

ON THE MINIMUM CORE MASS FOR GIANT PLANET FORMATION

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ABSTRACT

1. INTRODUCTION

Kelvin-Helmholz (KH) contraction considered previously (Ikoma, PapNel05). Say why useful. This paper develops a simplified model of KH contraction to both elucidate the relevant physical processes and explore a wide range of disk parameter space...

This paper is organized as follows. In section 2 we describe the assumptions of our atmosphere model, and derive the basic equations that govern the structure and evolution of the atmosphere. In section 3, we present a simplified analytic model that predicts the qualitative behavior of the numerical model. Results for atmospheric structure and evolution are presented in section 4, and implications for the critical core mass are presented in section 5. The discussion in section 6 addresses various approximations and neglected effects. We summarize our findings in section 7. The appendices contain...

2. ATMOSPHERE MODELS

To model the growth of planetary atmospheres around a solid core, we develop a simplified two layer model for time-dependent atmospheric cooling, i.e. Kelvin-Helmholtz (KH) contraction. With a convective interior and radiative exterior, this model is motivated by similar models of hot Jupiters (Arras & Bildsten 2006; Youdin & Mitchell 2010).

Our model can accurately model the growth of the atmosphere up to the crossover mass, when the atmosphere mass equals the mass of the core. Beyond the crossover mass, our approximate treatment of the radiative zone breaks down. Since subsequent growth is a rapid runaway process (*cite a numerical core accretion paper*), our model can investigate to good accuracy the timescale requirement of core accretion. Our simplified treatment is also inappropriate for hot, short-period planets, where dust sublimation gives deeper radiative zones that require more detailed models.

Our main assumptions are summarized as follows:

1. The atmosphere is spherically symmetric and remains in hydrostatic balance during its thermal evolution.
2. The core mass and radius are fixed in evolutionary calculations, neglecting ongoing planetesimal or dust accretion.
3. At the planet's Hill radius, the atmospheric temperature and pressure match the conditions of the disk midplane.

4. The only source of planetary luminosity is the gravitational contraction of the atmosphere.
5. In the radiative zone, luminosity generation is neglected, i.e. the luminosity is held constant.
6. A global cooling model connects independent static solutions into a time-dependent sequence.
7. A polytropic EOS is assumed for simplicity.
8. Dust grains provide the opacity in the radiative zone, which remains cool enough to avoid dust sublimation.
9. Because subsequent growth accelerates, the time to reach the crossover mass is a good approximation of the total time to form a gas giant.

The remainder of this section develops our model in more detail.

2.1. Disk and Opacity Model

We adopt a minimum mass solar nebula (MMSN) model for a passively irradiated disk (Chiang & Youdin 2010). With the semi-major axis a normalized to the outer disk as $a_{10\oplus} = a/(10 \text{ AU})$, the gas surface density and mid-plane temperature are

$$\Sigma_d = 70 F_\Sigma a_{10\oplus}^{-3/2} \text{ g cm}^{-2} \quad (1a)$$

$$T_d = 45 F_T a_{10\oplus}^{-3/7} \text{ K}. \quad (1b)$$

The normalization factors F_Σ and F_T adjust the model relative to the fiducial MMSN. We fix $F_\Sigma = F_T = 1$ unless noted otherwise.

For a vertically isothermal disk in hydrostatic balance (with no self-gravity), the mid-plane pressure of disk gas is

$$P_d = 6.9 \times 10^{-3} F_\Sigma \sqrt{F_T} a_{10\oplus}^{-45/14} \text{ dyne cm}^2 \quad (2)$$

for a molecular weight of $\mu = 2.35$ proton masses and a Solar mass star.

The (thermodynamically isothermal) sound speed in the disk is

$$c_d = \sqrt{\mathcal{R}T_d} = 0.4 \sqrt{F_T} a_{10\oplus}^{3/14} \text{ km s}^{-1} \quad (3)$$

in terms of the specific gas constant \mathcal{R} . The disk scale-height is

$$H_d = c_d/\Omega = 0.42 \sqrt{F_T} a_{10\oplus}^{9/7} \text{ AU}. \quad (4)$$

in terms of the Keplerian frequency $\Omega = \sqrt{GM_*/a^3}$ with G the gravitational constant and M_* the stellar (in this work Solar) mass.

We assume a dust opacity following Bell & Lin (1994):

$$\kappa = 2F_\kappa \left(\frac{T}{100 \text{ K}} \right)^\beta \text{ cm}^2 \text{ g}^{-1}, \quad (5)$$

with a powerlaw index $\beta = 2$ and normalization $F_\kappa = 1$ unless noted otherwise. Roughly speaking, grain growth will lower both F_κ and β , while dust abundance scales with F_κ . Section 6.2 discusses dust sublimation and more realistic opacity laws.

