**Ejection of Chondrules from Fluffy Matrices**

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**Abstract**

Chondritic meteorites primarily contain millimeter-sized spherical objects, chondrules; however, the co-accretion process of chondrules and matrix grains is not yet understood. In this study, we investigate the ejection process of chondrules via collisions of fluffy aggregates composed of chondrules and matrices. We reveal that fluffy aggregates cannot grow into planetesimals without losing chondrules if we assume that the chondrite parent bodies are formed via direct aggregation of similar-sized aggregates. Therefore, an examination of other growth pathways is necessary to explain the formation of rocky planetesimals in our solar system.

**Key words:** meteorites, meteors, meteoroids – planets and satellites: terrestrial planets

1. Introduction

Terrestrial planets are believed to form in protoplanetary disks via collisions and the coalescence of rocky planetesimals; however, the process by which interstellar dust grains evolve into planetesimals is still enigmatic. This is because there are several “barriers” to planetesimal formation, e.g., the bouncing barrier (e.g., Zsom et al. 2010), the fragmentation barrier (e.g., Blum & Wurm 2008), and the radial drift barrier (e.g., Weidenschilling 1977).

The classical proposed mechanism for planetesimal formation is the gravitational instability of thin dust layer (e.g., Goldreich & Ward 1973; Sekiya 1983); however, it is now believed that the Kelvin–Helmholtz instability generated by vertical shears and the differential motions of dust particles and gas prevents disks from entering a gravitationally unstable state (e.g., Cuzzi et al. 1993). Subsequently, several alternative scenarios have been proposed to explain planetesimal formation, for example, the streaming instability (e.g., Youdin & Goodman 2005; Johansen & Youdin 2007) and collisional growth via direct aggregation (e.g., Windmark et al. 2012; Kataoka et al. 2013b).

The streaming instability is driven by differences in the motions of the gas and dust particles in the disk, and this mechanism requires a large dust-to-gas ratio (e.g., Drążkowska & Dullemond 2014). In addition, the streaming instability requires grown dust aggregates whose motions are marginally decoupled from the disk gas (e.g., Bai & Stone 2010). Therefore, whether planetesimals can form via the streaming instability depends on the vertically integrated dust-to-gas mass ratio and the critical velocity for collisional growth.

The direct aggregation scenario is also restricted by the solid mass fraction and the critical velocity for growth. At least for aggregates composed of submicron ice grains, the critical velocity could exceed the maximum collisional velocity of dust aggregates in protoplanetary disks (e.g., Wada et al. 2009, 2013). In addition, if the fluffy growth of dust aggregates is taken into consideration, Okuzumi et al. (2012) have shown that these aggregates can avoid the radial drift barrier even if they grow in the minimum mass solar nebula (Hayashi 1981). Moreover, according to numerical simulations, fluffy aggregates are less likely to bounce when they collide (Wada et al. 2011; Seizinger & Kley 2013).

By contrast, the critical collision velocity of rocky aggregates is still open to argument. While the critical velocity strongly depends on the surface energy and the radius of the monomers (Dominik & Tielens 1997), both the surface energy (e.g., Yamamoto et al. 2014; Kimura et al. 2015) and the monomer radius (Arakawa & Nakamoto 2016b) are still under debate. Additional mechanisms for increasing the critical velocity by dissipating the kinetic energy when dust aggregates collide have also been suggested by Tanaka et al. (2012) and Krijt et al. (2013). Furthermore, the critical velocity increases if the silicate monomers are coated with sticky organic mantles (e.g., Piani et al. 2017) or if the collisions of the rocky aggregates occur in high-temperature environments (e.g., Hubbard 2017).

If we discuss the formation process of rocky planetesimals in our solar system, we must consider how the chondrite parent bodies formed. Chondrites are one of the major groups of meteorites, and it has been shown that S-type asteroids are the parent bodies of ordinary chondrites (Nakamura et al. 2011). Ordinary chondrites primarily consist of millimeter-sized spherical objects so-called chondrules and nanometer- and micron-sized matrix grains between the chondrules (Scott 2007, and references therein). Therefore, rocky planetesimals in our solar system must be formed via the co-accretion of chondrules and matrix grains; however, it is not yet understood how chondrite parent bodies are made.

There are a few studies that have investigated the co-accretion process of chondrules and matrices. Ormel et al. (2008) have shown that free-floating chondrules in a protoplanetary disk can obtain a porous dust layer around them, which absorbs the kinetic energy of collisions. Beitz et al. (2012) also suggested that chondrite parent bodies could have formed from clusters of dust-rimmed chondrules based on low-velocity collisional experiments using millimeter-sized glass beads coated with micron-sized silica particles. In addition, granular-mechanics simulations (e.g., Gunkelmann et al. 2017) revealed that dust layers around chondrules lead to a significant increase of the sticking velocity.

However, if fluffy aggregates collide at large velocities, the chondrules contained in the fluffy matrices might move in the aggregates. Moreover, a fluffy aggregate cannot retain its chondrules if the length for stopping the chondrules exceeds the radius of the aggregate. This ejection of chondrules from the fluffy matrices has the potential to prevent chondritic...
planetesimal formation. Therefore, in this study, we examine whether chondrules can be retained inside fluffy aggregates when aggregates collide and grow in a protoplanetary disk.

2. Models

2.1. Outline

The goal of this paper is to explore whether fluffy aggregates composed of chondrules and matrix grains can grow into chondrite parent bodies without chondrule losses. Therefore, it is necessary to calculate the structure of the dust aggregates, such as the radius and the bulk density, and the collision velocity. Here we assume that there are two types of aggregates, that is, chondrule-dust compound aggregates (hereinafter referred to as compound aggregates or CAs) and aggregates that are composed purely of matrix grains (hereinafter referred to as matrix aggregates or MAs). The formation processes of MAs and CAs are illustrated in Figure 1.

