choice of a_0 represents the reference radius for the J_n values. We adopt $a_0 = R_{1 \text{ bar}}$ so that CMS produces J_n at the volumetric reference radius determined by our spherical structure integration; the J_n are then renormalized to the equatorial reference (below) permitting comparisons to published values. For this purpose we apply a first-order oblate-spheroid bulge correction, where the flattening f is

$$f \approx \frac{3J_2 + q_{\text{rot}}}{2}, \quad (4)$$

and apply it using

$$R_{\text{eq,model}} = R_{1 \text{ bar}} (1 + f/3), \quad (5)$$

such that $J_n \rightarrow J_n (R_{1 \text{ bar}}/R_{\text{eq,model}})^n$. The first-order truncation in Equation (5) is accurate to $\mathcal{O}(f^2) \approx 3 \times 10^{-4}$, well below all reported uncertainties.

From the planet model realizations, including the shape models, we obtained solutions that minimize the residuals between the model and six observable or derived quantities for Neptune and Uranus: the gravitational harmonics J_2 and J_4 (compared to wind-corrected values described below), the normalized moment of inertia C/MR^2 , the 1 bar equatorial radius R_{eq} , the 1 bar temperature $T_{1 \text{ bar}}$, and the intrinsic luminosity L_{int} .

We modify our target J_2 and J_4 values for our static models by removing the dynamic component. For this purpose we perform a manual iteration between our planet realizations and models for the winds (Appendix E), the latter following methods described by W. Dietrich et al. (2021) and adapted to the ice giants as in D. Soyuer et al. (2023).

Most previous works linking Uranus' and Neptune's zonal winds to dynamic gravity contributions (e.g., Y. Kaspi et al. 2013) assume that $J_2^{\text{obs}} = J_2^{\text{stat}}$, so that only J_4 constrains the wind-induced component of the gravity signature. By assuming $J_2^{\text{dyn}} = 0$, deep-reaching winds are excluded a priori. However, deeper winds necessarily generate a nonnegligible J_2^{dyn} . As the problem is degenerate with only two measured gravity moments, we assume a decay function of the winds with depth that is physically motivated, and allow them to extend to the bottom of the convective atmosphere. Details of this calculation and the J_2^{dyn} and J_4^{dyn} values associated with the final models for Uranus and Neptune are presented in Appendix E.

In the results presented here, we find that for Uranus, with winds extending down to the base of the convective layer, at a depth of 4584 km, $\Delta J_2 = J_2^{\text{obs}} - J_2^{\text{stat}} = -13.47 \times 10^{-6}$ and $\Delta J_4 = -3.46 \times 10^{-6}$. For Neptune, winds extending to the base of the convective layer at a depth of 3838 km yield $\Delta J_2 = 0.87 \times 10^{-6}$ and $\Delta J_4 = -3.88 \times 10^{-6}$ (Appendix E). The dynamic contribution to J_2 is negative for Uranus but positive for Neptune, while both dynamic contributions to J_4 are negative. The correction to J_2 for Uranus is ~ 30 times the quoted measurement uncertainty (Table 2) while for Neptune it is 0.2 times the quoted uncertainty (Table 1). The corrections for J_4 for Uranus and Neptune are 7 and 0.4 times nominal uncertainties, respectively. The smaller correction and uncertainty ratios for Neptune reflect the Monte Carlo error analysis for observed values by B. Wang et al. (2023).

The low equilibrium temperatures for the ice giants results in the 1 bar temperature being relatively insensitive to the shortwave irradiation term. In the gray atmosphere calculation, variations in $\gamma = \kappa_{\text{vis}}/\kappa_{\text{IR}}$ have little effect on $T_{1 \text{ bar}}$, and the

Table 3
Best-fit Solution from the Neptune Interior MCMC

Quantity	Predicted	Target	χ^2
P_{binodal} (GPa)	5.42	...	...
Bulk H_2 (wt%)	12.98	...	...
Ra/Ra_{crit}	1.224×10^{12}	...	...
$J_2 (\times 10^{-6}, \text{static})$	3529.50	3528.05 ± 4.140	0.122
$J_4 (\times 10^{-6}, \text{static})$	-31.500	-31.952 ± 3.500	0.017
C/MR^2	0.2381	0.2410 ± 0.00482	0.362
R_{eq} (km)	24760.9	24766.0 ± 22.0	0.054
$T_{1 \text{ bar}}$ (K)	84.4	72.0 ± 5.0	6.150
L_{int} (W)	2.93×10^{15}	$3.30 \pm 0.35 \times 10^{15}$	1.118
χ_{total}^2			7.823
$\chi_{\text{red}}^2 (\nu = 3)$			2.608

Note. The reduced χ^2 is calculated with $n_{\text{obs}} = 6$ constraints and $n_{\text{free}} = 3$ free parameters. A one-sided prior on the melt-layer H_2 content (ceiling of 12.0 wt%; $\sigma = 1.0$ wt%) is included but inactive at the best fit (predicted = 9.27 wt%).

1 bar level is more closely coupled to the intrinsic thermal structure imposed by the intrinsic luminosity, the pressure–optical-depth relation, and the radiative–convective transition. We therefore include $T_{1 \text{ bar}}$ as a constraint on the coupled interior–atmosphere model, while assigning it an estimated systematic model uncertainty of $\pm 5^\circ$, larger than the formal observational error but commensurate with probable uncertainties in opacities (see Section 6). Previous models have generally either anchored an assumed adiabat at the observed $T_{1 \text{ bar}}$ (e.g., N. Nettelmann et al. 2013; E. Bailey & D. J. Stevenson 2021), adjusted the profile by a scale factor to match observations (e.g., N. Nettelmann et al. 2016), or omitted this constraint entirely where the interior is the sole focus. These options are not justified where the lower boundary conditions for the atmospheres are constrained, in this case by the binodal phase boundary.

The best fit is obtained using MCMC searches applied to the forward planet model code that explores the 3D parameter space to find the posterior distribution over P_{binodal} , x_{H_2} , and Ra/Ra_{crit} . These three parameters are used to obtain the maximum likelihood fit for the six observables.

5. Results

5.1. Neptune

Our best-fit model parameters for Neptune are summarized in Table 3. The best-fit model expressed in density versus radius and various features of the model are shown in Figure 2. The fixed input parameters are the total mass of $17.150 M_{\oplus}$ and the equilibrium temperature of 47 K. By adjusting the binodal pressure (the pressure at the phase transition between the magma ocean and the overlying envelope), the total hydrogen fraction of the planet by weight, and Ra/Ra_{crit} we fit the equatorial 1 bar radius, external gravitational harmonics coefficients, adjusted for dynamic effects, inferred normalized moment of inertia, 1 bar temperature, and intrinsic luminosity with a reduced χ^2 of 2.6, based on three degrees of freedom and the uncertainties shown in (Table 1). This corresponds to an upper-tail p value of $\simeq 0.05$. The reduced χ^2 value > 1 is driven entirely by a 12 K overestimate of the 1 bar temperature (Table 3).

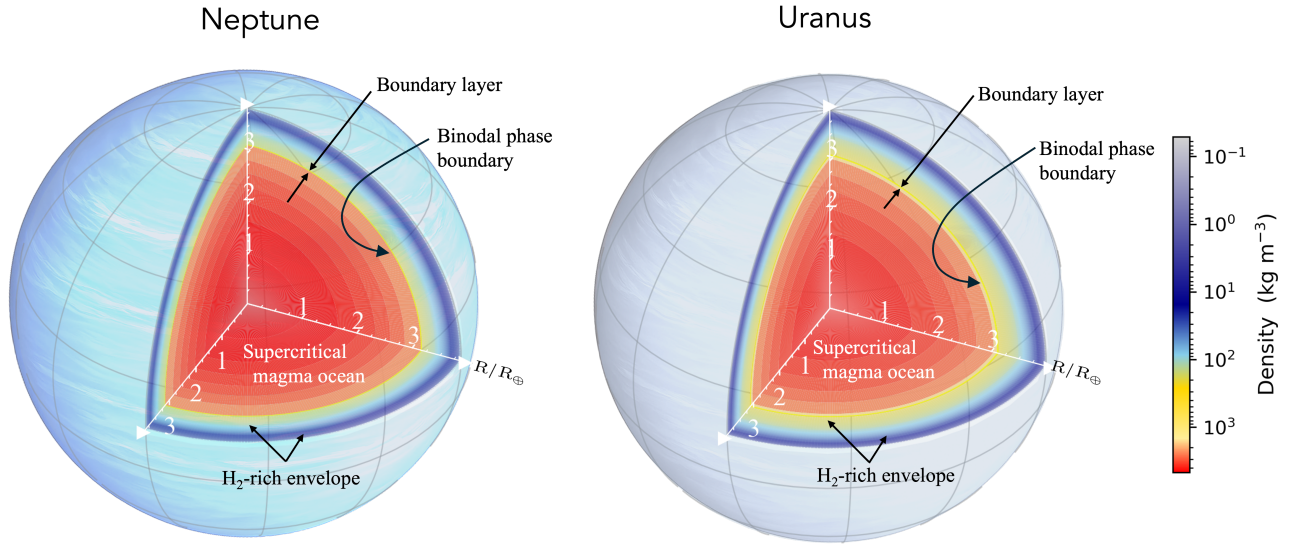


Figure 2. Models for Neptune and Uranus derived in this study expressed as density vs. radius. Contour colors show density. Key elements of the models are labeled. The bright narrow band at the base of the envelope in each case is the boundary region stable against convection, ~ 120 and ~ 180 km wide in the cases of Neptune and Uranus, respectively.

The uncertainties used to calculate the reduced χ^2 comprise a mixture of observational and model uncertainties, whereby purely observational values for σ are used for three of the variables, and systematic errors associated with modeling are used for the other three. The uncertainties on J_2 and L_{int} are purely observational, taken directly from the published values. We could also use $|\Delta J_2|$, the absolute value of our dynamic correction to J_2 , as a potential systematic uncertainty, but for Neptune the wind correction to J_2 is smaller than the observational uncertainty itself ($|\Delta J_2|/\sigma_{J_2, \text{obs}} \approx 0.2$), so no broadening was necessary. The uncertainty on R_{eq} is also observational, where the adopted σ_R of 22 km is intermediate between the reported uncertainties on the equatorial and polar 1 bar radii of ± 15 and ± 30 km respectively, both dominated by atmospheric rotation-period scatter (G. F. Lindal 1992). The uncertainties on J_4 , $T_{1 \text{ bar}}$, and C/MR^2 include known systematic model errors that are greater than the purely observational values. For J_4 , we adopt $\sigma_{J_4} = |\Delta J_4|$, approximately one-half the Monte Carlo uncertainty of B. Wang et al. (2023; Table 1), where ΔJ_4 is our calculated wind correction to the static value. For $T_{1 \text{ bar}}$, we adopt ± 5 K rather than Lindal’s observational ± 2 K, as a concession to systematic uncertainties in atmospheric opacities κ that propagate into the modeled temperature at the 1 bar level (L. N. Fletcher et al. 2018; L. Siebenaler & Y. Miguel 2026). For C/MR^2 , we adopt a 2% model uncertainty (N. Movshovitz & J. J. Fortney 2022; L. Siebenaler & Y. Miguel 2026). The best-fit solution fits all parameters to within the nominal observational uncertainties, excluding $T_{1 \text{ bar}}$, even where these limits were relaxed to account for model systematic errors in the search (compare Tables 3 and 1).

We also ran searches that matched the static (uncorrected) J_2 and J_4 values, using the larger Monte Carlo-based J_4 uncertainty of B. Wang et al. (2023, not reported here). In all cases, the best-fit model J_4 approached the wind-corrected value, about 4×10^{-6} greater (smaller absolute value) than the static observed value, in agreement with our independently calculated wind effect.

The maximum likelihood solution from the MCMC search is obtained for $P_{\text{binodal}} = 5.42$ GPa and a bulk planet H_2 fraction of 12.98% by mass. The overall best-fit structure is composed of a supercritical magma ocean extending out to $3.15 R_{\oplus}$ and an H_2 -rich envelope extending out to a 1 bar equatorial radius of $3.863 R_{\oplus}$ (Figure 2). These best-fit binodal pressure and hydrogen fraction values are strongly positively correlated, as shown in the search corner plot in Figure 3. We note that reasonable changes to the estimates for the EoS for the supercritical magma ocean (e.g., behavior of compressibility with hydrogen content) can change the precise values for these parameters without meaningfully degrading the quality of the fit; the magma-ocean adiabat derived by the MCMC search is effectively the same, but with modestly different values for x_{H_2} and P_{binodal} .

Interestingly, J_4 , a parameter most sensitive to the outer structure of the planet, exhibits a strong negative correlation with the H_2 concentration in the supercritical melt phase (Figure 3). This is an illustration of the integrated nature of the structure model. It can be traced to the links between P_{binodal} and T at the base of the hydrogen-rich envelope, and the concentration of H_2 in melt imposed by these conditions at the binodal boundary.