2.2. Length Scales

The characteristic length scales for protoplanetary atmospheres are crucial for choosing boundary conditions and for understanding the validity of spherical symmetry in a disk of scaleheight H_d . The radius of the solid core:

$$R_c \equiv \left(\frac{3M_c}{4\pi\rho_c} \right)^{1/3} \approx 10^{-4} m_{c10\oplus}^{1/3} \text{ AU}, \quad (6)$$

where the core mass, M_c , is normalized to 10 Earth masses as $m_{c10\oplus} \equiv M_c/(10 M_\oplus)$. The core density is held fixed at $\rho_c = 3.2 \text{ g cm}^{-3}$. We thus neglect the detailed equation of state of the solids core (Fortney et al. 2007).

A planet can bind a dense atmosphere if the escape speed exceeds the sound speed. This criterion is satisfied inside the Bondi radius

$$R_B \equiv \frac{GM_{\text{pl}}}{c_d^2} \approx 0.17 \frac{m_{c10\oplus} a_{10\oplus}^{3/7}}{F_T} \text{ AU} \quad (7)$$

where the planet mass, $M_{\text{pl}} = M_c + M_{\text{atm}}$, includes the mass of the core and atmosphere (defined as lying within the Bondi radius). The numerical estimates above and in Equation (8) assume the core mass dominates, $M_{\text{pl}} \sim M_c$.

Stellar tides dominate the planet's gravity beyond the Hill radius

$$R_H = \left(\frac{M_{\text{pl}}}{3M_*} \right)^{1/3} a \approx 0.22 m_{c10\oplus}^{1/3} a_{10\oplus} \text{ AU} \quad (8)$$

where hydrostatic balance clearly breaks down.

The relevant length scales of the atmosphere and disk satisfy the relation $R_B H_d^2 = 3R_H^3$. The lengthscales are roughly equal at the ‘‘thermal mass’’

$$M_{\text{th}} > \frac{c_d^3}{G\Omega} \approx 25 \frac{F_T^{3/2}}{\sqrt{m_*}} a_{10\oplus}^{6/7} M_\oplus. \quad (9)$$

In the low mass regime, $M_{\text{pl}} < M_{\text{th}}/\sqrt{3}$, the lengthscales order as $R_B < R_H < H_d$ (R06). In this regime, many studies assume the atmosphere matches the disk conditions at R_B . We however use R_H as the matching radius in both this low mass and other higher mass regimes. This choice is justified by the fact that, for hydrostatic solutions, the density at R_B exceeds the disk's background density by an order unity factor (R06). This modest density change has a similarly modest effect on our results.

For a finite range of intermediate masses, $M_{\text{th}}/\sqrt{3} < M_{\text{pl}} < 3M_{\text{th}}$, the Hill radius is the smallest scale, satisfying both $R_H < R_B$ and $R_H < H_d$. Spherical symmetry remains a good, if imperfect, approximation because the bound atmosphere within R_H doesn't see the vertical stratification of the disk on scales $\gtrsim H_d$.

At higher planet masses where $M_{\text{pl}} > 3M_{\text{th}}$ and $H_d < R_H < R_B$, spherical symmetry is no longer a good approximation, due to both the vertical stratification of the disk and gap opening. See §6.1 for discussion of neglected non-hydrostatic effects on all mass scales.

Note that while R_H is the outer boundary of our structure calculations, we define planet masses to include only the mass inside the smaller of R_B or R_H . This conservative choice in quoting planet masses is usually a minor distinction because (when $R_B < R_H$) the gas between R_B and R_H is weakly compressed.

2.3. Structure Equations and Boundary Conditions

Our atmosphere calculations use the standard structure equations of mass conservation, hydrostatic balance, thermal gradients, and energy conservation:

$$\frac{dm}{dr} = 4\pi r^2 \rho \quad (10a)$$

$$\frac{dP}{dr} = -\frac{Gm}{r^2} \rho \quad (10b)$$

$$\frac{dT}{dr} = \nabla \frac{T}{P} \frac{dP}{dr} \quad (10c)$$

$$\frac{dL}{dr} = 4\pi r^2 \rho \left(\epsilon - T \frac{\partial S}{\partial t} \Big|_m \right), \quad (10d)$$

where r is the radial coordinate, P , T , ρ and L are the gas pressure, temperature, density and luminosity, respectively. The enclosed mass at radius r is m . Equation (10c) simply defines the temperature gradient $\nabla \equiv d \ln T / d \ln P$. In radiative zones radiative diffusion gives a temperature gradient

$$\nabla_{\text{rad}} \equiv \frac{3\kappa P}{64\pi G m \sigma T^4} L, \quad (11)$$

where σ is the Stefan-Boltzmann constant. In convectively unstable regions, efficient convection gives an isentropic temperature gradient with $\nabla = \nabla_{\text{ad}}$, the adiabatic gradient. According to the Schwarzschild criterion, convective instability occurs when $\nabla_{\text{rad}} > \nabla_{\text{ad}}$. Thus $\nabla = \min(\nabla_{\text{rad}}, \nabla_{\text{ad}})$ set the temperature gradient.