In Stage I, the growth of MAs primarily occurs owing to Brownian motion, and they reduce their internal densities via fractal aggregation. We define Stage I as where the timescale of collisions between two MAs $t_{\text{MA-MA}}$ is shorter than the timescale for MAs to accrete onto chondrules $t_{\text{MA-CA}}$. Then, in Stage II, the masses of the CAs increase via the accretion of the grown MAs, and finally, in Stage III, the growth of CAs via the collision of two CAs with static compression takes place. We also define Stage III as where the growth timescale of CAs by the collision of two CAs $t_{\text{grow}}$ is shorter than the growth timescale of CAs by the accretion of the grown MAs $t_{\text{CA-MA}}$. The growth timescale of CAs by the collision of two CAs $t_{\text{grow}}$ is the same as the timescale of collisions between two CAs $t_{\text{CA-CA}}$. On the other hand, the growth timescale of CAs by the accretion of the grown MAs is $t_{\text{CA-MA}} = (m_{\text{CA}}/m_{\text{MA}})t_{\text{CA-MA}}$, where $m_{\text{MA}}$ and $m_{\text{CA}}$ are the masses of MAs and CAs, respectively, and $t_{\text{CA-MA}}$ is the timescale for CAs to collide with MAs. We can imagine that fine-grained matrices could help CAs grow in protoplanetary disks due to their large adhesive force. Forming highly porous matrices also encourages CAs to grow without significant radial drift because fluffy dust aggregates have large collisional cross-sections and can grow rapidly.

In this study, we focus on Stage III and only consider the existence and growth of CAs for simplicity. We calculate the density evolution of the CAs by considering the compression processes and examine whether the CAs can retain chondrules inside their fluffy matrices.

2.2. Structure of the Disk

We assume that the disk structure is in steady state, i.e., the mass accretion rate $\dot{M}$ is constant with distance from the Sun $R$. Because a large amount of fine dust particles must have been present when planetesimals formed in the solar nebula, we also assume that the disk is optically thick and that the main heating source at the midplane is viscous dissipation. Then the local surface density $\Sigma$ and the midplane temperature $T_m$ are given by (e.g., Pringle 1981; Oka et al. 2011)

$$\Sigma = \frac{M}{3\pi \nu_{\text{acc}}},$$

and

$$T_m = \left(\frac{9\kappa \Sigma GM_\odot \dot{M}}{32 \pi \sigma_{\text{SB}} R^3}\right)^{1/4},$$

where $\nu_{\text{acc}}$ is the effective viscosity of the accretion disk, $\kappa$ is the Rosseland mean opacity of the disk medium, $G = 6.67 \times 10^{-8}$ dyn cm$^2$ g$^{-2}$ is the gravitational constant, and $\sigma_{\text{SB}} = 5.67 \times 10^{-5}$ erg cm$^{-2}$ K$^{-4}$ s$^{-1}$ is the Stefan-Boltzmann constant. The solar mass $M_\odot$ is $M_\odot = 1.99 \times 10^{33}$ g. The Kepler frequency $\Omega_k$ is given by $\Omega_k = \sqrt{GM_\odot/R^3}$ and the isothermal sound velocity $c_s$ is given by $c_s = \sqrt{k_B T_m/m_g}$, where $k_B = 1.38 \times 10^{-16}$ erg K$^{-1}$ is the Boltzmann constant and $m_g$ is the mean molecular weight. We set the mean molecular weight to $m_g = 2.34 m_H$, where $m_H = 1.67 \times 10^{-24}$ g is the mass of a hydrogen atom. Then, the effective viscosity $\nu_{\text{acc}}$ is written as $\nu_{\text{acc}} = \alpha_{\text{eff}} c_s^2/\Omega_k$ (Shakura & Sunyaev 1973). Here, we assume $\alpha_{\text{eff}} = 10^{-2}$ in this study (e.g., Hartmann et al. 1998).

The surface densities of dust and gas, $\Sigma_d$ and $\Sigma_g$, respectively, are given by $\Sigma_d = Z \Sigma$ and $\Sigma_g = (1 - Z) \Sigma$, where $Z$ is the vertically integrated dust-to-total mass ratio. The vertical structure of the gas is assumed to be in hydrostatic equilibrium, and the midplane gas density $\rho_g$ is given by $\rho_g = \Sigma_g/(\sqrt{2 \pi} h_g)$, where $h_g = c_s/\sqrt{\kappa q_g}$ is the gas scale height. The mean-free path of gas molecules $\lambda_m$ is given by $\lambda_m = m_g/(\sigma_{\text{mas}} \rho_g)$, where $\sigma_{\text{mas}} = 2 \times 10^{-15}$ cm$^2$ is the collisional cross-section of the gas molecules (Okuzumi et al. 2012). The opacity of the disk $\chi$ is determined by the surface density of the fine dust, i.e., the matrix grains and small MAs. We assume that the mass fraction of matrix grains to all solids $\chi$ is equal in the disk and in the CAs. According to Scott (2007), the volume fraction of the matrices in ordinary chondrites is approximately 10%, and the mass fraction is also approximately 10%. Therefore, we set $\chi = 0.1$ as the canonical value and the opacity of the disk is given by $\chi = \chi Z_{\text{ds}}$, where $Z_{\text{ds}} = 10^3$ cm$^2$ g$^{-1}$ is the mass opacity of the fine silicate grains (e.g., Draine 1985). The mass opacity of highly porous aggregates...
3. Dust Dynamics in the Disk

We consider the Brownian motion, turbulence, and radial and azimuthal drift as sources of the relative velocity of the gas and dust particles. We write the relative velocity \( v \) as the root sum square of these contributions:

\[
\Delta v = \sqrt{(\Delta v_B)^2 + (\Delta v_t)^2 + (\Delta v_L)^2},
\]

where \( \Delta v_B, \Delta v_t, \) and \( \Delta v_L \) are the collisional velocities induced by Brownian motion, turbulence, radial drift, and azimuthal drift, respectively. In this study, we assume that the dust aggregates have no mass distribution and that aggregates that have the same mass have the same volume for simplicity. The collision velocity \( \Delta v \) is

\[
\Delta v = \sqrt{(\Delta v_B)^2 + (\Delta v_t)^2},
\]

where \( \Delta v_B \) and \( \Delta v_t \) are the collisional velocities induced by the Brownian motion and turbulence, respectively. The model of the dust dynamics is the same as that of Arakawa & Nakamoto (2016b).