The various model parameters that characterize the model Neptune are listed in Table 4. Salient features of the model include the binodal temperature of 3256 K, corresponding to the boundary between the supercritical magma ocean and the H_2 -rich envelope, an envelope mass fraction for the planet of 0.04 (note in our model we used a molecular weight of 2.3 g mol^{-1} to accommodate a nominal He fraction), and a Ledoux-stable boundary layer that is 120 km thick above the magma ocean. The model intrinsic luminosity of $2.9 \times 10^{15} \text{ W}$ is virtually indistinguishable from the measured value for Neptune of $3.3 \pm 0.7 \times 10^{15} \text{ W}$ for the nominal 1 bar equatorial radius (J. C. Pearl & B. J. Conrath 1991).

The temperature–pressure profile for the model Neptune is shown in Figure 4. The effect of the Ledoux-stable boundary layer on temperature is clearly evident. Our 1 bar temperature

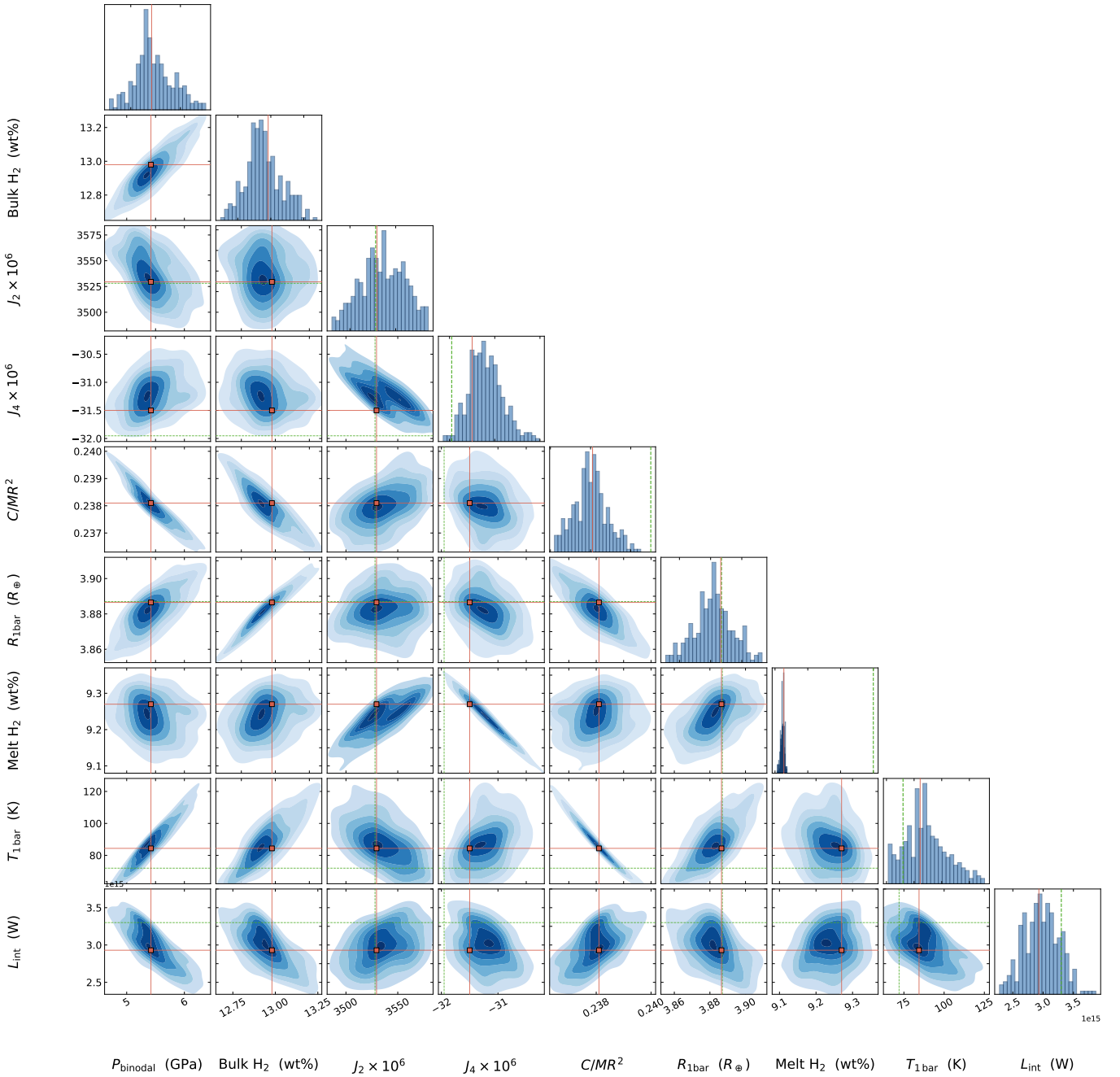


Figure 3. Corner plot for MCMC search over 1100 models for the best fit for Neptune showing the marginalized posterior densities for the parameters of interest. The melt H_2 weight fraction that was not part of the fit is shown for reference. Green dashed lines are target values, and the red squares and lines show the highest likelihood solution (line for melt H_2 is maximum allowed in search).

for Neptune is about 12 K greater than the measured value of 72 ± 2 K (G. F. Lindal 1992).

The influence of the binodal phase boundary on the structure and chemistry of the model Neptune is illustrated in Figure 5. The red curve in the figure is the adiabat within the miscible interior. This path intersects the binodal phase-boundary surface in T , P , and mole fraction of H_2 (x_{H_2}) space (blue dome), at which point the miscible silicate–hydrogen–iron mixture begins to exsolve into the two phases comprising the envelope: an H_2 -rich phase that becomes gas at low pressures, and liquid condensates that we model as silicate rain that settles back to the base of the envelope. At first contact with the binodal, the interior and incipient envelope are

compositionally continuous and the first condensates are displaced to lower H_2 concentrations. As T decreases moving outward into the envelope, the silicate liquids condense out, and in our model are removed from the envelope by settling. With further decreases in T and P higher in the envelope, the relative mass of silicate liquid that rains out of the envelope increases, while the condensate itself rapidly becomes depleted in H_2 (violet curve), leaving the coexisting H_2 -rich fluid phase increasingly enriched in hydrogen (light blue curve).

5.2. Uranus

The results for Uranus are broadly similar to those for Neptune but indicate a greater total hydrogen mass fraction

Table 4

Comparison of Best-fit Interior Model Outputs for Neptune and Uranus

Quantity	Neptune	Uranus
Bulk planet		
Mass ($M_{\oplus}$)	17.150	14.536
Radius at 1 bar ($R_{\oplus}$)	3.863	3.986
Core mass ($M_{\oplus}$)	16.449	13.894
Core radius ($R_{\oplus}$)	3.151	3.114
Core volume fraction	0.537	0.472
Atm. volume fraction (chord)	0.463	0.528
Composition (H_2 mass fraction)		
Whole planet	0.1298	0.1436
Core	0.0927	0.1041
Gas envelope ($H_2 + He$, $\mu = 2.3$)	0.9994	0.9981
Silicate condensates	0.0201	0.0173
Gas mass fraction of planet	0.0393	0.0424
Gas mass percent of planet (%)	3.927	4.235
Atmosphere/envelope		
Surface pressure (GPa)	5.420	7.040
Surface temperature (K)	3256.3	3071.5
Equilibrium temperature (K)	47.0	58.1
Mass of atmosphere (kg)	4.186×10^{24}	3.839×10^{24}
Atm. density, $R < 1$ bar ($kg\ m^{-3}$)	146.68	106.98
Atm. volume, $R < 1$ bar (m^3)	2.854×10^{22}	3.588×10^{22}
Pressure at R_{rcb} (bar)	5	7
$R(1\ bar)/R_c$	1.226	1.280
Ra/Ra_{crit} (envelope)	1.224×10^{12}	5.674×10^9
Boundary layer thickness (m)	105.4	741.1
Ledoux layer above surface (m)	1.219×10^5	1.842×10^5
Internal luminosity L_{int} (W)	2.933×10^{15}	6.646×10^{14}
Core		
Central temperature (K)	10,698.1	8516.8
Bulk core density ($kg\ m^{-3}$)	2899.4	2537.3
Core thermal energy (J)	1.240×10^{33}	8.617×10^{32}
Core grav. potential energy (J)	-2.063×10^{34}	-1.474×10^{34}
Energy budget		
Latent heat to form atmosphere (J)	9.665×10^{31}	8.102×10^{31}
Planet thermal energy (J)	1.276×10^{33}	9.135×10^{32}
Planet grav. potential energy (J)	-2.129×10^{34}	-1.561×10^{34}
Total energy (J)	-2.010×10^{34}	-1.444×10^{34}

(Figure 2). The input parameters are the total mass of $14.536 M_{\oplus}$ and the equilibrium temperature of 58 K.

The maximum likelihood solution is obtained for a binodal pressure of 7.04 GPa, a bulk planet H_2 fraction of 14.36% by mass (as in the case for Neptune, this includes the nominal He fraction by virtue of the adopted molecular weight for the nonmetal fraction), and an Ra/Ra_{crit} value corresponding to a model intrinsic luminosity of 6.65×10^{14} W, indistinguishable from recent estimates (Table 5). The reduced χ^2 for the fit of 0.2 is based on three degrees of freedom.

The Uranus uncertainties used for the MCMC searches follow a similar scheme as for Neptune, with values adjusted to reflect the different observational sources and model systematics. For J_4 , we adopt σ_{J_4} equal to the magnitude of our wind correction ΔJ_4 (Table 2), which is approximately 7 times the Monte Carlo uncertainty of R. G. French et al. (2024) but reflects the systematic uncertainty introduced by the wind

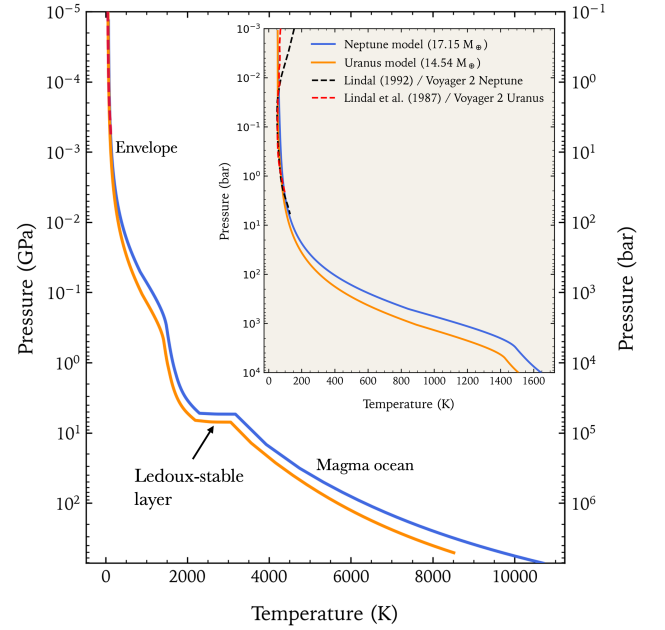


Figure 4. Model temperature–pressure profiles for Neptune and Uranus derived in this study. The regions corresponding to the supercritical silicate–hydrogen–iron magma ocean and the overlying Ledoux-stable layer (boundary layer) in the hydrogen-rich envelope are indicated. The inset shows the upper atmosphere compared with the Voyager 2 observations for both planets.