In the energy equation (10d), ϵ represents all local sources of heat input, which excludes the motion of the atmosphere itself. In stars, nuclear burning contributes to ϵ . In a protoplanetary atmosphere, dissipative drag on planetesimals contributes to ϵ . Our simplified models set $\epsilon = 0$, consistent with our neglect of planetesimal accretion luminosity at the base of the atmosphere. The $\epsilon_g = -T \partial S / \partial t$ term gives the energy input from gravitational contraction.¹ The partial time derivative would normally require our radial derivatives should be partial derivatives. However our subsequent developments will

¹ In general, any motion in a non-stationary atmosphere is accounted for by this term. The partial time derivative is performed on shells of fixed mass.

replace the local energy equation (10d) with global energy balance, reverting the structure equations to time-independent ordinary differential equations (ODEs).

To solve the equation set (10) an equation of state (EOS) is required for closure. In our study, we adopt an ideal gas law with a polytropic EOS

$$P = \rho \mathcal{R} T, \quad (12a)$$

$$P = K \rho^\gamma, \quad (12b)$$

where K is the adiabatic constant. The adiabatic index $\gamma = 1/(1 - \nabla_{\text{ad}})$. An ideal monatomic gas has $\nabla_{\text{ad}} = 2/5$. This work uses $\nabla_{\text{ad}} = 2/7$ for an ideal diatomic gas. While our reference mean molecular weight ($\mu = 2.35$ proton masses) includes Helium, we ignore Helium's effect on the EOS, which is already greatly simplified. The second law of thermodynamics gives the relative entropy as

$$S = \mathcal{R} \ln \left(\frac{T^{1/\nabla_{\text{ad}}}}{P} \right) \quad (13)$$

eliminating the need for K .

Boundary conditions must be satisfied at both the base and the top of the atmosphere with $m(R_c) = M_c$, $T(R_H) = T_d$ and $P(R_H) = P_d$. In principle our solutions describe atmospheres with $L(R_c) = 0$. In practice, since we do not directly integrate Equation (10d) we need not directly impose this boundary condition, as described in §2.5.

2.4. Global Cooling of an Embedded Planet

This section describes the global energy balance of a planet embedded in a gas disk, or more generally any spherical, hydrostatic object in pressure equilibrium with a background medium. The total atmospheric energy includes gravitational and internal energies, $E = E_G + U$:

$$E_G = - \int_{M_c}^M \frac{Gm}{r} dm, \quad (14a)$$

$$U = \int_{M_c}^M u dm. \quad (14b)$$

The specific internal energy $u = C_V T = \mathcal{R}(\nabla_{\text{ad}}^{-1} - 1)T$ for a polytropic EOS. For a star or coreless planet, $L_c = M_c = 0$.

We start with the global energy balance for an isolated planet with a free surface:

$$L_M = L_c + \Gamma - \dot{E}. \quad (15)$$

The surface luminosity, L_M , includes contributions from the core luminosity L_c – e.g. planetesimal accretion or radioactive decay – from the total heat generation Γ – given by the integral of ϵ over the object – and from the rate of change of atmospheric energy \dot{E} , a loss term.

For an object with no core luminosity (or no core) and no internal heat sources, the energy equation $L_M = -\dot{E}$ describes KH contraction in its simplest form. For a main sequence star $L_M = \Gamma$ as the internal heat of nuclear burning supplies the total luminosity.

For a protoplanetary atmosphere embedded in a gas disk, the full energy equation,

$$L_M = L_c + \Gamma - \dot{E} + e_M \dot{M} - P_M \left. \frac{\partial V_M}{\partial t} \right|_M, \quad (16)$$

acquires surface terms as derived in Appendix A. The energy accreted across the surface is given by the specific energy $e_M = u_M - GM/R$ and the mass accretion rate \dot{M} . The work done by the surface is $P_M \partial V_M / \partial t$, with the partial derivative performed at fixed mass. This generalized energy equation applies on any spherical shell where hydrostatic balance holds.² Thus M no longer refers to a uniquely defined total mass, but to the chosen mass level, where the instantaneous radius is R . Values on this shell are labelled by M subscripts.

For static solutions, which are not the focus of this Paper, the surface terms (and also \dot{E}) vanish. Static solutions are valid when imposed heat sources, i.e. L_c and Γ exceed the atmospheric losses. Quantitatively, static solutions apply when the evolutionary timescale exceeds the KH timescale,

$$\tau_{\text{KH}} \sim \frac{|E|}{L_M}, \quad (17)$$

where surface terms are assumed to be subdominant. Thus evolutionary calculations, including the quasistatic calculations of this Paper, are needed to consider the fastest possible evolution that occurs on τ_{KH} .