2.3.1. Brownian Motions

The Brownian-motion-induced velocities \( v_B \) and \( \Delta v_B \) are given by

\[
\Delta v_B = \sqrt{\frac{8k_BT}{\pi m}},
\]

\[
v_B = \sqrt{\frac{8k_BT}{\pi m}},
\]

and

where \( m \) is the mass of an aggregate.

2.3.2. Turbulent Motions

The dynamics of dust particles in a disk is affected by disk gas turbulence. The turnover time and the root-mean-squared random velocity of the largest turbulent eddies, \( t_L \) and \( v_L \), respectively, are given by \( t_L = \Omega_K^{-1} \) and \( v_L = \sqrt{\alpha_L c_s} \), where \( \alpha_L \) is the dimensionless turbulent parameter at the midplane (Cuzzi & Hogan 2003). The value of the dimensionless parameter \( \alpha_L \) for our solar nebula is unclear; however, measurements of CO emission lines reveal that some circumstellar disks, such as the disk around HD 163296, have a small turbulent parameter (\( \alpha_L \lesssim 10^{-3} \) Flaherty et al. 2015). The disk around HL Tau also has a small turbulent viscosity, which is equivalent to an \( \alpha_L \) at a few \( 10^{-4} \) when considering the dust settling suggested by observations of the gap structure (Pinte et al. 2016); meanwhile, \( \alpha_{acc} \sim 10^{-2} \) is indicated by the disk mass distribution and the accretion rate (8.7 \( \times \) 10\(^{-8}\) \( M_\odot\) yr\(^{-1}\); Beck et al. 2010). If the magnetorotational instability is active, \( \alpha_L \) may be up to \( 10^{-3} \) (e.g., Balbus & Hawley 1991), whereas \( \alpha_L \) might be lower than \( 10^{-5} \) to \( 10^{-4} \) if we focus on the low-ionization-fraction regions (e.g., Gammie 1996; Mori & Okuzumi 2016). In this study, we assume that the turbulence is weak in the midplane of the inner disk and that \( \alpha_L \approx 10^{-4} \). The turbulent viscosity \( v_L \) is given by \( v_L = \alpha_L c_s H_g \).

The extent of the gas turbulence is determined by the turbulent Reynolds number \( Re_t \). The turbulent Reynolds number \( Re_t \) is given by \( Re_t = \frac{v_t}{\nu_m} \), where \( v_t = \lambda_m \nu_{th}/2 \) is the molecular viscosity, and \( \nu_{th} = \sqrt{8/\pi \xi c_s} \) is the thermal velocity of the gas molecules. The turnover time and the root-mean-square random velocity of the smallest eddies, \( t_t \) and \( v_t \), respectively, are given by \( t_t = Re_t^{-1/2} t_L \) and \( v_t = Re_t^{-1/4} v_L \).

The key parameter that determines the motion of dust aggregates in the gas disk is the normalized stopping time, the Stokes number \( St \), given by

\[
St = \frac{\Omega_K t_s},
\]

where \( t_s \) is the stopping time of the dust aggregates. The stopping time depends on the radius of the aggregates \( a \) and the particle Reynolds number \( Re_p = 2av/\nu_m \). The stopping time \( t_s \) is given by Probstein & Fassio (1970) and Weidenschilling (1977) as follows. If the radius of an aggregate \( a \) is smaller than \( (9/4)\lambda_m \), then the stopping time of the dust aggregates is determined by Epstein’s law, \( t_s^{Ep} = 3m/(4\pi \rho_b \nu_{th} a^2) \). Conversely, when the radius \( a \) is larger than \( (9/4)\lambda_m \), the stopping time depends on the Reynolds number. If the particle Reynolds number \( Re_p \) is less than unity, the motion of the aggregate is determined by Stokes’ law, \( t_s^{St} = m/(6\pi \rho_b \nu_{th} a) \), while for the case of \( 1 \lesssim Re_p < 54^{1/3} \approx 800 \), the stopping time is obtained from Allen’s law, \( t_s^{Al} = (23/5 m)/(12\pi \rho_b \nu_{th} a^{3/5} u^{2/5}) \). Finally, for the case in which the aggregate moves fast and the particle Reynolds number is larger than \( 54^{1/3} \), we use Newton’s law to obtain the stopping time, \( t_s^{New} = (9m)/(2\pi \rho_b \nu_{th} a^2) \).

We use the analytic formulae derived by Ormel & Cuzzi (2007) for the turbulence-driven velocity. If the stopping time \( t_s \) is shorter than the turnover time of the smallest eddies \( t_t \), the relative velocity between the dust particle to disk gas \( v_t \) is given by \( v_t^S \) such that

\[
v_t^S = Re_t^{1/4} St v_L.
\]

For the case of \( t_t \ll t_s \ll t_L \), the relative velocity \( v_t \) is given by \( v_t^M \) such that

\[
v_t^M = 1.7 St^{1/2} v_L,
\]

and for the case of \( t_s \gg t_L \), the relative velocity \( v_t \) is given by \( v_t^L \) such that

\[
v_t^L = \left(1 + \frac{1}{1 + St}\right)^{1/2} v_L.
\]

Here, we assume that the turbulence-driven velocity \( v_t \) is given by the minimum of these three terms for continuity and better handling of \( v_t \):

\[
v_t = \min(v_t^S, v_t^M, v_t^L).
\]

The turbulence-driven velocity between two dust particles was also obtained by Ormel & Cuzzi (2007). We assume the collision velocity of two particles with a stopping time \( t < t_s \) to be

\[
\Delta v_t^S = \varepsilon Re_t^{1/4} St v_L,
\]

where \( \varepsilon = 0.1 \) is a nondimensional value representing the variation in the density fluctuation between same-mass
aggregates, which is neglected in this study (e.g., Okuzumi et al. 2011). Note that there is a large uncertainty in the estimation of the \( \varepsilon \) for CAs; the estimation by Okuzumi et al. (2011) is only valid for non-compressed fractal aggregates composed of \( 10^{-6} \) monomers. For the intermediate region, we assume \( \Delta v_i^M \) such that

\[
\Delta v_i^M = 1.4St^{1/2}v_L, \quad (13)
\]

and the relative velocity for high Stokes number aggregates \( \Delta v_i^L \) is

\[
\Delta v_i^L = \left( \frac{2}{1 + St} \right)^{1/2}v_L. \quad (14)
\]