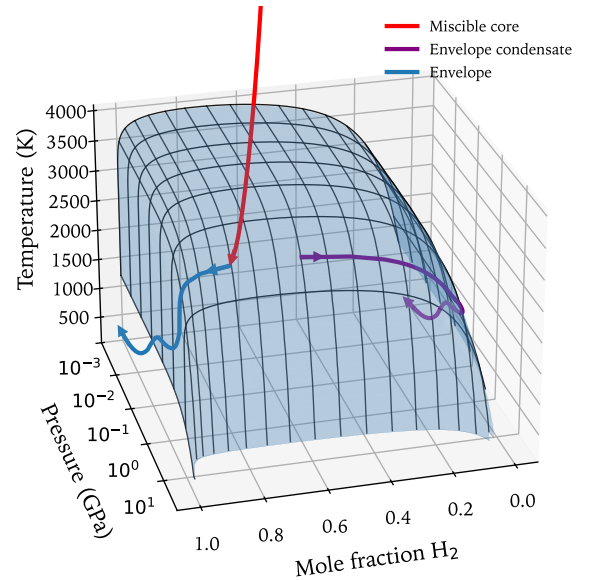


Figure 5. Plot of the binodal phase boundary in T , P , and mole fraction of H_2 space that separates the miscible interior and the outer envelope of Neptune as modeled here. The binodal phase boundary is shown as the blue dome structure. The curves terminating in arrow heads represent the miscible interior (red), H_2 -rich gas (light blue), and liquid silicate rain (violet) with the arrow heads pointing in the direction of greater radial position. The binodal is calculated using the $MgSiO_3$ – H_2 system (T. Gilmore & L. Stixrude 2026), a reasonable approximation to the ternary $MgSiO_3$ – Fe – H_2 system for this purpose (E. D. Young et al. 2025).

correction itself. We similarly broaden σ_{J_2} from the observational value of 4.12×10^{-7} to the magnitude of our wind correction, $|\Delta J_2| = 1.35 \times 10^{-5}$. The wind correction is more than 30 times the observational uncertainty in J_2 , and even a modest fractional uncertainty in the wind-induced contribution (R. G. French et al. 2024) dominates the formal observational

Table 5
Best-fit Solution from the Uranus Interior MCMC

Quantity	Predicted	Target	χ^2
P_{binodal} (GPa)	7.04	...	...
Bulk H_2 (wt%)	14.36	...	...
Ra/Ra_{crit}	5.674×10^9	...	...
J_2 ($\times 10^{-6}$, static)	3522.80	$35,22.76 \pm 13.470$	0.000
J_4 ($\times 10^{-6}$, static)	-30.880	-32.062 ± 3.000	0.155
C/MR^2	0.2275	0.2250 ± 0.00450	0.309
R_{eq} (km)	25561.7	$25,559.0 \pm 20.0$	0.019
$T_{1 \text{ bar}}$ (K)	75.0	76.0 ± 5.0	0.040
L_{int} (W)	6.65×10^{14}	$6.30 \pm 1.5 \times 10^{14}$	0.054
χ_{total}^2			0.577
χ_{red}^2 ($\nu = 3$)			0.192

Note. The reduced χ^2 is computed with $n_{\text{obs}} = 6$ constraints and $n_{\text{free}} = 3$ free parameters. A one-sided prior on the melt-layer H_2 content (ceiling of 12.0 wt%; $\sigma = 1.0$ wt%) is included but inactive at the best fit (predicted = 10.41 wt%).

error. This broadening of σ_{J_2} is specific to Uranus but parallels the treatment of J_4 for both planets ($\sigma_{J_4} = |\Delta J_4|$). For R_{eq} , we adopt $\sigma_R = 20$ km, obtained by adding in quadrature the disparate equatorial and polar 1 bar radius uncertainties reported by G. F. Lindal et al. (1987, ± 4 and ± 20 km respectively). For $T_{1 \text{ bar}}$, we adopt ± 5 K rather than the Lindal observational ± 2 K, matching the Neptune treatment. For C/MR^2 , we adopt a 2% model uncertainty (N. Movshovitz & J. J. Fortney 2022; L. Siebenaler & Y. Miguel 2026), noting that C/MR^2 is not directly measured for Uranus and the adopted value of 0.2250 derives from N. Nettelmann et al. (2013). Five of the six parameters fit by the MCMC match their target values within the original observational 1σ values, regardless of whether those values were increased for the search to include potential model systematic errors. The exception is J_4 , which deviates from the wind-corrected value by 2.5 times the cited observational uncertainty (see Tables 5 and 2), although it is well within potential model uncertainties (Figure 6). This is not surprising given inherent uncertainties in modeling the dynamic effects on J_{2n} values.

The core density of Uranus in our models is less than that for Neptune by about 12% due to the larger mass fraction of H_2 in the core (Table 4). Overall these results show that within this model structure, Uranus is “puffier” than Neptune, commensurate with an overall greater mass fraction of hydrogen. At the same time, the model Ledoux-stable boundary layer is thicker than that for Neptune, with a total boundary layer depth of 186 km for Uranus compared with 129 km for Neptune. This is a physical explanation for the lower intrinsic luminosity for Uranus relative to Neptune provided by our model framework, but is not a prediction. Rather, it is prescribed by our best-fit Rayleigh number ratio.

6. Discussion

When assessing the quality of the fits obtained in these magma-ocean models, it is useful to distinguish between a constrained inverse structural problem, in which observed planetary properties are imposed as constraints on an assumed interior architecture, and a forward physical model, in which a specified structure and set of physical parameters generate observable quantities that are then compared with the data. In many classical ice-giant models, the inverse problem consists

of prescribing observed bulk and atmospheric properties, such as M_p , $R_{1 \text{ bar}}$, $T_{1 \text{ bar}}$, and ω , and then solving for internal structural parameters that reproduce the measured gravity field. For example, N. Nettelmann et al. (2013) use this framework to constrain the metallicities of two envelope layers Z_1 and Z_2 , the core mass M_c , and the pressure P_{1-2} at the transition between the two envelope layers for their model structure $\mathcal{M}$. Formally, this may be viewed as a constrained inverse problem in which one solves for Θ_{model} such that $\mathbf{d}_{\text{pred}}(\Theta_{\text{model}}, \mathbf{x}_{\text{fixed}}) = \mathbf{d}_{\text{obs}}$, where $\Theta_{\text{model}} = \{Z_1, Z_2, M_c, P_{1-2}\}$, the fixed inputs are $\mathbf{x}_{\text{fixed}} = \{M_p, R_{1 \text{ bar}}, T_{1 \text{ bar}}, \omega\}$, and the data to match are the gravity harmonics $\mathbf{d}_{\text{obs}} = \{J_2, J_4\}$. The fixed inputs are in effect boundary conditions rather than data to fit, so by matching two gravity harmonics $\mathbf{d}_{\text{obs}} = \{J_2, J_4\}$ while solving for the four structural features comprising Θ_{model} , the system is underdetermined by two degrees of freedom.

Because the physical chemistry of miscibility links structure and chemistry in our models, our approach is subtly different. In our calculation the observables are not imposed. Rather, for each trial set of fit parameters, the model returns a forward prediction for the full observable data vector, and the quality of the fit is determined by the likelihood $\mathcal{L}(\Theta_{\text{model}}) = p(\mathbf{d}_{\text{obs}} | \Theta_{\text{model}}, \mathcal{M})$ given the magma-ocean model $\mathcal{M}$ with fit parameters $\Theta_{\text{model}} = \{P_{\text{binodal}}, x\text{H}_2, Ra/Ra_{\text{crit}}\}$. Simultaneous agreement between the predicted and observed values, $\mathbf{d}_{\text{obs}} = \{R_{1 \text{ bar}}, T_{1 \text{ bar}}, J_2, J_4, C/MR^2, L_{\text{int}}\}$, is therefore not enforced a priori, but is instead a consequence of finding regions of parameter space with high likelihood. This predictive approach has precedent in the ice-giant literature. E. Bailey & D. J. Stevenson (2021) and M. Cano Amoros et al. (2024) let H_2 - H_2O immiscibility physics determine the locations of compositional gradients, rather than imposing them as free parameters, and L. Scheibe et al. (2021) computed the 1 bar radius as a forward output of thermal-evolution calculations, retaining models that did not match $R_{1 \text{ bar}}$ to illustrate the physics. They invoked a thermal boundary layer similar in concept to that at the base of the H_2 envelope in this study. Our approach is in the spirit of this same Bayesian model evaluation, in which the data are used to assess the likelihood of a specified physical hypothesis, rather than to invert for the parameters of an assumed structure.

It is in this context that we can evaluate the fits of the models to both planets. The model reproduces Uranus’ six target observable values to well within their respective target uncertainties that include potential systematic errors, and five of the six are within purely observational uncertainties, the very tight nominal observational uncertainty for J_4 being the exception. The model also reproduces Neptune’s observables to within their observational uncertainties, with a single exception: the predicted 1 bar temperature is high by approximately 12 K.

The ~ 12 K excess in $T_{1 \text{ bar}}$ for Neptune is consistent with potential systematic uncertainties in the gray thermal opacity κ_{th} used in our atmosphere model, exacerbated by L_{int} . In the two-stream gray solution at large optical depth, we have

$$T_{1 \text{ bar}}^4 \simeq T_{\text{eq}}^4 + \frac{3 L_{\text{int}} \kappa_{\text{th}} P_{1 \text{ bar}}}{16 \pi \sigma G M}. \quad (6)$$

While $T_{1 \text{ bar}}$ scales weakly with opacity ($T_{1 \text{ bar}} \propto \kappa_{\text{th}}^{1/4}$ in the internal-heating term), the underlying κ_{th} uncertainties are large. As an example, an unconstrained difference in the Rosseland-mean opacity of $2.5\times$, well within the spread of

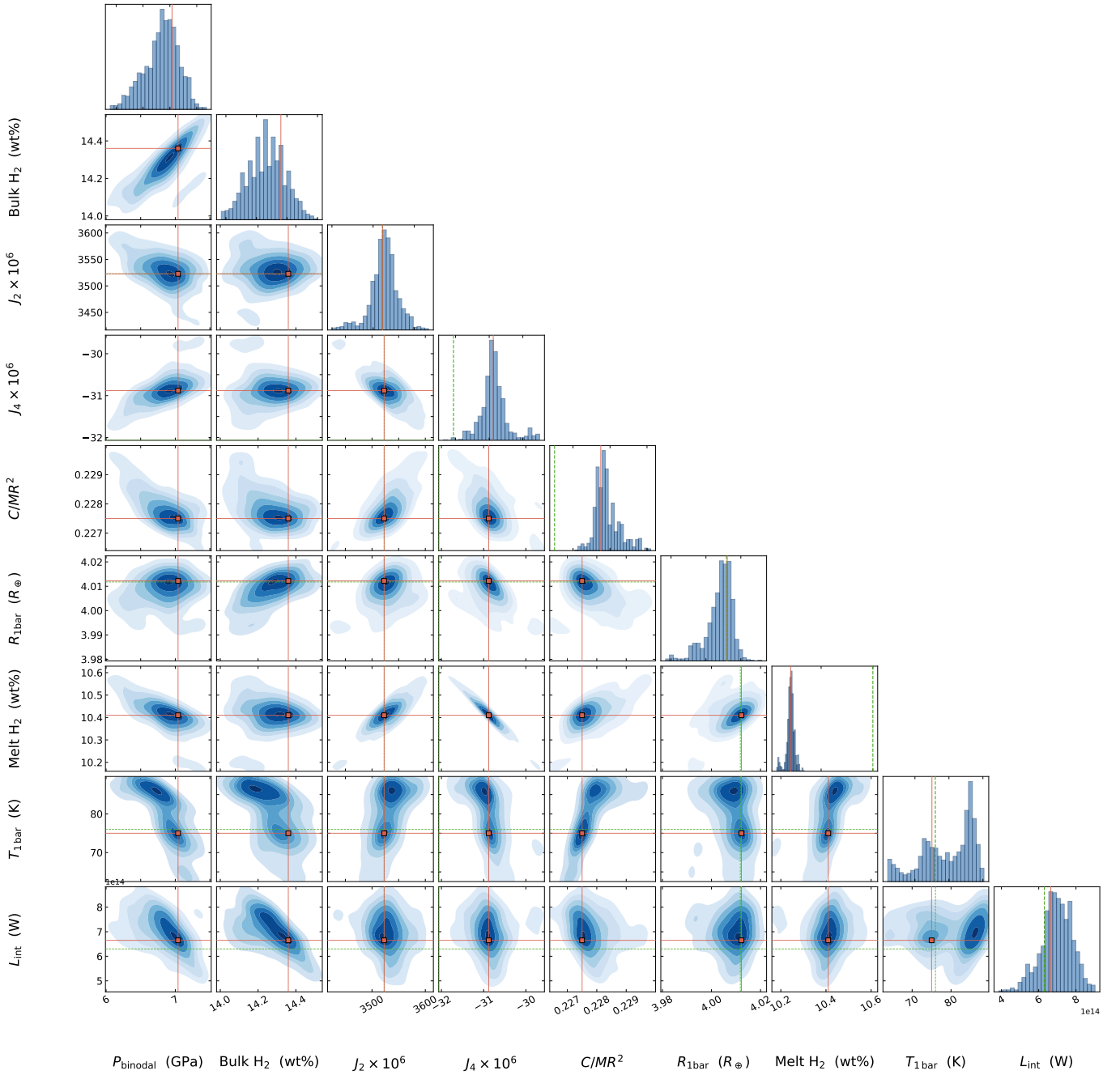


Figure 6. Corner plot for MCMC search over 1008 models for the best fit for Uranus showing the marginalized posterior densities for the parameters of interest. The melt H_2 weight fraction that was not part of the fit is shown for reference. Green dashed lines are target values, and the red squares and red lines show the highest likelihood solution.

plausible gray averages for H_2 – H_2 collision-induced absorption (CIA), CH_4 near-infrared bands, and aerosol contributions (D. Valencia et al. 2013), produces an ~ 15 K shift in $T_{1\text{ bar}}$ that is comparable in magnitude to the residual observed for Neptune. We interpret the $T_{1\text{ bar}}$ excess as the signature of inherent uncertainties in atmosphere-specific opacities rather than a clear structural defect in the interior model.