2.5. The Two-Layer Model

To simplify our calculations of atmospheric contraction we use a two layer model with a bottom convective region and an upper radiative layer. The existence of such a structure is well known from previous studies (R06) and can be readily understood. Before the protoplanetary atmosphere can cool, it has the entropy of the disk. As the atmosphere cools the deep interior remains convective. Convective interiors are a common feature of low mass cool objects (brown dwarfs and planets) that results from the behavior of ∇_{rad} for realistic opacity laws. However the entropy of the deep interior decreases as the atmosphere cools. A region of outwardly increasing entropy, i.e. a radiative layer, is required to connect the convective interior to the disk. A more complicated structure, with radiative windows in the convection zone, is possible as discussed in §6.2.

In convective regions, the adiabatic structure is independent of luminosity and can be calculated without local energy balance, Equation (10d). Thus for fully convective objects, a cooling sequence can be established by connecting a series of adiabatic solutions using a global energy equation, $L_M = -\dot{E}$ or Equation (15). Such methods are commonly used for their computational efficiency and are sometime referred to as “following the adiabats,” since the steady state solutions evolve in order of decreasing entropy (Marleau & Cumming 2013).

In the radiative zone, local energy balance, Equation (10d), does affect the atmospheric structure. We proceed by assuming that the majority of energy is lost from the convective interior, and thus the luminosity can be treated as constant in the outer radiative zone. With this approximation we can construct solutions from equa-

² Equation (16) also applies in the interior of objects with a free surface. The work term, $P_M \partial V_M / \partial t$, vanishes at the free surface. The accretion energy, $e_M \dot{M}$, vanishes for a truly isolated object, but would be included to account for accretion onto (or mass loss from) an otherwise free surface.

tions (10a – c) that “follow the mass,” i.e. gradually increase the atmospheric mass. We then use the global energy balance, Equation (16) to connect these solution in a cooling sequence. The validity of neglecting luminosity generation in the radiative zone can be checked *a posteriori*.

To make a single atmosphere model (indexed by i) we choose a planet mass M_i . At the outer boundary, at $R_H(M_i)$, the temperature and pressure are set to the disk values. To integrate equations (10a–c), the luminosity is required to compute ∇_{rad} . The correct value of the luminosity is the eigenvalue of the problem, which we find by the shooting method. Only for the eigenvalue of luminosity does the integrated value of mass at the core, $m(R_c)$, match the actual core mass, M_c .

To establish the time difference between neighboring solutions, we apply Equation (16) at the radiative-convective boundary (RCB) of the solutions. Using exact solutions, energy balance could be evaluated at any level. The approximation of constant luminosity in the radiative zone makes the RCB the self-consistent choice. After setting $\Gamma = L_c = 0$, we solve for the elapsed time Δt between states i and $i+1$ by finite differencing to give

$$\Delta t = \frac{-\Delta E + \langle e \rangle \Delta M - \langle P \rangle \Delta V_{\langle M \rangle}}{\langle L \rangle}. \quad (18)$$

Brackets indicate an average of, and Delta (Δ) indicates a difference between, neighboring solutions. All values are evaluated at the RCB. Due to the partial derivative in Equation (16), the volume difference $\Delta V_{\langle M \rangle}$ is performed at fixed mass, here the average of the masses at the RCB.

3. ANALYTIC COOLING MODEL

This section develops the analytic version of our two-layer model, which is useful for predicting and explaining the results of numerical evolution models. To obtain our analytic results, we neglect self-gravity. As shown in section 4, self-gravity becomes quantitatively important at surprisingly low atmospheric masses, $M_{\text{atm}} \gtrsim 0.2M_c$. Thus the analytic model cannot accurately model the later stages of evolution. Nevertheless, the analytic model does successfully explain general trends with e.g. opacity and disk temperature and pressure.

The analytic model also assumes that the upper radiative radiative layer is thick, with a pressure depth $P_{\text{RCB}} \gg P_d$. This approximation limits the ability to study the initial cooling of very low mass atmospheres. However since early evolution is quite rapid, this approximation negligibly affects total cooling times and critical core masses, our main foci.

3.1. Two Layer Structure

To apply the two layer cooling model analytically, we require expressions for the atmospheric structure. Conditions at the RCB are crucial as they set the interior adiabat and the radiative losses. (Recall that luminosity generation in the radiative zone is neglected.) We express the temperature and pressure of the RCB, at the radius R_{RCB} , as

$$T_{\text{RCB}} = \chi T_d \quad (19a)$$

$$P_{\text{RCB}} = \theta P_d e^{R_B/R_{\text{RCB}}}. \quad (19b)$$

If the radiative zone were replaced by an isothermal layer, then hydrostatic balance would give $\chi = \theta = 1$. Standard radiative structure calculations (see Appendix B for details) give

$$\chi \simeq \left(1 - \frac{\nabla_{\text{ad}}}{\nabla_{\infty}}\right)^{-\frac{1}{4-\beta}} \simeq 1.53, \quad (20)$$

for $P_{\text{RCB}} \gg P_d$, our regime of interest, and with the radiative temperature gradient at infinite depth, $\nabla_{\infty} = 1/2$, for our dust opacity. The radiative zones are thus nearly isothermal, as also found by R06. While a simple closed form expression is not possible, the factor $\theta \simeq 0.556$ for our parameters. According to Equation (19b) our pressure criterion also gives $R_{\text{RCB}} \ll R_B$, but the inequality is logarithmically weaker.