Here, we also assume that the turbulence-driven relative velocity \( \Delta v_i \) is given by the minimum of these three terms:

\[
\Delta v_i = \min(\Delta v_i^S, \Delta v_i^M, \Delta v_i^L). \quad (15)
\]

2.3.3. Systematic Motions

Dust aggregates in the gas disk have systematic velocities such as the radial drift velocity, the azimuthal velocity, and the settling velocity. In this study, we consider the radial drift velocity \( v_r \) and the azimuthal velocity \( v_\phi \) given by (Weidenschilling 1977)

\[
v_r = -\frac{2St}{1 + St^2}\eta v_K, \quad (16)
\]

and

\[
v_\phi = \left( 1 - \frac{1}{1 + St^2} \right)^{\eta v_K}, \quad (17)
\]

where \( \eta v_K \) is the difference between the Keplerian velocity and the gas rotational velocity. The Keplerian velocity \( v_K \) is given by \( v_K = R\Omega_K \), and the dimensionless coefficient \( \eta \) is given by (Nakagawa et al. 1986)

\[
\eta = -\frac{1}{2} \left( \frac{c_s}{v_K} \right)^2 \frac{\partial \ln(\rho_s c_s^2)}{\partial \ln R}. \quad (18)
\]

2.4. Density of the Compound Aggregates

The aerodynamic properties of the aggregates depend on their bulk densities. In this paper, we assume that all CAs have the same mass and density and that the matrix regions of the CAs have a uniform density. Then, the bulk filling factor of the CAs \( \phi_{\text{CA}} \) is given by

\[
\phi_{\text{CA}} = \left( \frac{\chi}{\phi_{\text{mat}}} + \frac{1 - \chi}{1} \right)^{-1}, \quad (19)
\]

where \( \phi_{\text{mat}} \) is the filling factor of the matrix regions and \( \chi \) is the mass fraction of the matrices in a CA. The bulk density of the CAs \( \rho_{\text{CA}} \) is

\[
\rho_{\text{CA}} = \rho_0 \phi_{\text{CA}}, \quad (20)
\]

where \( \rho_0 = 3 \text{ g cm}^{-3} \) is the material density of the chondrules and matrix grains.

There are several processes that affect the density of the dust aggregates: the hit-and-stick growth of two colliding aggregates without compaction (e.g., Meakin 1991; Okuzumi et al. 2009), the dynamical compression via high-speed collisions (e.g., Paszun & Dominik 2009; Suyama et al. 2012), and static compression via the ram pressure of the disk gas and/or via the self-gravity of large planetary bodies (e.g., Seizinger et al. 2012; Kataoka et al. 2013a).

Dynamical compression, however, does not work in our calculations, and the CAs are compressed by the static pressure. Note that Ormel et al. (2008) assumed that dust rims around chondrules are formed via direct accretion of monomer grains. Therefore, Ormel et al. (2008) assumed that the initial filling factor of the matrix regions of CAs is \( \phi_{\text{mat}} = 0.15 \), which is the filling factor obtained by the particle-cluster aggregation. By contrast, we assume chondrules obtain fluffy dust rims via the accretion of highly porous MAs, and \( \phi_{\text{mat}} \) is far smaller than 0.15 at the start of Stage III in our calculations. In addition, Ormel et al. (2008) employs the collisional compression model obtained by Ormel et al. (2007), however, Wada et al. (2008) revealed that Ormel et al. (2007) overestimates the density of fluffy aggregates at the maximum compression.

Kataoka et al. (2013a) revealed that the static compression of a highly porous aggregate depends on the pressure that works on the aggregate \( P \) and the rolling energy \( E_{\text{roll}} \). The rolling energy is given by (Dominik & Tielens 1995)

\[
E_{\text{roll}} = 6\pi^2 a_0 \xi, \quad (21)
\]

where \( \gamma \) is the surface energy of the silicate, \( a_0 \) is the radius of the monomers (i.e., matrix grains), and \( \xi \) is the critical displacement for rolling. We assume that the radius of the matrix grains is \( a_0 = 0.025 \mu m \) in this study, which is the typical radius of opaque grains in pristine chondrites such as Acfer 094 and QUE 99177 (Vaccaro et al. 2015). We employ the canonical value of the surface energy of the silicate, which is \( \gamma = 25 \text{ erg cm}^{-2} \) (Kendall et al. 1987). The lower limit of \( \xi \) is 0.2 nm, which is equal to the interatomic distance (Dominik & Tielens 1997), and the upper limit is the radius of the contact surface area of two matrix grains \( a_{\text{cont}} \) (Heim et al. 1999). Krijt et al. (2014) showed that \( \xi \) is approximately 0.3–1 nm for micron-sized amorphous silica spheres, and we assume that the critical displacement for rolling is \( \xi = 0.3 \text{ nm} \) for submicron-sized monomers. Therefore, we can calculate the filling factor of a highly porous dust aggregate using the relation obtained by Kataoka et al. (2013a), and the filling factor of matrices in CAs is given by

\[
\phi_{\text{mat}} = \left( \frac{a_0^3 P}{E_{\text{roll}}} \right)^{1/3}. \quad (22)
\]

The sources of the pressure \( P \) are the ram pressure \( P_{\text{gas}} \) and the self-gravitational pressure \( P_{\text{grav}} \) (Kataoka et al. 2013b). For CAs whose radii are \( a_{\text{CA}} \) and masses are \( m_{\text{CA}} \), the ram pressure of the disk gas is estimated to be \( P_{\text{gas}} = (m_{\text{CA}} v^2)/(2\rho_{\text{gas}} a_{\text{CA}}^2) \), whereas the self-gravitational pressure is \( P_{\text{grav}} = (G m_{\text{CA}}^2)/(\pi a_{\text{CA}}^4) \). We calculate the filling factors and densities assuming that \( P = \max(P_{\text{gas}}, P_{\text{grav}}) \).

2.5. Penetration of Chondrules in Matrices

Chondrules contained in fluffy matrices move in the aggregates when fluffy aggregates collide at large velocities. We estimate that these CAs cannot retain the chondrules inside their matrices if the stopping length of the chondrules exceeds...
the radius of the CAs (Figure 2). Here, we assume that the stopping length is the penetration length for stopping via deceleration forces, \( L_{\text{pen}} \).