That the miscible magma-ocean model satisfies the joint constraints of bulk density, intrinsic luminosity, atmospheric thermal structure, and gravity harmonics for Uranus is notable in light of the long-recognized difficulty of constructing interior models that simultaneously reproduce these quantities (W. B. Hubbard & M. S. Marley 1989; R. Helled &

J. J. Fortney 2020). We emphasize that our model reproduces the observed low L_{int} by fitting the width of the boundary layer at the base of the atmosphere and so it is not a prediction of the model. The long-standing question of why Uranus is so much fainter than Neptune (K. Kurosaki & M. Ikoma 2017; L. Scheibe et al. 2021) remains an open question. Our result does lend additional impetus to the concept of compositional gradients suppressing convection in the envelope (A. Vazan & R. Helled 2020) as the source of the low luminosity.

Aside from simply providing another model for Uranus and Neptune that fits their observed characteristics, the magma-ocean giant model has several ancillary features that recommend it. One is that it highlights a potential connection

between gas-dwarf planets in general, from sub-Neptunes to Neptunes, with no requirement for a fundamental distinction in their formation mechanisms and resulting structures. There is merit to deriving successful models that depend as much as possible on physical chemistry and material properties; it is preferable for the EoSs to do the work of matching the planet properties. Another feature to recommend the magma-ocean giant model is that the basic features of the planets are reproduced using just three tunable parameters, pressure at the binodal, the bulk hydrogen fraction, and the Ra/Ra_{crit} ratio for the envelope that influences the intrinsic luminosity (i.e., the boundary layer). The Bayesian information criterion (BIC) applied to Bayes factors is an example of a quantitative means of assessing the relative merits of complexity versus parsimony in model selection (K. Aho et al. 2014). While fewer parameters is not definitive evidence for veracity, it is a benefit when choosing among models based on material properties. Additional evidence for the presence of extant supercritical silicate–hydrogen magma oceans in the interiors of Uranus and Neptune and connections to extrasolar planets are discussed in what follows.

6.1. Atmosphere Chemistry

The atmospheric chemistries of both Uranus and Neptune are consistent with predictions for magma-ocean worlds. James Webb Space Telescope (JWST) observations of the sub-Neptunes K2-18 b and TOI-270 d detect abundant CH_4 and CO_2 but no NH_3 (N. Madhusudhan et al. 2023; B. Benneke et al. 2024; M. Holmberg & N. Madhusudhan 2024). NH_3 depletion is a natural consequence of the high solubility of nitrogen in silicate melt under the reducing conditions imposed by a thick H_2 envelope in contact with a magma ocean (O. Shorttle et al. 2024; A. Werlen et al. 2026), precisely the scenario modeled here, and the atmospheric C/O ratio of such planets is set by equilibrium with the underlying melt rather than inherited from the protoplanetary disk, with CH_4 being the dominant carbon carrier in the deep atmosphere (A. Werlen et al. 2025b). The nondetection of NH_3 in both Uranus and Neptune (I. de Pater et al. 1991) and their observed CH_4 abundances are hence qualitatively consistent with magma-ocean interiors, lending independent support for the hypothesis that the solar system ice giants belong to the same class of objects as the sub-Neptune gas dwarfs characterized by JWST.

To test this quantitatively, we applied the `ChemicalGlobalEquilibrium` code³ (S. L. Grimm et al. 2026), a global chemical equilibrium code originally developed for super-Earths and sub-Neptunes (H. E. Schlichting & E. D. Young 2022), to the conditions at the binodal interface in our models. The code equilibrates a hydrogen-rich atmosphere with coexisting silicate and metal phases. For the ice giants, we exclude a separate metal phase, since metal is expected to be miscible with the silicate-rich interior under Uranus- and Neptune-like conditions (E. D. Young et al. 2025). This reduces the problem to a two-phase silicate–gas system. We base our chemical network on the one used by A. Werlen et al. (2026), which includes sulfur-, nitrogen-, and oxygen-bearing species, but exclude the metal phase in our calculations.

The mantle composition is dominated by MgSiO_3 (mole fraction = 0.9211), with trace contributions from MgO , SiO_2 ,

FeO , FeSiO_3 , Na_2O , and Na_2SiO_3 . The gas phase is initialized with carbon, sulfur, and nitrogen in solar proportions ($\text{CO} = 5.4 \times 10^{-4}$, $\text{H}_2\text{S} = 2.64 \times 10^{-5}$, and $\text{N}_2 = 6.76 \times 10^{-5}$). The resulting equilibrium mixing ratios above the magma ocean are dominated by H_2 ($x = 0.9782$) and H_2O ($x = 0.0613$), with CH_4 ($x = 0.0149$) the principal carbon carrier, H_2S ($x = 5.08 \times 10^{-4}$) the principal sulfur carrier, and nitrogen-bearing species essentially absent (NH_3 : $x = 4.16 \times 10^{-9}$; N_2 : $x = 4.16 \times 10^{-9}$).

These results reproduce the key observational signatures of the ice-giant atmospheres. CH_4 is the dominant carbon carrier with an enrichment of $\sim 28\times$ solar relative to hydrogen, broadly consistent with the $40\text{--}50\times$ and $\sim 80\times$ solar enrichment observed in Uranus and Neptune, respectively. Sulfur appears primarily as H_2S at $20\times$ solar enrichment, matching the $20\text{--}30\times$ solar values inferred for both planets (I. de Pater et al. 1991). Nitrogen-bearing species are virtually absent, consistent with the nondetection of NH_3 in both planets. The model also produces water above the magma ocean, but it is expected to remain unobservable. Even under vigorous vertical mixing—which can loft water-rich vapor upward from the deep atmosphere—the steep decline of water’s saturation vapor pressure with temperature imposes a cold trap: rising atmosphere saturates and condenses at the level (hundreds of bar) where the temperature profile crosses the water condensation curve, limiting the water partial pressure at its low saturation value regardless of the mixing strength embodied by eddy diffusion K_{zz} , for example. CH_4 , by contrast, has a high enough saturation vapor pressure at relevant low temperatures to persist as a detectable vapor in the observable 1–10 bar region (N. Fray & B. Schmitt 2009). The equilibrium chemistry of a magma-ocean interface therefore reproduces the observed pattern of CH_4 and H_2S enrichment together with nitrogen depletion, without requiring fine-tuning of atmospheric or disk conditions. Of course nondetection in the upper atmospheres is not proof of extant magma oceans nor is it proof of depletion of the atmospheres as a whole (C. Li et al. 2017; T. Guillot 2022), but it is consistent with expectations where magma oceans are present.

6.2. Hydrogen Accretion and Links to Sub-Neptunes

Later stages of planet formation are likely dominated by the merging of planetary embryos, each with their own H_2 -rich primary atmosphere. The observed spread in inferred H_2 mass fractions among sub-Neptunes spans roughly an order of magnitude, from $\sim 1\%$ to $\sim 10\%$ by mass (J. L. Bean et al. 2021), providing some clue as to the connection between Neptune and Uranus and the broader sub-Neptune population.

To explore this connection, we performed a simple stochastic numerical accretion experiment in which planets of five target masses (2, 5, 10, 15, and $20 M_{\oplus}$) are assembled by sequential addition of embryos drawn at random from a log-uniform mass distribution (minimum mass of $0.1 M_{\oplus}$; maximum mass of one-half the target mass), reflecting the embryo-dominated late stages of accretion. The characteristic disk temperature at the planet’s location is assigned to reflect formation distance, ranging from 300 K for $2 M_{\oplus}$ to 50 K for the most massive planets motivated by the architecture of the solar system. The precise radial temperature profile is not essential to the salient features of the calculations presented here. During each accretion event two processes occur. First, the impact partially strips the preexisting H_2 envelope, so that

³ https://github.com/ExoInteriors/GlobalChemicalEquilibrium_Release

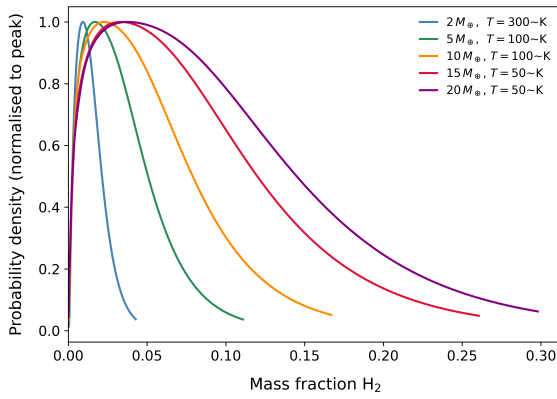


Figure 7. Results of a toy accretion model illustrating the expected relationship between planet mass and hydrogen fraction.

only a fraction x_{r_i} of the envelope mass survives. Second, the accreted embryo contributes additional H_2 according to the cooling-limited scaling of S. Ginzburg et al. (2016):

$$x_{H_2}(m) = 0.02 \left(\frac{m}{M_{\oplus}} \right)^{0.8} \left(\frac{T}{1000 \text{ K}} \right)^{-1/4} \left(\frac{t_{\text{disk}}}{1 \text{ Myr}} \right)^{1/2}, \quad (7)$$

where T is the characteristic disk temperature and t_{disk} is the gas disk lifetime. The envelope mass therefore evolves as

$$M_{\text{env}}^{(k+1)} = x_{r_k} M_{\text{env}}^{(k)} + \Delta M_k, \quad (8)$$

where $\Delta M_k = f(m_k)m_k$ is the H_2 increment contributed by the k th embryo of mass m_k . Iterating through N accretion events gives

$$M_{\text{env}}^{(N)} = \left(\prod_{j=0}^{N-1} x_{r_j} \right) M_{\text{env}}^{(0)} + \sum_{k=0}^{N-1} \left[\Delta M_k \prod_{j=k+1}^{N-1} x_{r_j} \right], \quad (9)$$

so that the final envelope is the sum of all accreted increments, each weighted by the retention factors of all subsequent impacts. The final H_2 mass fraction is then

$$x_{H_2} = \frac{M_{\text{env}}}{M_{\text{core}} + M_{\text{env}}}. \quad (10)$$

Because the retention factor x_{r_k} enters multiplicatively at each step, the central limit theorem implies that $\ln(M_{\text{env}})$ behaves approximately as a sum of random contributions, and the resulting H_2 mass fraction distribution is approximately lognormal. The width of the distribution at each target mass reflects the stochastic variability in both the sequence of embryo masses and the sequence of retention factors.

The results, shown in Figure 7, demonstrate that the median H_2 fraction increases systematically with planet mass, from $\sim 1\%$ for $2 M_{\oplus}$ to $\sim 9\%$ for $20 M_{\oplus}$. Two effects drive this trend. First, and most importantly, more massive planets accrete larger individual embryos, and the scaling $x_{H_2} \propto m^{0.8}$ means each event contributes disproportionately more H_2 per unit solid mass as embryo masses grow (S. Ginzburg et al. 2016). Second, more massive planets form at cooler disk temperatures in this model, further boosting the accreted envelope fraction via the $T^{-1/4}$ dependence. The $\sim 13\%$ H_2 mass fractions inferred for Uranus and Neptune from our interior models are fully consistent with this picture, sitting at the high-mass end of the predicted distribution. The magma-ocean giant model is

thus both parsimonious—reproducing the structure, chemistry, and thermal states of both planets with just three free parameters—and physically motivated, connecting directly to the broader population of gas-dwarf planets and their formation.

6.3. Implications for Magnetic Fields

Measurements of Uranus’ and Neptune’s magnetic fields also suffer from the limitations of only single flybys by Voyager II. However, despite only having information on the large-scale structure of their fields, it is evident that they differ dramatically from those of the two gas giants, and Earth, as they are not dominated by dipoles (N. F. Ness et al. 1986, 1989). Dynamo studies suggest that the generation of so-called “multipolar” magnetic fields does not require a specific dynamo region geometry, or a more slowly rotating system (e.g., T. Gastine et al. 2012; R. K. Yadav et al. 2013; K. M. Soderlund et al. 2025).

DFT-MD simulations have shown that the magnetic Reynolds numbers for molten silicates at relevant conditions are sufficient to produce magnetic fields (F. Soubiran & B. Militzer 2018). Silicate melts enriched in iron exhibit a further increase in magnetic Reynolds number (F. Dragulev & L. Stixrude 2025), and we anticipate that the presence of hydrogen would amplify this effect. Thus, there is no reason to doubt that multipolar magnetic fields may be generated in the models presented here.

6.4. Evolution

We calculated the time evolution of Neptune based on our model. The quasi-static evolution curves are obtained by specifying the planet mass M_p , bulk hydrogen mass fraction, equilibrium temperature T_{eq} , and an initial binodal pressure that results in the model binodal pressure today. The binodal pressure is uniquely correlated with total energy and thermal state. For each step we calculated the total energy of the planet E_p and the associated intrinsic luminosity L_{int} . The time step associated with a step from P_i to P_{i+1} was obtained from energy balance:

$$dt = \frac{|E_p(P_{i+1}) - E_p(P_i)|}{L_{\text{int}}}, \quad (11)$$

where L_{int} is output from the prior time step. Total energy E_p is calculated from

$$E_p = \int_0^{R_p} \left(-\frac{GM(r)}{r} + cT(r) \right) \rho(r) 4\pi r^2 dr, \quad (12)$$

where c is the specific heat capacity for the melt or envelope using the EoSs and species described above. Pressure steps are sufficiently small to ensure that dt is approximately 4 to 5 orders of magnitude smaller than the cooling time constant $\tau_{\text{cool}} = |E_p|/L_{\text{int}}$.