Given the conditions at the RCB, hydrostatic balance gives the density and temperature profiles along the adiabat,

$$\rho = \rho_{\text{RCB}} \left[1 + \frac{R'_B}{r} - \frac{R'_B}{R_{\text{RCB}}}\right]^{1/(\gamma-1)} \quad (21a)$$

$$T = T_{\text{RCB}} \left[1 + \frac{R'_B}{r} - \frac{R'_B}{R_{\text{RCB}}}\right]. \quad (21b)$$

The effective Bondi radius

$$R'_B \equiv \frac{GM_c}{C_P T_{\text{RCB}}} = \frac{\nabla_{\text{ad}}}{\chi} R_B \quad (22)$$

simplifies expressions, with $C_P = \mathcal{R}/\nabla_{\text{ad}}$ the specific heat at constant pressure.

Even simpler expressions hold deep in the adiabatic interior, where $r \ll R_{\text{RCB}} \ll R'_B$ gives

$$\rho \simeq \rho_{\text{RCB}} \left[\frac{R'_B}{r}\right]^{1/(\gamma-1)} \propto r^{-5/2} \quad (23a)$$

$$T \simeq T_{\text{RCB}} \frac{R'_B}{r} = \frac{GM_c}{C_P r}. \quad (23b)$$

While the radial density profile depends on the adiabatic index, the temperature scaling does not. The total specific energy at depth

$$e = e_g + u = -\nabla_{\text{ad}} \frac{GM_c}{r}. \quad (24)$$

is thus generically proportional to the gravitational potential, $e_G = -GM_c/r$.

3.2. Mass, Energy and Luminosity

The most relevant quantities for the global cooling model are the integrated energy, luminosity, and atmospheric mass. In our non-self-gravitating limit, the mass of our nearly isothermal radiative zones is less than the convective interior, as shown in Appendix B.

Equation (23a) thus integrates to give the atmospheric mass

$$M_{\text{atm}} = \frac{5\pi^2}{4} \rho_{\text{RCB}} R_B'^{5/2} \sqrt{R_{\text{RCB}}}, \quad (25)$$

in the relevant limit $R_c \ll R_{\text{RCB}} \ll R'_B$ and for $\gamma = 7/5$. Mass is concentrated near the outer regions of the convective zone, a result that holds more generally for $\gamma > 4/3$.

Using Equation (19b) to eliminate R_{RCB} , the critical ratio of atmosphere to core mass becomes

$$\frac{M_{\text{atm}}}{M_c} = \frac{P_{\text{RCB}}}{\xi P_M}, \quad (26)$$

where we define a characteristic pressure and a logarithmic factor

$$P_M \equiv \frac{4\nabla_{\text{ad}}^{3/2} GM_c^2}{5\pi^2 \sqrt{\chi} R_B'^4} \quad (27a)$$

$$\xi \equiv \sqrt{\ln[P_{\text{RCB}}/(\theta P_d)]}. \quad (27b)$$

The atmosphere mass increases as radiative losses lower the internal adiabat and increase P_{RCB} . The crossover mass, $M_{\text{atm}} = M_c$, is reached when

$$P_{\text{RCB}} = \xi P_M, \quad (28)$$

i.e. near the characteristic pressure P_M . The critical value of the order unity factor ξ is found by eliminating P_{RCB} from Equations (27b) and (28). This logarithmic factor complicates our analytic description, but since it remains order unity we simply hold it fixed in our scalings.

The total energy is concentrated towards the core if $u\rho r^3 \propto \rho r^2$ drops with increasing r . This condition requires $\gamma < 3/2$, which our choice of $\gamma = 7/5$ satisfies, but $\gamma = 5/3$ does not. Integration of Equation (24) over the mass of the atmosphere thus gives

$$E = -4\pi \nabla_{\text{ad}} GM_c \int_{R_c}^{R_{\text{RCB}}} \rho r dr \quad (29a)$$

$$\approx -4\pi P_{\text{RCB}} R_B' \frac{1}{\nabla_{\text{ad}}} \left(\frac{\gamma - 1}{3 - 2\gamma} \right) R_c^{\frac{2\gamma-3}{\gamma-1}} \quad (29b)$$

$$\approx -8\pi P_{\text{RCB}} \frac{R_B'^{7/2}}{\sqrt{R_c}} \quad (29c)$$

with $\gamma < 3/2$ and $\gamma = 7/5$ in Equations (29b) and (29c), respectively.