The penetration of chondrules in fluffy matrices is the same process as that for the penetration of spherical projectiles in a silica aerogel (Niimi et al. 2011) and the intrusion of chondrule-analogs into matrix analogs (Machii et al. 2013). Therefore, the equation of motion is

\[
F_{\text{pen}} = \frac{m_{\text{ch}} dv_{\text{pen}}}{dt} = F_{\text{pen}},
\]

where \( m_{\text{ch}} \) is the mass of the chondrules. In this study, we assume that the radius of the chondrules is \( a_{\text{ch}} = 0.01 \text{ cm} \), which is a typical size of chondrules in H ordinary chondrites and EH enstatite chondrites (Rubin 2000). The mass of the chondrules is \( m_{\text{ch}} = (4\pi/3)a_{\text{ch}}^3 \rho_0 \). Note that some other groups (e.g., L and LL ordinary chondrites) contain chondrules whose radii are a few times larger than that, which we assume here.

There are two mechanisms for decelerating chondrules in matrices: “hydrodynamic” drag for high-speed penetration and “crushing” drag for low-speed penetration (e.g., Niimi et al. 2011). In the case of high-speed penetration, the deceleration process is due to the “hydrodynamic” drag caused by the ram pressure, and the drag force \( F_{\text{pen}} \) is given by

\[
F_{\text{HD}} = -\frac{C_d}{2} \pi a_{\text{ch}}^2 \rho_{\text{mat}} v_{\text{pen}}^2,
\]

where \( C_d = 1 \) is the drag coefficient obtained from impact experiments (e.g., Trucano & Grady 1995). Conversely, in the case of low-speed penetration, the deceleration process is due to “crushing” against the compressive strength of the matrices and the equation of motion is given by

\[
F_{\text{CR}} = -\pi a_{\text{ch}}^2 P_c,
\]

where \( P_c \) is the compressive strength of the matrices (e.g., Domínguez et al. 2004). The compressive strength \( P_c \) has been studied in experiments (e.g., Blum & Schräpler 2004; Güttler et al. 2009) and in numerical simulations (e.g., Seizinger et al. 2012; Kataoka et al. 2013a). In this study, we calculate the \( P_c \) of highly porous matrices using the following equation (Kataoka et al. 2013a):

\[
P_c = \frac{E_{\text{mat}}}{\alpha_0^3} \phi_{\text{mat}}^3.
\]

Here, we can rewrite the equation of motion as

\[
\frac{dv_{\text{pen}}}{dt} = -\max(Av_{\text{pen}}^2, B),
\]

where \( A = 3C_d \phi_{\text{mat}}/(8a_{\text{ch}}) \) and \( B = 3P_c/(4\rho_0 a_{\text{ch}}) \). Then we can obtain the penetration length \( L_{\text{pen}} \) as a function of the initial penetration velocity \( v_0 \). In the case of high-speed penetration (\( v_0 \geq \sqrt{B/A} \)), \( L_{\text{pen}} \) is given by

\[
L_{\text{pen}} = \frac{1}{A} \ln(v_0 \sqrt{A/B}) + \frac{1}{2A},
\]

and in the case of low-speed penetration (\( v_0 < \sqrt{B/A} \)), \( L_{\text{pen}} \) is given by

\[
L_{\text{pen}} = \frac{v_0^2}{2B}.
\]

If two colliding CAs have the same mass and the matrices of the two colliding aggregates stop immediately after the collision, then the initial penetration velocity \( v_0 \) can be evaluated as \( \Delta v/2 \). However, the motion of the matrices after the collision is not yet understood. We examine two cases in this study: \( v_0 = \Delta v/2 \) and \( v_0 = \Delta v/8 \). Note that, if the mass ratio of the projectile to the target is small, then \( v_0 \) of the target decreases while \( v_0 \) of the impactor increases when considering the motions relative to the center of mass.

2.6. Critical Velocity for Collisional Growth

We also estimate the critical velocity for the collisional growth of aggregates \( \Delta v_{\text{crit}} \) from the scaling relation obtained by Wada et al. (2009). The impact energy of two colliding aggregates \( E_{\text{imp}} \) is \( E_{\text{imp}} = m(\Delta v)^2/4 \), and the critical impact energy for growth \( E_{\text{imp,crit}} \) is \( E_{\text{imp,crit}} = 30N_{\text{total}}E_{\text{break}} \), where \( N_{\text{total}} \) is the total number of monomers contained in the colliding aggregates, and \( E_{\text{break}} \) is the energy for breaking a single contact between two monomers (Wada et al. 2009). The energy for breaking a monomer–monomer contact \( E_{\text{break}} \) is given by Wada et al. (2007). Clearly, the total binding energy of a CA is dominated by the contacts between matrix grains. Therefore, the total number of monomers in the two colliding CAs is \( N_{\text{total}} = 2\chi m_{\text{CA}}/m_0 \), where \( m_0 = (4\pi/3)a_0^3 \rho_0 \) is the mass of monomers, and the critical velocity for the collisional growth of CAs \( \Delta v_{\text{crit}} \) is given by

\[
\Delta v_{\text{crit}} = \sqrt{\frac{240\chi E_{\text{break}}}{m_0}} = 5.5 \times 10^5 \text{ cm s}^{-1}.
\]

Note, however, that the critical collision velocity of dust aggregates is still open to argument (e.g., Tanaka et al. 2012; Krijt et al. 2013). Hence we do not consider fragmentation of CAs in this study.