Results indicate that the changes in radius, temperatures, and partitioning of hydrogen have been relatively modest over the last 4.5 Gyr (Figure 8). Comparing the temporal changes in radius with the mass fraction of H_2 in the core illustrates that the exsolution of hydrogen plays a role in limiting the decrease in radius during cooling. Unknown, highly irreversible events that may have occurred very early in the formation of the planet are not included here by necessity.

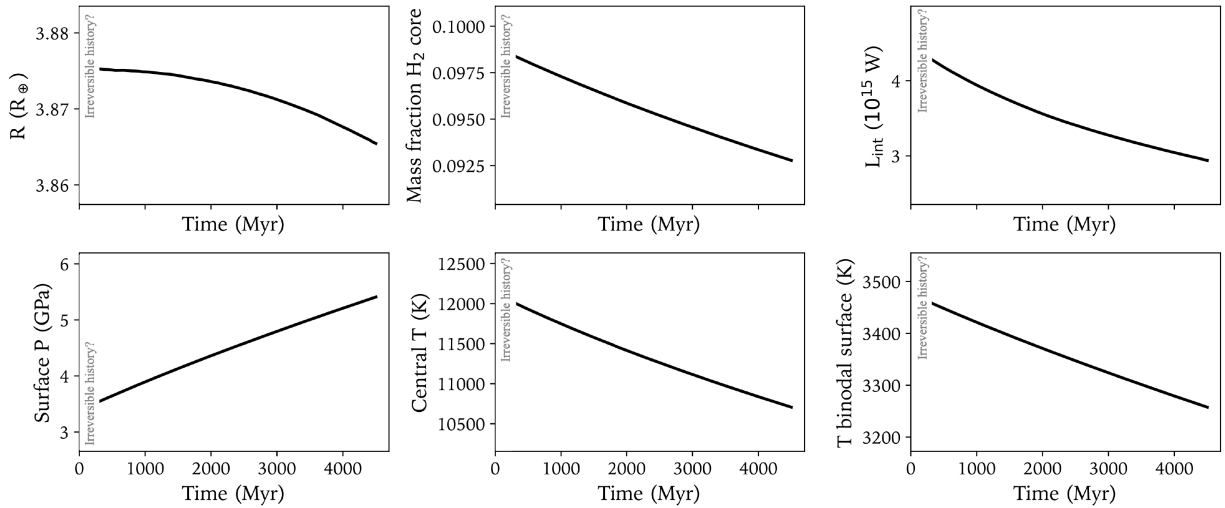


Figure 8. Calculated evolution of Neptune over the last 4.5 Gyr. Irreversible processes that predate the present-day structure of the planet are not resolved with this quasi-static calculation.

6.5. Limitations

The specifics of the magma-ocean giant models for Uranus and Neptune are not unique. Our numerical experiments show that reasonable changes to the supercritical melt EoS parameters result in similarly successful fits to the planet observables but with similar, although still distinct, P_{binodal} and x_{H_2} values. These models with different binodal pressures and bulk hydrogen fractions have essentially the same pressure–density curves as those obtained here. The precise parameters that determine these curves are affected by the supercritical melt EoS (see Appendix D). The uniqueness of the solutions will be improved as the EoS for miscible MgSiO_3 , H_2 , and Fe becomes better constrained (work in progress). Until then, we can conclude with confidence that the observable characteristics of these planets can be fit with extant, supercritical, hydrogen-rich magma oceans comprising the inner three-quarters of their radii, a boundary layer stable against convection due to molecular-weight gradients above the magma oceans, and an overlying hydrogen-rich envelope. The precise concentrations of hydrogen, and precise pressures and temperatures at the interfaces between the magma oceans and the envelopes will require refinement of the melt phase EoS. For now, we can say that there is a systematic absolute uncertainty of order 2% (e.g., mass fraction of H_2 for the core of Neptune is $9.5\% \pm 2\%$) in the mass fraction of total hydrogen for both planets in our models based on uncertainties in the EoS for the supercritical melt phase (Appendix D).

7. Conclusions

Modern cosmochemical and astrophysical constraints on the relative abundances of silicates and ices suggest that the progenitors of the ice-giant planets in the solar system may have been “rockier” than previously asserted. At the same time, our understanding of the physical chemistry of mixtures of rocky material and hydrogen at high temperatures and pressures has evolved. Against this backdrop, we have the opportunity to use the ice giants as a test of models put forward previously for sub-Neptune planets. We have asked and answered the question of whether the observable constraints for Neptune and Uranus can be explained by interiors composed of extant supercritical mixtures of silicate and

hydrogen, overlain by H_2 -rich envelopes. This same basic structure has been proposed for sub-Neptunes for the same reasons; temperatures and pressures are sufficiently high in these bodies that miscibility between rock and hydrogen should obtain.

The apparent disparity in the orbital distances from their host stars does not necessarily translate to fundamentally distinct interior structures for the solar system ice giants and sub-Neptunes. Accreted water ice is reworked by magma ocean–hydrogen envelope reactions, partly, if not entirely, erasing the signature of those ices (A. Werlen et al. 2025a, 2026). Based on the chemical model presented here and elsewhere, the primary manifestation of the progenitor ices will be to provide an additional source of hydrogen indistinguishable from that provided by H_2 and enhance the overall oxidation states of these planets, rather than altering their basic structure.

While this is just one of a number of models that successfully reproduce the observed features of Neptune and Uranus, this model has several aspects to recommend it. One is the connection with other gas-dwarf planets: it is not clear that ice giants and sub-Neptunes should be fundamentally different simply because of their distances from their host star. Related to this is the fact that the most basic chemical features of the ice giants resemble those of gaseous sub-Neptunes, perhaps indicating similar boundary conditions for the chemistry of the atmospheres imposed by the magma oceans. Another potential favorable feature of the magma-ocean giant proposition is that the use of the phase equilibria involving silicates and hydrogen, arguably the dominant chemical constituents of the planets, limits the number of tunable parameters to reproduce their observed densities, external gravity, and calculated moments of inertia to just two. Simpler models are in many ways easier to test, and are less susceptible to overfitting. However, more work is required to improve relevant EoSs and address the initial conditions for gas-dwarf planet formation.

Acknowledgments

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Appendix A Planet Structure

Many of the methods for calculating the planet models used here have been described previously (J. G. Rogers et al. 2025; E. D. Young et al. 2025; E. D. Young & A. Werlen 2026). However, there are some modifications to the prior published work, and so the models are explained here. To derive the core, we solve the system of equations (S. Seager et al. 2007):

$$\frac{dm}{dr} = 4\pi r^2 \rho, \quad (\text{A1})$$

$$\frac{dP}{dr} = -\frac{Gm\rho}{r^2}, \quad (\text{A2})$$

and

$$\left(\frac{dT}{dr}\right)_s = -\frac{\gamma(\eta) T}{K_S(\eta)} \rho g, \quad (\text{A3})$$

where m is the mass contained within radius r , ρ is mass density, $\gamma(\eta)$ is the Grüneisen parameter that varies with P and T through the compressibility factor $\eta = \rho/\rho_0$, and $K_S(\eta)$ is the isentropic bulk modulus. We assume the core is fully convective, with a limiting isentropic dT/dP gradient. In practice, we evaluate the bulk modulus using $K_S(\eta) \equiv \eta d(P(\eta, T_S))/d\eta$ where $P(\eta, T_S)$ is the total isentropic pressure along the adiabat with temperature T_S . Numerically integrating Equations (A1)–(A3) through one or more layers (with or without a metal core) yields a density and temperature profile for the planet interior.

For the structure of the envelope we integrate numerically the hydrostatic and mass conservation equations as in the interior (Equations (A1) and (A2)) using the tabulated function from G. Chabrier et al. (2019) for the EoSs. The temperature structure of the envelope is obtained by numerically integrating the coupled hydrostatic, mass conservation, and energy transport equations outward from the magma-ocean surface, followed by an iterative inward integration from the optically thin outer atmosphere. Throughout most of the envelope, energy transport is convective, and the temperature gradient is therefore taken to be the convective (pseudoadiabatic) gradient

$$\frac{dT}{dr} = -\nabla T_{\text{conv}}, \quad (\text{A4})$$

where ∇T_{conv} is evaluated using a moist pseudoadiabat appropriate for an H₂–silicate vapor mixture (R. J. Graham et al. 2021), scaled to the H₂ EoS (G. Chabrier et al. 2019). Above the convective layer where the gas is still collisional, radiative diffusion applies:

$$\nabla T_{\text{rad}} = -\frac{3\kappa\rho L_{\text{int}}}{64\pi\sigma T^3 r^2}, \quad (\text{A5})$$

where κ is the local opacity, ρ is the density, and L_{int} is the intrinsic luminosity carried upward from the interior. The envelope transitions smoothly into a radiative regime by relaxing the temperature gradient with Eddington’s two-stream solution for a plane-parallel gray atmosphere

(T. Guillot 2010):

$$T^4(\tau) = \frac{3}{4} T_{\text{int}}^4 \left(\tau + \frac{2}{3} \right) + \frac{3}{4} T_{\text{eq}}^4 \left[\frac{2}{3} + \frac{2}{3\gamma_{\text{vis}}} + \left(\frac{\gamma_{\text{vis}}}{3} - \frac{2}{3\gamma_{\text{vis}}} \right) e^{-\gamma_{\text{vis}}\tau} \right], \quad (\text{A6})$$

where τ is the optical depth measured downward from the top of the atmosphere, T_{eq} is the equilibrium temperature set by stellar irradiation, T_{int} is the intrinsic temperature associated with the interior luminosity, and γ_{vis} in this context is the visible-to-thermal opacity ratio (we used one for these calculations, but the results do not depend critically on this value). We impose conservation of intrinsic luminosity, rather than a constant intrinsic flux as assumed in the plane-parallel equations, by accounting for the area dependence of the intrinsic flux $f_{\text{int}}(r) = \sigma T_{\text{int}}^4(r)$ such that $T_{\text{int}}(r) = T_{\text{int},s}(r_s/r)^{1/2}$, where subscript s signifies evaluation just above the supercritical magma ocean. At large infrared optical depth ($\tau \gg 1$), Equation (A6) approaches the radiative-diffusion limit, yielding a radiative temperature gradient consistent with Equation (A5) and thus ensuring consistency with deep-envelope energy transport. This approach affords a smooth and numerically stable transition from the deep convective adiabat to the irradiated upper atmosphere.

Appendix B Luminosity

Specifying an initial entropy and intrinsic luminosity, L_{int} , based on inferences about retention of heat during accretion does not guarantee self-consistency; a specified L_{int} may exceed, or fall short of, the flux the atmosphere can transmit given its opacity structure and boundary conditions, leading to either interior heating or unphysical cooling timescales. A closure for this problem, and thus self-consistent values for L_{int} , can be obtained from consideration of boundary layer effects.

We include the expected formation of a boundary layer at the base of the hydrogen-rich envelope between the supercritical melt and the overlying convective region in the envelope (G. Ahlers et al. 2009), augmented by the hindrance of convection due to the mass loading of heavy elements at relatively high temperatures (W. Misener et al. 2023). The lower boundary layer restricts the intrinsic upward luminosity that comprises heat transfer from the core to the atmosphere, forming a thermal bottleneck.

The minimum thickness of the basal boundary layer separating the interior (core) and the overlying envelope, δ , comes from consideration of the boundary for Rayleigh–Bénard convection in the envelope. The width of this boundary is controlled by the vigor and scale of the convection and is given by scalings similar to

$$\delta \sim nh(Ra/Ra_{\text{crit}})^{-1/3}, \quad (\text{B1})$$

where h is a characteristic convecting length scale, taken to be some multiple n of the pressure scale height evaluated at the base of the atmosphere, Ra is the Rayleigh number, and Ra_{crit} is the critical Rayleigh number for the convecting H₂–rich envelope (S. Grossmann & D. Lohse 2000; R. S. Long et al.

2020). The Rayleigh number is computed as

$$Ra = \frac{g\alpha\Delta Th^3}{\nu\kappa}, \quad (\text{B2})$$

where h is here taken to be $5H$ for scale height H , ΔT is the temperature contrast across the layer, ν is the kinematic viscosity, κ is the thermal diffusivity, and α is the thermal expansivity. Use of the adiabat for ΔT , rather than super-adiabatic temperatures, likely overestimates Ra , resulting in underestimates for the widths of the boundary layers (see below). Kinematic viscosities for H_2 are estimated here using a Sutherland-law fit for dynamic viscosities η :

$$\eta(T) = \eta_0 \left(\frac{T}{T_0} \right)^{3/2} \frac{T_0 + S_T}{T + S_T}, \quad (\text{B3})$$

with reference viscosity η_0 at reference temperature T_0 and Sutherland temperature S_T for H_2 (J. Braun et al. 2018). The kinematic viscosities are then

$$\nu = \frac{\eta}{\rho}. \quad (\text{B4})$$

Thermal expansivities required to evaluate Ra are obtained using

$$\alpha = -\frac{1}{T} \left(\frac{\partial \log \rho}{\partial \log T} \right)_p, \quad (\text{B5})$$

based on the EoS tables for hydrogen (G. Chabrier et al. 2019). We adopt a thermal diffusivity κ of $1.0 \times 10^{-5} \text{ m}^2 \text{ s}^{-1}$ for the purpose of evaluating approximate values for the Rayleigh number. For nominal parameters for Neptune, e.g., $\Delta T = 3000 \text{ K}$, surface pressure $P_s = 6 \text{ GPa}$, planet mass $M_p = 17.1 M_\oplus$, and density $\rho = 230 \text{ kg m}^{-3}$, one obtains $Ra \simeq 10^{32}$.