Applying marginal convective stability, $\nabla_{\text{rad}} = \nabla_{\text{ad}}$, to Equation (11) gives the emergent luminosity from the radiative-convective boundary,

$$L_{\text{RCB}} = \frac{64\pi GM_{\text{RCB}} \sigma T_{\text{RCB}}^4}{3\kappa P_{\text{RCB}}} \nabla_{\text{ad}} \approx L_d \frac{P_d}{P_{\text{RCB}}}, \quad (30)$$

where we define

$$L_d \equiv \frac{64\pi GM_{\text{RCB}} \sigma T_d^4}{3\kappa(T_d) P_d} \nabla_{\text{ad}} \chi^{4-\beta}. \quad (31)$$

The inverse scaling $L_{\text{RCB}} \propto 1/P_{\text{RCB}}$ shows that luminosity drops as the atmosphere cools and P_{RCB} deepens. This result relies on the pressure independence of dust opacities. For fully non-self-gravitating results, we replace M_{RCB} , the mass up to the RCB, with the core mass, but the mass of the convective atmosphere could be included for a slightly higher order estimate.

3.3. Simplified Cooling Model

Our analytic cooling model uses $L = -\dot{E}$, neglecting the surface terms in Equation (16).³ Applying Equations

³ Appendix B shows that these terms are negligible for a non-self-gravitating model, but they do become relevant when self-gravity

(29c) and (30) the time it takes to cool the atmosphere until the RCB reaches a given pressure depth, P_{RCB} , is

$$t_{\text{cool}} = - \int_{P_d}^{P_{\text{RCB}}} \frac{dE/dP_{\text{RCB}}}{L} dP_{\text{RCB}} \quad (32a)$$

$$\approx 4\pi \frac{P_{\text{RCB}}^2}{P_d} \frac{R_B'^{7/2}}{L_d \sqrt{R_c}}. \quad (32b)$$

The initial RCB depth, set to P_d as a formality, is of little importance because the cooling slows it proceeds with $t_{\text{cool}} \propto P_{\text{RCB}}^2$.

The most relevant cooling time is the crossover time, t_{co} , required to reach the crossover mass, $M_{\text{atm}} = M_c$. Equations (32) and (28) give this time as

$$t_{\text{co}} \approx 2 \times 10^8 \frac{F_T^{5/2} F_\kappa \left(\frac{\xi}{3.4} \right)^2}{m_{c10\oplus}^{5/3} a_{10\oplus}^{15/14}} \text{ yr}. \quad (33)$$

The corresponding critical core mass, M_{crit} , has a crossover time that equals the typical lifetime of a protoplanetary disk

$$t_d = 3 \times 10^6 \text{ yr}. \quad (34)$$

Setting $t_{\text{co}} = t_d$ gives

$$M_{\text{crit}} \approx 100 \frac{F_T^{3/2} F_\kappa^{3/5} \left(\frac{\xi}{2.6} \right)^{6/5}}{a_{10\oplus}^{9/14}} M_\oplus. \quad (35)$$

Both t_{co} and M_{crit} are too large to be interesting, or to be correct based on previous results and the numerical results in this Paper. The main reason for this discrepancy is the neglect of self-gravity, as shown in section 4. The magnitude of the discrepancy may be surprising as one might expect self-gravity to be a modest factor for $M_{\text{atm}} \leq M_c$. Clearly this seemingly reasonable expectation must be applied with caution, an interesting result in itself.

The analytic model is also useful because, despite the amplitude error, it explains the basic scaling of our (and other) numerical models. As is well known, a lower opacity, here scaling with F_κ , accelerates growth and gives a lower critical core mass. Lower disk temperatures also decrease both t_{co} and M_{crit} . Furthermore the decline in both quantities with semimajor axis is completely generated by the disk temperature law (ignoring ξ). The temperature dependence are a competition between two main effects. The higher luminosity, $\propto T^{4-\beta}$, at higher temperatures gives faster cooling, which opposes the overall effect. The larger Bondi radius, and the effectively stronger gravity of the core, is the dominant effect that gives faster cooling at lower temperatures. For different β values, $t_{\text{co}} \propto F_T^{\beta+1/2}$, so the slope of the dust opacity is an important moderating effect on the cooling.

The cooling timescale and critical core mass depend only weakly on the disk pressure, via the logarithmic factor ξ . Section 4.0.1 shows that these analytic predictions for disk temperature and pressure dependence compare favorably with numerical results.

is important and included.

TODO: consider adding "revised" analytic expressions for t_{co} and M_{crit} (e.g. lower Matm criterion). People like scalings, but the obviously wrong amplitudes will prevent people from using them unless we correct.