3. Results

First, we calculate the pathway of the dust aggregate growth in radius–density space and investigate whether the growth of the CAs is sufficiently rapid to avoid non-negligible radial drift. Here, we neglect bouncing, erosion, and fragmentation for simplicity; therefore, the timescale of collisional growth \( t_{\text{grow}} \) is
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defined as \( t_{\text{grow}} = m/(\pi a^2 \rho_d \Delta v) \), where \( \rho_d = \sum_d/(\sqrt{2\pi} h_d) \) is the midplane dust density and \( h_d \) is the dust scale height. The dust scale height is given by (Youdin & Litchwill 2007)

\[
h_d = h_g \left(1 + \frac{St}{\alpha} + \frac{2St}{1 + St}\right)^{-1/2}.
\]

The timescale of radial drift \( t_{\text{drift}} \) is defined as the orbital radius divided by the radial drift velocity, \( t_{\text{drift}} = R/v_r \). Aggregates grow without significant radial drift if the condition \( t_{\text{grow}} < (1/30) t_{\text{drift}} \) is satisfied (Okuzumi et al. 2012). Therefore, the radial drift barrier (the pink region) and the growth pathway of the CAs (the black and violet dashed lines) can be drawn as shown in Figure 3. The black lines represent where the collision velocity \( \Delta v \) is below the critical velocity \( \Delta v_{\text{crit}} \), while the violet lines represent where \( \Delta v > \Delta v_{\text{crit}} \). We also calculated the penetration length of the chondrules in CAs \( L_{\text{pen}} \) as a function of \( a_{\text{CA}} \) and \( \rho_{\text{CA}} \). In Figure 3, we show the “ejection barrier,” which represents the region where the penetration length \( L_{\text{pen}} \) exceeds the radius of the CAs \( a_{\text{CA}} \). The forest-green region represents the ejection barrier obtained by assuming that the initial penetration velocity of \( v_0 = \Delta v/2 \), while the yellow-green region represents the case where \( v_0 = \Delta v/8 \).

The upper panel of Figure 3 illustrates that the pathway of the CAs overcomes the radial drift barrier but fails to avoid the ejection barrier. In this case, we assume \( M = 10^{-7} M_\odot \, \text{yr}^{-1} \) (which is comparable to the accretion rate of the disk around HL Tau; Beck et al. 2010), \( R = 1 \, \text{au} \) (which is the distance from the Sun to the Earth), and \( Z = 0.0043 \) (e.g., Miyake & Nakagawa 1993) in this case. The radius and the density of the chondrules are represented by the circular marker, and the radius and the density of the CAs, which contain one chondrule, are represented by the square marker. Thus the square marker is equivalent to the start of Stage III. If we do not consider the ejection of chondrules from the fluffy matrices, the CAs can grow into planetesimals without radial drift and the pentagonal marker represents the onset of runaway growth (\( \Delta v = \sqrt{2GM_{\text{CA}}/a_{\text{CA}}}; \) Wetherill & Stewart 1989; Kobayashi et al. 2016). Note that the onset of runaway growth might be affected by turbulence-induced density fluctuations (Okuzumi & Ormel 2013; Ormel & Okuzumi 2013) and/or the gravitational instability of the porous planetesimal disk (Michikoshi & kokubo 2016, 2017). Therefore, we do not discuss the onset of runaway growth in this study but focus on the growth of CAs whose Stokes numbers are less than unity.

The filled diamond indicates the onset of critical ejection, \( L_{\text{pen}} = a_{\text{CA}} \), for the severe estimate (\( v_0 = \Delta v/2 \)), and the open diamond indicates the end of critical ejection for the severe estimate (\( v_0 = \Delta v/8 \)). Our results show that, if the density of the CAs is determined by static compression, then CAs cannot grow into planetesimals without losing chondrules, even if we evaluate the penetration length using a lenient assumption. Our calculations also reveal that CAs face the ejection barrier before they grow into meter-sized bodies and face the fragmentation barrier.

In addition, the recapture of chondrules by chondrule-free aggregates is not sufficient to explain the formation of chondrite parent bodies. The lower panel in Figure 3 shows that the Stokes numbers of the chondrules and the CAs. The Stokes number of the chondrules is \( St_{\text{ch}} = 5.5 \times 10^{-5} \), and the Stokes number of the CAs when they reach the end of critical ejection is \( St_{\text{CA}} = 0.12 \) for the severe estimate and \( St_{\text{CA}} = 6.0 \times 10^{-2} \) for the lenient estimate. According to Equation (31), the scale height of the CAs is \( h_{\text{CA}} \sim (1/30) h_g \) while the scale height of the chondrules is \( h_{\text{ch}} \sim h_g \). The collision velocity of two CAs is \( 1.4St_{\text{CA}}^{1/2}v_{\text{L}} \) (Equation (13)) and the relative velocity between the colliding chondrule to the CA is approximately \( 1.7St_{\text{ch}}^{1/2}v_{\text{L}} \) (Equation (9)). This is because the Stokes number of the chondrules \( St_{\text{ch}} \) is far smaller than the Stokes numbers of the CAs \( St_{\text{CA}} \). Therefore, the timescale for chondrules to accrete onto CAs \( t_{\text{ch-CA}} \) is dozens

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**Figure 3.** Evolutionary track of the CAs (the black and violet dashed lines), the radial drift barrier (the pink region), and the ejection barrier (the forest-green regions), for the case of \( M = 10^{-7} M_\odot \, \text{yr}^{-1} \), \( R = 1 \, \text{au} \), and \( Z = 0.0043 \). The black lines represent where \( \Delta v < \Delta v_{\text{crit}} \) and the violet lines represent where \( \Delta v > \Delta v_{\text{crit}} \). The upper panel shows that the density evolution track of CAs is determined by static compression. We show two barriers in radius-density space. The lower panel represents the change in the Stokes number as a function of the radius. It also indicates the appropriate area for driving the streaming instability in a disk with a moderate dust-to-total disk mass ratio, \( 10^{-2} \lesssim St_{\text{CA}} \lesssim 1 \) (the yellow region). The circles indicate the values of single chondrules, the filled triangles and diamonds indicate the onset of critical ejection (\( L_{\text{pen}} = a_{\text{CA}} \)) while the open triangles and diamonds represent the ends. Triangles correspond to the severe estimate (\( v_0 = \Delta v/2 \)) and diamonds correspond to the lenient estimate (\( v_0 = \Delta v/8 \)). The pentagonal markers indicate the onset of runaway growth (\( \Delta v = \sqrt{2GM_{\text{CA}}/a_{\text{CA}}} \)).
of times longer than the timescale of collisions between two CAs. This result indicates that CAs will have already grown much more than $2^{10}$ times in mass and more than 200 times in radii, before the chondrules accrete onto the CAs. Conversely, the penetration length $L_{\text{pen}}$ is given in Equation (28) and $L_{\text{pen}} \sim 1/A$ hardly changes even when the CAs grow 10 times in radii. Strictly speaking, the initial penetration velocity $v_0$ is not $\Delta v/2$ but approximately $v$ for the case of recapture; however, this would not significantly affect the penetration length for the reason described above. Therefore, taking into consideration that the radius of the CAs is on the order of 100 cm at the end of critical ejection, the chondrite parent bodies would consist of rocks with meter-sized chondrule-free matrices. This is, however, inconsistent with observational evidence. One of the most important pieces of evidence is that the most numerous type of meteorites is chondrule-free matrices. This is, however, inconsistent with observational evidence. One of the most important pieces of evidence is that the most numerous type of meteorites is chondrites (Scott 2007), and the typical size of the recovered chondrites is 10 cm. This fact excludes the possibility of the recapture of chondrules by meter-sized fluffy aggregates.