The absolute value of Ra_{crit} is subject to large systematic uncertainty. Rotating convection theory predicts $Ra_{\text{crit}} \approx 8.7 \text{ Ek}^{-4/3}$ as $\text{Ek} \rightarrow 0$ (S. Chandrasekhar 1961), where the Ekman number is $\text{Ek} = \nu/(2\Omega H^2)$ and Ω is the angular rotation rate. At the base of the planetary envelope $\text{Ek} \sim 10^{-18}$, so that extrapolating the Chandrasekhar scaling involves 7 orders of magnitude beyond the regime where it is known to hold. The actual onset and vigor of convection in a compressible, compositionally stratified planetary envelope is an open theoretical problem. The values of Ra/Ra_{crit} derived in our models should therefore be understood as effective parameters that reproduce the observed heat flows, rather than as first-principles predictions from rotating convection theory. For Neptune's rotation rate we obtain a value for Ra_{crit} of 10^{22} , and thus $Ra/Ra_{\text{crit}} \sim 10^9$. However, reproducing Neptune's observed internal luminosity $L_{\text{int}} \sim 3 \times 10^{15} \text{ W}$ requires $Ra/Ra_{\text{crit}} \simeq 10^{12}$, yielding $\delta \approx 100 \text{ m}$. The analogous fit for Uranus, whose internal heat flux is several times smaller, requires $Ra/Ra_{\text{crit}} \simeq 10^{10}$, yielding $\delta \approx 700 \text{ m}$. The physically meaningful and robust result is the ratio $(Ra/Ra_{\text{crit}})_{\text{Neptune}}/(Ra/Ra_{\text{crit}})_{\text{Uranus}} \sim 100$. The ratio is independent of the absolute calibration of the Chandrasekhar scaling. The factor of ~ 100 difference in criticality for convection between the two planets is physically plausible if their deep envelopes differ in viscosity, detailed composition, or super-adiabatic temperature excess, all of which would affect the boundary layer thickness.

The model intrinsic luminosity L_{int} is determined from the blackbody emission $L_{\text{bb}} = 4\pi r_s^2 \sigma T_s^4$ using the ability of the magma ocean to radiate through the boundary layer immediately above its surface, which for a gray atmosphere is

$$L_{\text{int}} = L_{\text{bb}} \left(\frac{4}{3(\tau_{\text{BL}} + 1)} \right). \quad (\text{B6})$$

The intrinsic temperature at the surface satisfies

$$L_{\text{int}} = 4\pi r_s^2 \sigma T_{\text{int},s}^4. \quad (\text{B7})$$

This formulation allows for self-consistency between the surface of the magma ocean and the upper atmosphere, accommodating the link between intrinsic luminosity in the regime of vigorous convection. Rather than resolving the temperature gradient in this often very narrow boundary, its influence as a governor on radiative flux is captured through Equation (B6).

The basal radiative layer is further stabilized and enhanced in elevation by the molecular-weight gradients that attend the interaction between the supercritical magma ocean and the overlying envelope. As a result, the boundary between the surface of the magma ocean and the convective layer in the envelope can be both pronounced and explicitly resolvable in our models. Rather than allowing arbitrarily large intrinsic luminosities that would dynamically destroy stratification, we adopt a quasi-static closure in which a Ledoux-stable layer, when permitted thermodynamically, forms, persists, and therefore acts as a finite thermal resistance that limits the transmitted flux.

Instability against convection in the presence of composition gradients is governed by the Ledoux criterion:

$$\nabla_{\text{rad}} > \nabla_{\text{ad}} + \frac{\phi}{\delta} \nabla_{\mu}, \quad (\text{B8})$$

where

$$\nabla_{\mu} \equiv \frac{d \ln \mu}{d \ln P}, \quad (\text{B9})$$

$$\delta \equiv - \left(\frac{\partial \ln \rho}{\partial \ln T} \right)_{P,\mu}, \quad (\text{B10})$$

and

$$\phi \equiv \left(\frac{\partial \ln \rho}{\partial \ln \mu} \right)_{P,T}. \quad (\text{B11})$$

Here μ is the mean molecular weight. A positive molecular-weight gradient ($\nabla_{\mu} > 0$) increases the critical temperature gradient required for convective instability by the factor $(\phi/\delta)\nabla_{\mu}$. The ratio (ϕ/δ) is unity for an ideal gas and remains of order unity for reasonable values of the compressibility factor Z , so in practice we simplify the Ledoux criterion (Equation (B8)) to $\nabla_{\text{rad}} > \nabla_{\text{ad}} + \nabla_{\mu}$. Where the Ledoux-stable region spans a finite radial extent, it further limits L_{int} for the overlying atmosphere layers. The intrinsic luminosity is limited by the requirement that the layer remain marginally Ledoux stable. In this radiative layer, the upward energy flux is given by the diffusion approximation, written in terms of

dimensionless gradient ∇ as

$$F_{\text{rad}} = \frac{4\sigma T^4}{3\kappa\rho H} \nabla, \quad (\text{B12})$$

where $H = P/(\rho g)$ is the pressure scale height and $\nabla \equiv d \ln T / d \ln P$. The maximum radiative flux consistent with stability against convection is obtained by evaluating this expression at the largest temperature gradient permitted by the Ledoux criterion:

$$\nabla_{\text{max}} = \nabla_{\text{ad}} + \nabla_{\mu}. \quad (\text{B13})$$

In the radiative layer, we therefore have

$$\nabla_{\text{rad}} = \nabla_{\text{max}} = \frac{3\kappa PL}{64\pi\sigma GMT^4}, \quad (\text{B14})$$

leading to a definition for the maximum luminosity that can be transported without violating Ledoux stability:

$$L_{\text{Ledoux}} = \left(\frac{64\pi\sigma GMT^4}{3\kappa P} \right) \nabla_{\text{max}}, \quad (\text{B15})$$

evaluated at a representative location near the top of the Ledoux-stable region. The intrinsic luminosity is then capped such that $L_{\text{int}} = \min(L_{\text{int}}, L_{\text{Ledoux}})$, where L_{int} is the value in the absence of the stable boundary. This upper limit is applied iteratively during the solution procedure until a self-consistent thermal structure satisfying Ledoux stability is obtained. Any larger luminosity would require a radiative temperature gradient steeper than ∇_{max} , causing ∇_{rad} to exceed the Ledoux-stability limit. In that case the stabilizing molecular-weight gradient would no longer suppress convection, and the layer would become convectively unstable.

Expressing the radiative flux as a finite difference across the Ledoux-stable region, we have

$$L_{\text{Ledoux}} \simeq 4\pi r_{\nabla_{\mu,\text{max}}}^2 \frac{4\sigma}{3} \frac{T_{\text{lower}}^4 - T_{\text{upper}}^4}{\Delta\tau_{\text{Ledoux}}}, \quad (\text{B16})$$

which makes explicit that the transmissible luminosity is controlled by the integrated optical-depth thickness across the inhibited region $\Delta\tau_{\text{Ledoux}}$. Here T_{lower} and T_{upper} are the temperatures at the lower and upper boundaries of the layer, respectively, and $r_{\nabla_{\mu,\text{max}}}$ is a representative radius within the region where ∇_{μ} is maximized. Equation (B16) therefore represents the integrated radiative throughput of the Ledoux-stable layer and provides a global maximum for the intrinsic luminosity. It also avoids numerical oscillations that can arise from high-frequency layer-by-layer adjustments of the temperature gradient.

Molecular-weight gradients specify where stratification is stable against overturn. In the presence of a phase change, an additional threshold must obtain for stability due to the thermal effects of the enthalpy of reaction. Accordingly, we include a microphysical inhibition threshold (S. Markham et al. 2022), based on a critical “heavy species” mole-fraction threshold x_{inhib} :

$$x_{\text{inhib}} = \left[\left(\frac{\mu_{\text{melt}} \Delta H}{RT} - 1 \right) (\epsilon - 1) \right]^{-1}, \quad (\text{B17})$$

where μ_{melt} is the mean molecular weight of the condensed phase, $\epsilon \equiv \mu_{\text{vap}}/\mu_{\text{H}_2}$, where μ_{vap} is the mean molecular weight of the heavy condensed vapor species (e.g., Mg or SiO), and ΔH is the effective latent heat per unit mass for condensation. Physically, x_{inhib} encodes stable stratification in the presence of the phase change. In practice, the Ledoux-stable, inhibited region is activated only when both a stabilizing molecular-weight gradient exists ($\nabla_{\mu} > 0$) and the heavy-component abundance exceeds the compositional threshold ($x_{\text{heavy}} > x_{\text{inhib}}$). We combine these criteria into a smooth inhibition “strength” $w \in [0, 1]$ that controls how strongly the local temperature gradient is driven toward the radiative (diffusion-limited) gradient:

$$\begin{aligned} w &= w_{\mu} w_{x_{\text{inhib}}}, \\ w_{\mu} &= \mathcal{H}(\nabla_{\mu}), \\ w_{x_{\text{inhib}}} &= \mathcal{H}\left(\frac{x_{\text{heavy}}}{x_{\text{inhib}}}\right), \end{aligned} \quad (\text{B18})$$

where $\mathcal{H}$ is a Hill function used to smooth the transition into and out of the inhibited regime and to avoid numerical noise.

The temperature gradient is evaluated by interpolating between a baseline gradient and the radiative-required gradient:

$$\nabla_{\text{eff}} = (1 - w) \nabla + w \nabla_{\text{rad}}, \quad (\text{B19})$$

where $\nabla = \nabla_{\text{ad}}$ in convective regions and $\nabla = \nabla_{\text{rad}}$ otherwise. In the limit $w \rightarrow 1$, the temperature gradient approaches the diffusion-limited value, while for $w \rightarrow 0$ the solution reverts to the transport regime outside the Ledoux-stability region.

Both types of basal boundary layers are global bottlenecks for the intrinsic luminosity. When a Ledoux-stable compositional gradient is present, L_{int} is capped at the maximum luminosity that can be transmitted through the inhibited layer without violating Ledoux stability. This constraint is conceptually analogous to the luminosity limit imposed by the basal thermal boundary layer of thickness δ_{BL} , which bounds the net heat flux delivered from the magma ocean into the envelope. The distinction between the two lies in resolvability: the Ledoux-stable region is sufficiently extended that its stabilizing gradient can be described explicitly, whereas the basal thermal boundary layer is often too thin to resolve and is therefore parameterized as a global flux limit. In the cases of Uranus and Neptune, the Ledoux-stable region of convective inhibition is by far the dominant governor on intrinsic luminosity, with the thin boundary layer reducing the luminosity by factors of order 5 and the Ledoux region reducing the latter by factors of order 50.

Appendix C Physical Chemistry

Recent work on the physical chemistry of H_2 -silicate-Fe metal systems applied to sub-Neptunes (E. D. Young et al. 2024, 2025; J. G. Rogers et al. 2025; T. Gilmore & L. Stixrude 2026) shows that the interface separating outer envelopes and underlying magma oceans is a phase boundary that can be well approximated as the binodal (or solvus) surface in the MgSiO_3 - H_2 system (E. D. Young et al. 2024; J. G. Rogers et al. 2025; T. Gilmore & L. Stixrude 2026). This boundary moves inward toward the center as the planet cools (J. G. Rogers et al. 2025). Interior to the boundary, the magma ocean is a

supercritical mixture of silicate (modeled as MgSiO_3) and hydrogen with or without a distinct Fe-rich metal phase. For planets with sufficient hydrogen that their cores have greater than about 1% by mass H_2 , Fe metal, MgSiO_3 (silicate), and H_2 are entirely miscible (E. D. Young et al. 2025). In these cases, the cores are a single phase throughout except perhaps at the most shallow depths.