4. QUASI-STATIC KELVIN-HELMHOLTZ CONTRACTION

4.0.1. Disk Temperature and Pressure

5. CRITICAL CORE MASS

6. NEGLECTED EFFECTS

6.1. Hydrodynamic Effects

The neglect of hydrodynamical effects in our model is best discussed in terms of the thermal mass, M_{th} , and the lengthscales introduced in §2.2. In the low mass regime, $M_{\text{pl}} < M_{\text{th}}/\sqrt{3}$, where $R_{\text{B}} < R_{\text{H}}$ we assume that hydrostatic balance holds out to the outer boundary at R_{H} . In this low mass regime, Ormel (2013) calculated the flow patterns driven by stellar tides and disk headwinds. On scales $\gtrsim R_{\text{B}}$ the flows no longer circulate the planet: they belong to the disk. Nevertheless, the density structure outside R_{B} remains spherically symmetric and hydrostatic. Even if these flows do not destroy hydrostatic balance, they could affect the planet's cooling. We expect such effects to be weak, as heat losses at greater depths dominate planetary cooling, but more study is needed.

At higher masses, non-hydrostatic effects become more severe. At $M_{\text{pl}} \gtrsim M_{\text{th}}$ planets can open significant gaps (Zhu et al. 2013). At yet higher masses accretion insta-

bilities could occur (Ayliffe & Bate 2012). However in this high mass regime we already know that the spherically symmetric approximation breaks down, as does our approximate cooling model (*discuss elsewhere and cite*).

Thus by restricting our attention to low masses, neglected hydrodynamic effects should be minor. Moreover since $M_{\text{th}} \propto a^{6/7}$ increases with disk radius, spherical hydrostatic models like ours have a greater range of applicability in the outer regions of disks.

6.2. Realistic Opacities and EOS

some comments on sublimation and radiative windows here (current text rough). upshot is the outer disk is good.

Real dust opacities exhibit a more complicated behavior that depends on grain composition. Opacities drop by order unity when ice grains sublimate for $T \gtrsim 150$ K and they drop by orders of magnitude when silicate grains evaporate above $T \gtrsim 1500$ K (Semenov et al. 2003; Ferguson et al. 2005). While protoplanetary atmospheres get significantly hotter than 1500 K, our model depends on opacity only in the upper radiative zones, which are cool enough to remain dusty. For our opacity approximations to be valid, we must restrict our model to the lower temperatures of the outer disk.

A second radiative zone inside the convection zone, a "radiative window," is possible. Dust sublimation at shallow depths – giving a low κ and P in Equation (11) – could cause such a window. Our two layer model ignores such complications.

7. CONCLUSIONS

REFERENCES

- Arras, P., & Bildsten, L. 2006, ApJ, 650, 394
 Ayliffe, B. A., & Bate, M. R. 2012, MNRAS, 427, 2597
 Bell, K. R., & Lin, D. N. C. 1994, ApJ, 427, 987
 Chiang, E., & Youdin, A. N. 2010, Annual Review of Earth and Planetary Sciences, 38, 493
 Ferguson, J. W., Alexander, D. R., Allard, F., Barman, T., Bodnarik, J. G., Hauschildt, P. H., Heffner-Wong, A., & Tamanai, A. 2005, ApJ, 623, 585
 Fortney, J. J., Marley, M. S., & Barnes, J. W. 2007, ApJ, 659, 1661
 Kippenhahn, R., & Weigert, A. 1990, Stellar Structure and Evolution
 Marleau, G.-D., & Cumming, A. 2013, arXiv:1302.1517
 Ormel, C. W. 2013, MNRAS, 428, 3526
 Semenov, D., Henning, T., Helling, C., Ilgner, M., & Sedlmayr, E. 2003, A&A, 410, 611
 Youdin, A. N., & Mitchell, J. L. 2010, ApJ, 721, 1113
 Zhu, Z., Stone, J. M., & Rafikov, R. R. 2013, ApJ, 768, 143

APPENDIX

A. DERIVATION OF THE GLOBAL ENERGY EQUATION

To derive the global energy equation (16) for an embedded protoplanets, we generalize the analogous calculations in stellar structure theory, e.g. in §4.3 of Kippenhahn & Weigert (1990). For our problem we add the effects of finite core radius, surface pressure and mass accretion. We start with the local energy equation (10d) whose more natural form in Lagrangian (mass) coordinates is $\partial L / \partial m = \epsilon - T \partial S / \partial t$. Integrating from the core to a higher shell with enclosed mass M gives:

$$L - L_{\text{c}} = \int_{M_{\text{c}}}^M \frac{\partial L}{\partial m} dm \quad (\text{A1a})$$

$$= \int_{M_{\text{c}}}^M \left(\epsilon - T \frac{\partial S}{\partial t} \right) dm \quad (\text{A1b})$$

$$= \Gamma - \int_{M_{\text{c}}}^M \frac{\partial u}{\partial t} dm + \int_{M_{\text{c}}}^M \frac{P}{\rho^2} \frac{\partial \rho}{\partial t} dm. \quad (\text{A1c})$$

with $\Gamma = \int \epsilon dm$ the integral of the direct heating rate and applying the law of thermodynamics in the final step.