We show the results for other disk parameters in Figure 4. Figure 4(a) shows the result for a case with $M = 10^{-8} M_\odot$, $R = 1$ au, and $Z = 0.0043$. This is a low-accretion-rate case compared to the case in Figure 3. The result shows that CAs cannot retain chondrules even in this case, and that the onset of critical ejection strongly depends on the size of the matrix and the chondrules. Conversely, the end of critical ejection depends little on the initial penetration velocity $v_0$. This is because the onset is determined by Equation (29), so the penetration length depends on $v_0^2$, while the end is determined by Equation (28), so the dependence of $L_{\text{pen}}$ on $v_0$ is logarithmic.

Figure 4(b) shows the result for the case of $M = 10^{-7} M_\odot$, $R = 1.5$ au, and $Z = 0.0043$. The difference between this case and the case in Figure 3 is the distance from the Sun. The radial drift barrier is sensitive to the radial distance from the Sun, so CAs can grow into planetesimals only within $R \lesssim 1.5$ au even if we do not consider the ejection and fragmentation problems. Similarly, Figure 4(c) shows the result for the case of $M = 10^{-7} M_\odot$, $R = 1$ au, and $Z = 0.043$ and the difference between this case and the case in Figure 3 is the mass fraction of dust. The radial drift barrier is sensitive to the dust mass fraction $Z$, and there is no radial drift barrier if $Z$ is approximately 10 times higher than the canonical silicate abundance inferred from the solar metallicity ($Z = 0.0043$; Miyake & Nakagawa 1993). Nevertheless, CAs might lose chondrules before their Stokes number reaches approximately $10^{-2}$. The possibility of the recapture of chondrules by grown fluffy aggregates is, however, denied in either case for the same reason as that in the case of Figure 3.

### 4. Analytic Estimates

From the results shown in Figures 3 and 4, we can divide the ejection barrier into four regions; the Brownian-motion region, the Epstein-drag region, the Stokes-drag region, and the hydrodynamical region. In this section, we derive analytical formulae for the Epstein-drag region, the Stokes-drag region, and the hydrodynamical region.

#### 4.1. Epstein-drag Region

First, we define the “Epstein-drag region” as the region in which the CA affects the Epstein drag and the chondrules are decelerated by the compressive stress of the matrices. In the upper panel in Figure 3, this region appears within the range between approximately 0.3 and 8 cm for the lenient case and between approximately 0.1 and 8 cm for the severe case. From our calculation, the collision velocity is given by Equation (12) and we can rewrite Equation (29) as

$$L_{\text{pen}} \approx \frac{\pi^2}{6} \left( \frac{v_0}{\Delta v} \right)^2 \frac{a M_{\text{ch}}}{n^2} \frac{\rho_{\text{CA}}}{\rho_0} \frac{\lambda^2}{\lambda^2} \left( \frac{E_{\text{roll}}}{a_{\text{ch}}} \right)^{-1} \frac{1}{a_{\text{ch}}}. \quad (32)$$

Here, we use the fact that the volume of a fluffy CA is dominated by its matrices; therefore, the filling factor of the
The critical bulk density for retaining chondrules is
\[ \rho_{\text{CA}} \gtrsim \frac{8 \chi \rho_0 a_{\text{ch}}}{3 a_{\text{CA}}}. \]  

The critical bulk density depends little on the properties of matrix grains in the hydrodynamic region, while the properties of matrix grains (i.e., \( E_{\text{roll}}/a_0^3 \)) strongly affect the critical bulk density in the Epstein-drag and Stokes-drag regions.

In all regions, the penetration length is proportional to the radius of the chondrules. We assume \( a_{\text{ch}} = 0.01 \) cm in this study, however, the sizes of chondrules vary between different types of chondrites. The mean radius of chondrules in H ordinary chondrites and EH enstatite chondrites is approximately 0.01 cm, whereas the chondrules in L and LL ordinary chondrites and EL ordinary chondrites are approximately 0.03 cm in radii (Rubin 2000, and references therein). By contrast, the spread of sizes in individual groups is relatively small. Approximately 70% of chondrules in Qingzhen, Kota-Kota, and Allan Hills A77156 EH3 chondrites have similar radii within a factor of two (Rubin & Grossman 1987). It might be possible to derive some constraints on the accretion environment of CAEs from the differences in chondrule size in the different groups of chondrites. We will discuss this topic in the future.

5. Discussion

5.1. The Complementarity of Chondrules and Matrices

Recent high-precision measurements of chondrules, matrices, and bulk chondrites reveal that chondrites have chemical (e.g., Palme et al. 2015; Ebel et al. 2016) and isotopic (e.g., Budde et al. 2016a, 2016b) complementarities. These complementarities are therefore a strong constraint, indicating that the chondrules and the matrices must have formed from a single reservoir and that after their formation, neither the chondrules nor the matrix grains were lost. The simplest way to explain the chondrule-matrix complementarity is that the chondrules and the matrix grains were formed in the same heating events and that some parts of the matrix grains are condensates of evaporated dust. Miura et al. (2010) showed that a shock wave heating model for chondrule formation can predict the evaporation and formation of fine dust grains via chondrule formation; in addition, Miura & Nakamoto (2005) revealed that the minimum size of chondrules that avoids the complete evaporation is consistent with the observations (e.g., Eisenhour 1996; Nelson & Rubin 2002).