By determining the common tangents to the Gibbs free energy of mixing in the $\text{MgSiO}_3\text{--H}_2$ system, using the convex hull for the function for example, the values for the mole fractions of hydrogen, x_{H_2} , for melt and vapor in equilibrium are obtained. These coexisting compositions at various P and T values define the binodal surface that separates cores from overlying envelopes. The nonideal mixing between MgSiO_3 and H_2 calculated by T. Gilmore & L. Stixrude (2026) is used here. Their DFT-MD simulations yield the binary free energy of mixing expression

$$\begin{aligned} \Delta\hat{G}_{\text{binary mix}} = & (L_{\text{CB}}x_{\text{H}_2} + L_{\text{BC}}(1 - x_{\text{H}_2}))x_{\text{H}_2}(1 - x_{\text{H}_2}) \\ & \times \left(1 - \frac{T}{\tau} + \frac{P}{\pi}\right) \\ & + RT(x_{\text{H}_2} \ln x_{\text{H}_2} + (1 - x_{\text{H}_2}) \ln(1 - x_{\text{H}_2})), \end{aligned} \quad (\text{C1})$$

where L_{CB} and L_{BC} are the binary interaction parameters and x_{H_2} is the mole fraction of H_2 along the binary join. The values for the subregular mixing parameters that fit the DFT-MD simulations are $L_{\text{CB}} = 622,000 \text{ J mol}^{-1}$, $L_{\text{BC}} = -4950 \text{ J mol}^{-1}$, $\tau = 4800 \text{ K}$, and $\pi = -35 \text{ GPa}$. Addition of minor elements like Al has minimal effects on the position of the binodal surface in $P\text{--}T\text{--}x_{\text{H}_2}$ space (T. Gilmore & L. Stixrude 2026).

In order to determine the state of the core given the planet's pressure–temperature structure and its mass fraction of hydrogen we make use of the ternary phase equilibria that includes iron (E. D. Young et al. 2025). With this approach, we construct a thermodynamically consistent ternary mixing model for miscible $\text{MgSiO}_3\text{--Fe--H}_2$ melts by extrapolating three independently constrained binary systems ($\text{MgSiO}_3\text{--H}_2$, $\text{MgSiO}_3\text{--Fe}$, and Fe--H_2) into ternary composition space. Phase stability is determined by minimizing the molar Gibbs free energy of mixing at fixed pressure and temperature. The total free energy of mixing is written as the sum of an ideal configurational entropy term and a nonideal (excess) mixing term. The ideal contribution is given by $RT\sum_i x_i \ln x_i$, where x_i are the mole fractions of Fe (A), H_2 (B), and MgSiO_3 (C). Nonideal interactions are parameterized using a subregular solution formalism, with interaction parameters taken directly from published DFT-MD simulations and experimental constraints on the three binary joins.

The resulting Gibbs free energy of mixing for the ternary system is obtained following the Muggianu–Jacob projection method:

$$\begin{aligned} \Delta\hat{G}_{\text{mix}} = & RT(x_{\text{A}} \ln x_{\text{A}} + x_{\text{B}} \ln x_{\text{B}} + x_{\text{C}} \ln x_{\text{C}}) \\ & + (1/2)x_{\text{A}}x_{\text{B}}(L_{\text{AB}}(1 + x_{\text{B}} - x_{\text{A}}) \\ & + L_{\text{BA}}(1 + x_{\text{A}} - x_{\text{B}})) \\ & + (1/2)x_{\text{B}}x_{\text{C}}(L_{\text{BC}}(1 + x_{\text{C}} - x_{\text{B}}) \\ & + L_{\text{CB}}(1 + x_{\text{B}} - x_{\text{C}})) \\ & \times (1 - T/\tau + P/\pi) \\ & + (1/2)x_{\text{C}}x_{\text{A}}(L_{\text{CA}}(1 + x_{\text{A}} - x_{\text{C}}) \\ & + L_{\text{AC}}(1 + x_{\text{C}} - x_{\text{A}})), \end{aligned} \quad (\text{C2})$$

where x_{A} , x_{B} , and x_{C} denote the mole fractions of Fe, H_2 , and MgSiO_3 , respectively, with $x_{\text{A}} + x_{\text{B}} + x_{\text{C}} = 1$. The $\text{MgSiO}_3\text{--H}_2$ interaction parameters are $L_{\text{CB}} = 6.22 \times 10^5 \text{ J mol}^{-1}$ and $L_{\text{BC}} = -4.95 \times 10^3 \text{ J mol}^{-1}$, with temperature and pressure dependence governed by $\tau = 4800 \text{ K}$ and $\pi = -35 \text{ GPa}$ (T. Gilmore & L. Stixrude 2026). The $\text{MgSiO}_3\text{--Fe}$ interaction is treated as a regular solution with $L_{\text{AC}} = L_{\text{CA}} = 2.4 \times 10^5 - 28T + 1116P \text{ J mol}^{-1}$, based on DFT-MD calculations for MgO--Fe mixing (E. D. Young et al. 2025). The Fe--H_2 interaction parameters are $L_{\text{AB}} = 1.38 \times 10^5 - 9500P \text{ J mol}^{-1}$ and $L_{\text{BA}} = 1.7 \times 10^4 - 9500P \text{ J mol}^{-1}$, calibrated to reproduce experimental and ab initio constraints on hydrogen solubility in molten iron. No explicit, and unknown, ternary interaction parameter is included ($L_{\text{ABC}} = 0$), as is often the case in such applications.

With this free energy of mixing we can map the regions of miscibility and immiscibility between MgSiO_3 , Fe, and H_2 as a function of pressure and temperature.

Appendix D Material Properties

For liquid silicate densities, we use a previously derived approach (E. D. Young et al. 2025). Briefly, for the silicate melt we use an EoS fit to the MgSiO_3 liquid properties (N. de Koker & L. Stixrude 2009) using the algorithms of A. S. Wolf & D. J. Bower (2018). Total pressure is computed by combining the elastic (cold) compression term from a Vinet EoS with thermal pressure contributions (A. S. Wolf & D. J. Bower 2018). For liquid iron, we use a Vinet EoS modified to include thermal pressure from Y. Kuwayama et al. (2020). The Grüneisen parameter is a simple function of the compression ratio, which in turn determines the thermal contribution to pressure.

We find that addition of H (and by fiat, Fe) to the supercritical magma ocean changes its density. DFT-MD simulations show that the presence of hydrogen in the silicate melt reduces its density such that a linear mixture of the compressed densities of MgSiO_3 and H_2 by volume reproduces the density of the mixture at the conditions of the binodal (E. D. Young et al. 2025). We therefore include the effects of H_2 on the density of the supercritical melt by calculating the density of the mixture at the pressures and temperatures of the surface of the magma ocean, ρ_0 . The density of the mixture is

$$\frac{1}{\hat{V}_{\text{mix}}} = x_{\text{H}_2} \frac{1}{\hat{v}_{\text{H}_2}} + x_{\text{sil}} \frac{1}{\hat{v}_{\text{sil}}}, \quad (\text{D1})$$

where $\hat{v}_{\text{H}_2}$ and $\hat{v}_{\text{sil}}$ are the molar volumes, and where x_{H_2} and x_{sil} are the mole fractions of hydrogen and silicate in the binary mixture. Densities are obtained from $\text{MW}_{\text{mix}}/\hat{V}_{\text{mix}}$, where MW_{mix} is the molecular weight of the mixture. Molar volumes are obtained from densities and the EoSs using ρ_i/MW_i . We fixed the densities to be 0.09 and 2.5 g cm^{-3} for H_2 and silicate, respectively, at the surface of the magma ocean. These values come from the EoSs for hydrogen (G. Chabrier et al. 2019) and ab initio MD simulations of MgSiO_3 melt at 6000 K and 3.5 GPa (E. D. Young et al. 2025).

Deeper in the supercritical magma ocean, the effects of hydrogen on the compressibilities must be accounted for. Based on a complete EoS for the mixture of MgSiO_3 and 4% by weight H_2 (S. P. Marcum et al. 2026) we determined that

the reference bulk modulus, K_T , is similar for pure MgSiO_3 liquid and for the supercritical mixture containing 4% H_2 . However, the pressure derivative of the bulk modulus, K' , decreases from 5.71 for MgSiO_3 liquid to 4.32 for the H-bearing mixture (S. P. Marcum et al. 2026). The smaller value of K' indicates that the supercritical mixture stiffens less rapidly with pressure, so that added hydrogen both lowers the density and increases the compressibility of the phase at high pressure. For H-bearing compositions, we therefore approximate the pressure derivative of the bulk modulus by linear interpolation and extrapolation in H_2 mass fraction:

$$K'(w_{\text{H}_2}) = K'_{\text{pure}} + (K'_{4\% \text{ H}_2} - K'_{\text{pure}}) \left(\frac{w_{\text{H}_2}}{0.04} \right). \quad (\text{D2})$$

Here K'_{pure} denotes pure MgSiO_3 liquid, with a value of 7.42, $K'_{4\% \text{ H}_2}$ denotes the supercritical mixture with 4% H_2 by mass, with a value of 4.47, and w_{H_2} is the H_2 mass fraction. This approximation is tested against densities from additional DFT-MD simulations at conditions near to those obtained for the interiors of Neptune and Uranus using the methods described previously (E. D. Young et al. 2025; S. P. Marcum et al. 2026). We find that using Equation (D2), we reproduce the densities at similar temperatures and pressures, albeit with some differences in the weight fractions of hydrogen in the melt.

A comparison between experimental data for MgSiO_3 , the complete EoS for the mixture with 4% H_2 , and the extrapolation to higher concentrations of hydrogen up to 12% by mass is shown in Figure D1. The plot illustrates that our densities are broadly similar to those predicted by simple volume-additive mixing of densities for pure MgSiO_3 and pure H_2 using their respective EoSs, except at the lowest pressures (compare dotted and solid curves in Figure D1). In greater detail, however, one can also see that for the 10.5% H_2 adiabat, for example, similar to our Uranus model, the DFT-MD simulations produce densities similar to our extrapolated values but with 12% H_2 rather than 10.5% H_2 . Accounting for this discrepancy would require the same pressure–density curves to be calculated using greater total mass fractions of H_2 for our models. We observe that changing the parameters of the EoS has this effect: the same pressure–density curves fit the planets but with higher total mass fractions of hydrogen. The higher mass fractions of hydrogen are associated with lower temperatures along the adiabats. Until the full EoS of this supercritical melt phase as a function of H_2 concentration is firmly established, we conclude that going beyond the linear extrapolation of K' is neither warranted nor justified, and accept that there is a systematic uncertainty in the absolute magnitude of the total hydrogen mass fraction in our models.

We model Fe dissolved in a supercritical silicate– H_2 melt by assigning Fe a partial molar volume that reflects Fe occupation of Mg-like cation sites within the silicate melt framework (E. D. Young et al. 2025). The Fe-site partial molar volume is

$$\bar{V}_{\text{Fe site}}^{\text{eff}} = \phi \alpha \bar{V}_{\text{MgSiO}_3}^{\text{pure}}(P, T), \quad (\text{D3})$$

and the corresponding effective partial Fe density is

$$\rho_{\text{Fe site}}^{\text{eff}} = \frac{MW_{\text{Fe}}}{\bar{V}_{\text{Fe site}}^{\text{eff}}}, \quad (\text{D4})$$

where ϕ denotes the fraction of the pure MgSiO_3 formula-unit molar volume attributable to the Mg site, $\alpha = (r_{\text{Fe}}/r_{\text{Mg}})^3$ scales

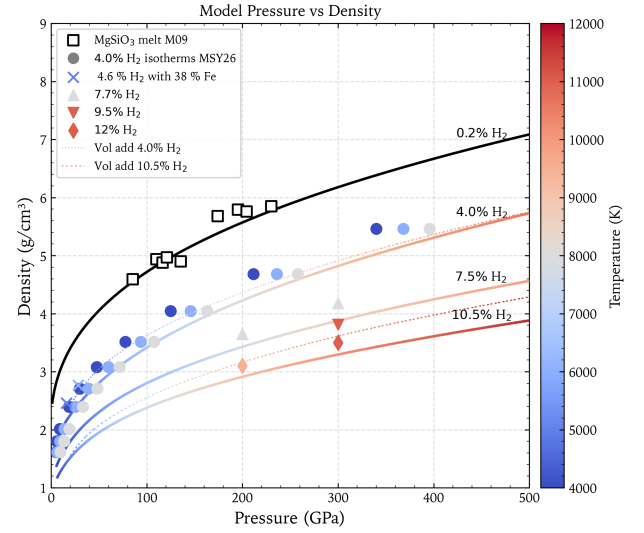


Figure D1. EoS and density data for MgSiO_3 and supercritical mixtures of MgSiO_3 and H_2 from DFT-MD simulations (circles and diamonds) compared with experiments (squares; J. L. Mosenfelder et al. 2009, denoted “M09”) and model adiabats using our EoS for this study. Curves for simple volume-additive densities are shown for comparison. Adiabats curves using the EoS used here are labeled with the H_2 concentrations in weight percent for the miscible silicate– H_2 mixtures. Corresponding volume-additive curves are shown for the 4% and 10.5% cases. Note adiabats are shallower than isotherms on this diagram.

the site volume by the ratio of octahedrally coordinated ionic radii (R. D. Shannon 1976), and $\bar{V}_{\text{MgSiO}_3}^{\text{pure}} = MW_{\text{MgSiO}_3} / \rho_{\text{MgSiO}_3}^{\text{pure}}$ is the molar volume of pure MgSiO_3 at the local P and T . Assuming additivity of specific volumes (no excess mixing volume), the density of the miscible melt with Fe mass fraction w_{Fe} is given by

$$\frac{1}{\rho_{\text{mix}}} = \frac{1 - w_{\text{Fe}}}{\rho_{\text{MgSiO}_3 + \text{H}_2}} + \frac{w_{\text{Fe}}}{\rho_{\text{Fe}}^{\text{eff}}}, \quad (\text{D5})$$

where $\rho_{\text{MgSiO}_3 + \text{H}_2}$ is the density of the $\text{MgSiO}_3 + \text{H}_2$ mixture derived at P and T . An empirical value for ϕ of 0.5 is obtained from the MgO –Fe system (L. Insixiengmay & L. Stixrude 2025). The effect of Fe on the supercritical mixture as determined using DFT-MD with methods similar to L. Insixiengmay & L. Stixrude (2025) is illustrated by the two crosses in Figure D1 where it is seen to be relatively minor.