The global energy equation is derived by eliminating the partial time derivatives in Equation (A1c), which are performed at a fixed mass, in favor of total time derivatives, denoted with overdots. For instance the surface radius, R , evolves as

$$\dot{R} = \frac{\partial R}{\partial t} + \frac{\dot{M}}{4\pi R^2 \rho_M} \quad (\text{A2})$$

where $\partial R/\partial t$ gives the Lagrangian contraction of the “original” shell, and mass accretion through the upper boundary at rate \dot{M} also changes the shell location. Similarly the volume, $V = (4\pi/3)r^3$, and pressure at the outer shell evolve as

$$\dot{V}_M = \frac{\partial V_M}{\partial t} + \frac{\dot{M}}{\rho_M} \quad (\text{A3a})$$

$$\dot{P}_M = \frac{\partial P_M}{\partial t} + \frac{\partial P_M}{\partial m} \dot{M} = \frac{\partial P_M}{\partial t} - \frac{GM}{4\pi R^4} \dot{M}. \quad (\text{A3b})$$

This derivation holds the core mass and radius fixed, $\dot{M}_c = \dot{R}_c = 0$. Therefore the core pressure satisfies

$$\dot{P}_c = \partial P_c / \partial t. \quad (\text{A4})$$

The internal energy integral follows simply from Leibniz’s rule as

$$\int_{M_c}^{M(t)} \frac{\partial u}{\partial t} dm = \dot{U} - \dot{M} u_M. \quad (\text{A5})$$

To make further progress we require the virial theorem:

$$E_G = -3 \int_{M_c}^M \frac{P}{\rho} dm + 4\pi(R^3 P_M - R_c^3 P_c) \quad (\text{A6})$$

which follows from Equations (10a), (10b) and (14) by integrating hydrostatic balance in Lagrangian coordinates. As an aside, the integral in equation (A6) can be evaluated for a polytropic EOS to give simple expressions for the total energy:

$$E = (1 - \zeta)U + 4\pi(R^3 P_M - R_c^3 P_c) \quad (\text{A7a})$$

$$= \frac{\zeta - 1}{\zeta} E_G + \frac{4\pi}{\zeta} (R^3 P_M - R_c^3 P_c) \quad (\text{A7b})$$

where $\zeta \equiv 3(\gamma - 1)$. We will not make this assumption and will keep the EOS general.

To express the work integral, the final term in Equation (A1c), in terms of changes to gravitational energy we first take the time derivative of Equation (A6):

$$\dot{E}_G = 3 \int_{M_c}^M \frac{P}{\rho^2} \frac{\partial \rho}{\partial t} dm - 3 \int_{M_c}^M \frac{\partial P}{\partial t} \frac{dm}{\rho} - 3 \frac{P_M}{\rho_M} \dot{M} + 3 \dot{P}_M V_M - 3 \dot{P}_c V_c + 3 P_M \dot{V}_M. \quad (\text{A8})$$

The first integral in Equation (A8) is the one we want, but the next one must be eliminated. The time derivative of Equation (14) (times four) gives

$$4\dot{E}_G = -4 \frac{GM\dot{M}}{R} + 4 \int_{M_c}^M \frac{Gm}{r^2} \frac{\partial r}{\partial t} dm \quad (\text{A9a})$$

$$= -4 \frac{GM\dot{M}}{R} + 4\pi \int_{M_c}^M r^3 \frac{\partial}{\partial m} \frac{\partial P}{\partial t} dm \quad (\text{A9b})$$

$$= -4 \frac{GM\dot{M}}{R} - 3 \int_{M_c}^M \frac{\partial P}{\partial t} \frac{dm}{\rho} + 3 V_M \frac{\partial P_M}{\partial t} - 3 V_c \frac{\partial P_c}{\partial t} \quad (\text{A9c})$$

where Equations (A9b) and (A9c) use hydrostatic balance and integration by parts.

Subtracting Equations (A5) and (A9c) and rearranging terms with the help of Equations (A2), (A3) and (A4) gives

$$\int_{M_c}^M \frac{P}{\rho^2} \frac{\partial \rho}{\partial t} dm = -\dot{E}_G - \frac{GM\dot{M}}{R} - P_M \frac{\partial V_M}{\partial t}. \quad (\text{A10})$$

Combining Equations (A1c), (A5) and (A10), we reproduce Equation (16) with the accreted specific energy $e_M \equiv u_M - GM/R$.

B. ANALYTIC COOLING MODEL DETAILS

TODO (academic): Figure out if theta $\neq 1$ is actually required? Should be but need to check