The densities versus normalized planet radii for both Neptune and Uranus based on the EoSs described above are shown in Figure D2.

We estimate the depression of the MgSiO_3 solidus due to dissolved H_2 using the cryoscopic equation for an ideal solution:

$$\Delta T_m = - \frac{RT_m^2}{\Delta H_{\text{fus}}} \ln X_{\text{MgSiO}_3}, \quad (\text{D6})$$

where T_m is the pure-component melting temperature, ΔH_{fus} is the molar enthalpy of fusion, and X_{MgSiO_3} is the mole fraction of MgSiO_3 in the melt. At a binodal condition of $P_{\text{binodal}} \approx 5$ GPa and $T \approx 3300$ K, the model H_2 mass fraction in the silicate melt is $f_{\text{H}_2}^{\text{core}} \approx 3$ wt%, corresponding to $X_{\text{MgSiO}_3} = 0.38$. Taking $T_m = 3000$ K and $\Delta H_{\text{fus}} = 75$ kJ mol $^{-1}$ at 5 GPa (L. Stixrude & B. Karki 2005; L. Stixrude 2014), the ideal-mixing estimate

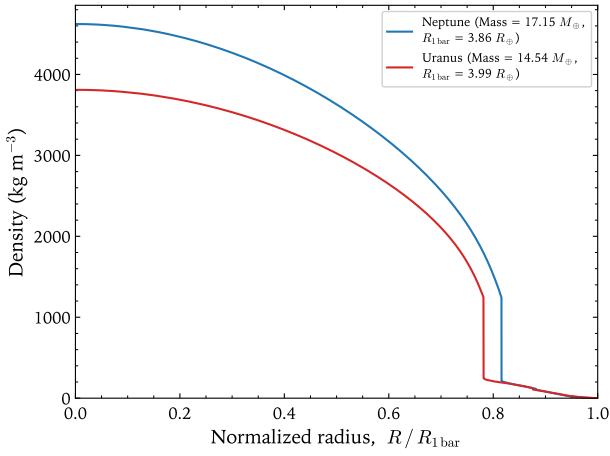


Figure D2. Density vs. normalized radial position for the Neptune and Uranus best-fit models.

yields $\Delta T_m \approx 965$ K, depressing the effective MgSiO_3 solidus from ~ 3000 to ~ 2035 K and placing the binodal conditions well within the melt field. The enthalpy of fusion carries substantial uncertainty (published ab initio values span $50\text{--}200$ kJ mol $^{-1}$), but even at the upper bound the depression exceeds several hundred degrees, making this conclusion reasonably robust.

The melting-point depression at fixed composition in Equation (D6) depends on pressure through $T_m(P)$ and $\Delta H_{\text{fus}}(P)$. The MgSiO_3 Clausius–Clapeyron slope is large and positive (L. Stixrude & B. Karki 2005), with T_m rising from ~ 3000 K at 5 GPa to $\sim 7000\text{--}8000$ K at 300 GPa. Since $\Delta T_m \propto T_m^2 / \Delta H_{\text{fus}}$, two competing effects operate with increasing pressure: T_m^2 grows rapidly, amplifying the depression, while ΔH_{fus} also increases moderately as the liquid–solid density contrast diminishes (L. Stixrude & B. Karki 2005). The T_m^2 term dominates, so the ratio $T_m^2 / \Delta H_{\text{fus}}$ grows with pressure and the cryoscopic depression is expected to increase throughout the interior. Taking $T_m \approx 7500$ K and $\Delta H_{\text{fus}} \approx 150$ kJ mol $^{-1}$ at 300 GPa as rough ab initio extrapolations, the ideal-mixing estimate yields $\Delta T_m \approx 3016$ K. This estimate is illustrative only given the uncertainties in ΔH_{fus} and H_2 solubility at these extreme conditions, but suggests that interiors modeled here should be entirely molten.

Iron mixing and nonideal $\text{H}_2\text{--MgSiO}_3$ mixing should further lower the melting point. For example, the activity coefficient $\gamma_{\text{MgSiO}_3} = 0.42$ from the regular-solution model used in our study (T. Gilmore & L. Stixrude 2026). Replacing X_{MgSiO_3} with the activity $\gamma_{\text{MgSiO}_3} X_{\text{MgSiO}_3} = 0.16$ in Equation (D6), the estimated depression in melting temperature is ~ 1830 K, indicating that nonideality acts in the direction of stabilizing the melt phase.

Appendix E Winds

Our approach in calculating the dynamic components of the gravity signature follows that outlined in J. Wicht et al. (2020) and W. Dietrich et al. (2021) and is adapted to the even gravity harmonics available for the ice giants as in D. Soyuer et al. (2023). We refer the reader to these works for further details and only outline the specifics of how the method was applied here.

The zonal gravity harmonic components due to the winds are defined as

$$J_n^{\text{dyn}} = -\frac{1}{MR^n} \int_V \rho^{\text{dyn}}(\mathbf{r}) P_n(\cos \theta) r^n dV, \quad (\text{E1})$$

having expanded both the density and gravity in terms of a static (background state) and dynamic (due to the winds) component:

$$\rho = \rho^{\text{stat}} + \rho^{\text{dyn}}, \quad J_n^{\text{obs}} = J_n^{\text{stat}} + J_n^{\text{dyn}}. \quad (\text{E2})$$

Thus, to obtain J_n^{dyn} , we must link the zonal flows with their associated density perturbations. For this, we follow exactly the steps detailed in W. Dietrich et al. (2021). This includes the dynamic self-gravity (J. Wicht et al. 2020) and omits the centrifugal potential (H. Cao & D. J. Stevenson 2017; E. Galanti et al. 2017).

We use the analytical forms of the surface wind profiles, as a function of the colatitude θ , which are the fits to the observed cloud-tracking measurements. For Neptune this is given by

$$U_N(\theta) = \sum_{i=1}^6 c_i \sin((2i-1)\theta), \quad (\text{E3})$$

from D. Soyuer et al. (2023), while for Uranus we have

$$U_U(\theta) = 170 \times (0.6 \sin \theta + \sin 3\theta) \text{ m s}^{-1}, \quad (\text{E4})$$

from H. B. Hammel et al. (2001). These are extrapolated downward, with the surface structures taken to be invariant with respect to the axis of rotation (see Figures E1(a) and (d)). We then assume a radially dependent decay profile of the winds, with the functional form

$$Q(r) = \frac{\tanh((r - (1-h))/\delta_h) + 1}{\tanh(h/\delta_h) + 1}, \quad (\text{E5})$$

where we fix the width of the hyperbolic tangent $\delta_h = 0.01$ and vary the depth h (see Figures E1(b) and (e)).

As pointed out by H. Cao & D. J. Stevenson (2017), to reliably constrain the wind structure with only even gravity harmonics, the contributions beyond J_6 are required. For the lower harmonics, small uncertainties in the background profiles match the amplitudes of the dynamic contributions.

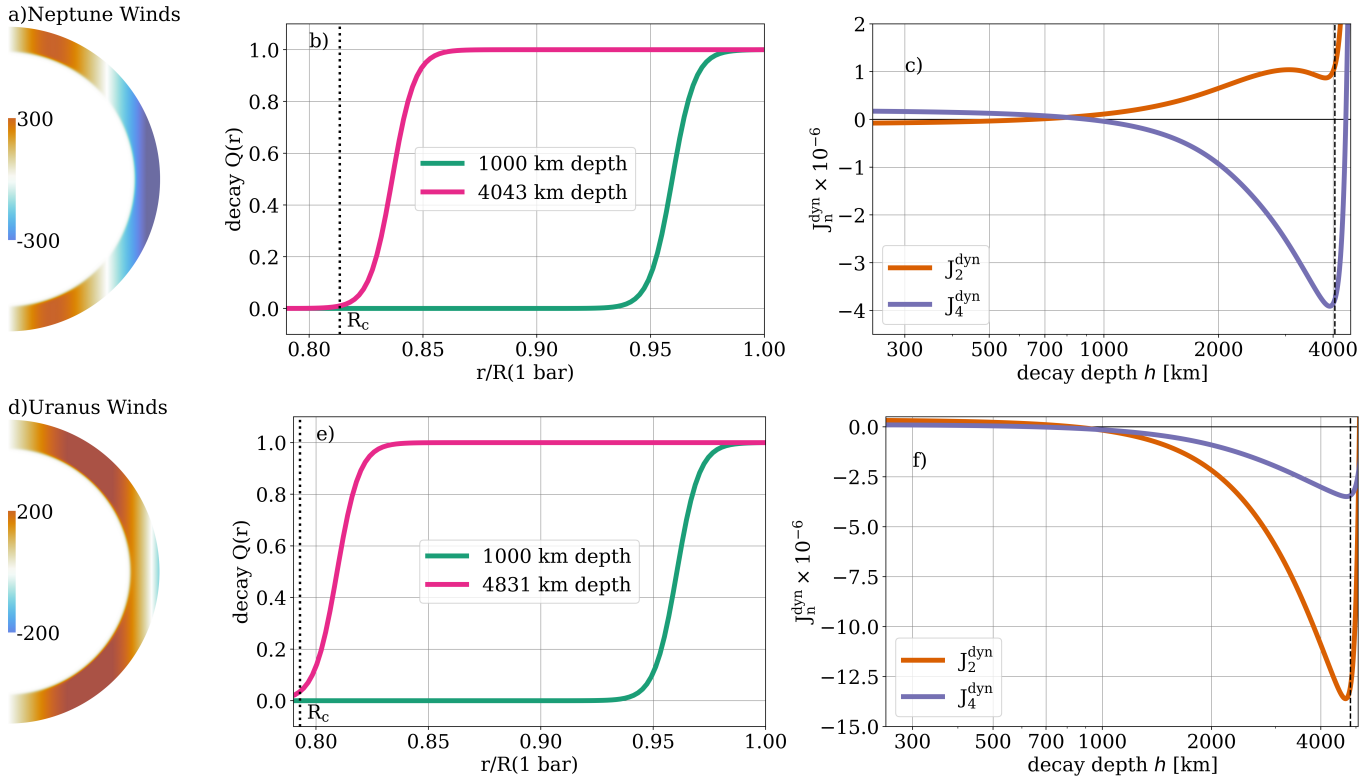


Figure E1. Plots illustrating wind models. Panels (a) and (d) show meridional cuts of the surface-measured zonal winds extrapolated downward—aligned with the axis of rotation—and reaching the bottom of the convective atmosphere of Neptune and Uranus, respectively. Red (blue) indicates prograde (retrograde) flow in meters per second. Panels (b) and (e) show the radial decay profile of the winds $Q(r)$, with radius on the x -axis normalized with respect to the 1 bar equatorial radius of the model. We show both our assumed, deep profiles, as well as profiles reaching 1000 km, similar to those inferred by Y. Kaspi et al. (2013). The dotted vertical line indicates the location of the core radius. Panels (c) and (f) are the values of J_2^{dyn} and J_4^{dyn} as a function of wind depth, where the dashed vertical line indicates 3838 and 4584 km depths, respectively.

Table E1

Dynamic Gravity Moments Based on Winds Extending to near the Bottom of Each Planet's Convective Atmosphere

Planet	Depth (km)	$J_2^{\text{dyn}} (\times 10^{-6})$	$J_4^{\text{dyn}} (\times 10^{-6})$
Neptune	3838	0.87	−3.88
Uranus	4584	−13.47	−3.46

Therefore, we do not take an optimization approach to the evaluation of the wind depths, and only iterate once between the calculation of J^{dyn} and adjustment of the static modeling. We show the calculated values of J_2^{dyn} and J_4^{dyn} for both planets as a function of zonal wind depth in Figures E1(c) and (f) (from model ρ^{stat} and g^{stat} as inputs).

For our reference values of J_2^{dyn} and J_4^{dyn} we take those corresponding to a wind depth based on the bottom of the convecting atmosphere in a previous iteration of the static models. These are given in Table E1.

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