There are many models for chondrule formation, and these models are classified into two groups, that is, impact origin models (e.g., Sanders & Scott 2012; Dullemond et al. 2014; Wakita et al. 2017) and nebular origin models (e.g., Muranushi 2010; Boley et al. 2013; McNally et al. 2013). To determine the formation process(es) of chondrules, the key constraints are the heating and cooling timescales. According to Tachibana & Huss (2005), the chondrule-forming events must be flash heating events, i.e., a heating rate of \( 10^4-10^6 \) K h\(^{-1}\) would be required to avoid the isotopic fractionation of sulfur in chondrule precursors. In addition, although there are huge uncertainties, several studies suggest that the cooling rate of chondrules is as large as \( 10^2-10^6 \) K h\(^{-1}\) from observations of iron–magnesium and oxygen isotopic exchange (e.g., Yurimoto & Wasson 2002) and the crystal growth of olivine phenocrysts (e.g., Wasson & Rubin 2003; Miura & Yamamoto 2014).
These estimates are larger by several orders of magnitude than the cooling rates of $10^{-3}$ K h$^{-1}$ obtained from furnace-faceted experiments (Desch et al. 2012, and references therein).

Even though it is true that furnace-faceted experiments have succeeded in reproducing some textural features of chondrules (e.g., Hewins & Fox 2004; Tsuchiyama et al. 2004), the solidification processes might be affected by the contact with a furnace wall (Nagashima et al. 2006). Furthermore, rapid cooling of levitating supercooled precursors also succeeds in reproducing some textures of chondrules (e.g., Srivastava et al. 2010). The supercooling of chondrule melts is preferable to explain the textural features of compound chondrules, which might form via the collision of chondrule precursors (Arakawa & Nakamoto 2016a). The rapid cooling of the chondrules indicates that the evaporated dust would also be cooled rapidly, and the cooling rate of the dust vapor affects the size and the morphology of the condensates (e.g., Miura et al. 2010).

### 5.2. Influence of the Monomer Properties

Understanding of the physical properties of the matrix grains is important when considering the growth process of CAs. As shown in Section 4, the penetration length $L_{\text{pen}}$ is inversely proportional to $E_{\text{roll}}/a_0^3$, and the rolling energy of two sticking matrix grains $E_{\text{roll}} = 6\pi^2\gamma a_0\xi$. Therefore, we can obtain the relationship between the surface energy, the radius, and the penetration length, $L_{\text{pen}} \propto \gamma^{-1}a_0^2$, and large surface energies and small grain sizes are preferable when the bulk density is fixed. We append that if the bulk density of the CAs is determined by static compression, $\rho_{\text{CA}}$ also depends on $E_{\text{roll}}/a_0^3$; however, the dependence is weak (e.g., Kataoka et al. 2013b; Arakawa & Nakamoto 2016b).

Here we mention the fact that the exact value of $E_{\text{roll}}$ might differ from $6\pi^2\gamma a_0\xi$. This is because the equation $E_{\text{roll}} = 6\pi^2\gamma a_0\xi$ obtained by Dominik & Tielens (1995) is based on the so-called JKR contact model (Johnson et al. 1971), and the JKR model is valid only for large values of the Tabor parameter $\mu$ (Tabor 1977):

$$\mu = \left( \frac{2a_0\gamma^2(1 - \nu^2)}{Y^2z_0^3} \right)^{1/3} = 0.11 \left( \frac{a_0}{0.025 \ \mu\text{m}} \right)^{1/3},$$  \hspace{1cm} (40)

where $\nu = 0.17$ is the Poisson’s ratio, $Y = 5.4 \times 10^{11}$ dyne cm$^{-2}$ is the Young’s modulus, and $z_0 = 0.2$ nm is the interatomic distance. The JKR model can be used only for $\mu \gtrsim 5$ (Johnson & Greenwood 1997), and thus other contact models should be used for submicron-sized silicate spheres. In addition, the dependence on the sphere radius is weak, thus we should not adopt the JKR theory not only for submicron-sized spheres but also micron- or millimeter-sized silicate grains. Even if we consider chondrule-sized grains, the Tabor parameter is still small; the Tabor parameter is $\mu = 1.7$ when the sphere radius is $a_0 = 0.01$ cm. The Maugis-Dugdale model (Maugis 1992) might be better than the JKR model where the Tabor parameter is $0.1 \lesssim \mu \lesssim 5$, and several contact models such as the so-called DMT model (Derjaguin et al. 1975) and the semi-rigid sphere model (Greenwood 2007) are suggested for $\mu \lesssim 0.1$. However, the maximum force needed to separate the two contact spheres $F_c$ is hardly dependent on the selection of contact models: $F_c = (3/2)\pi\gamma a_0$ for the JKR model and $F_c = 2\pi\gamma a_0$ for the DMT model. Additionally, the rolling energy $E_{\text{roll}}$ is proportional to $F_c\xi$ (e.g., $E_{\text{roll}} = 4\pi F_c\xi$ for the JKR model), therefore the rolling energy might be also subequal among contact models. Either way, future studies of the rolling resistance between two stiff spheres are necessary.

The morphology of the matrix grains also influences both the penetration process and the density evolution because the irregular shape of the monomers affects the rolling energy. In addition, the irregular shape affects not only the density and the penetration process but also the critical velocity for collisional growth (Poppe et al. 2000). It has been suggested that at least some parts of the matrix grains are condensates of evaporated dust and that the size of the condensate ranges from nanometer to micrometer (e.g., Yamamoto & Hasegawa 1977; Miura et al. 2010) depending on the cooling rate and the mass density of the dust vapor. Ishizuka et al. (2015) experimentally revealed that there are two types of nanometer-sized silicate condensates obtained from the re-condensation of magnesium and silicon monoxide vapor, that is, spherical forsterite particles and cubic MgO particles. However, there is no reliable contact theory that can be applied to non-spherical monomers. Therefore, improvements on the contact mechanical theory for non-spherical monomers are also necessary.

### 5.3. Dust Rims around the Chondrules

In this study, we assume the uniform density of the matrix regions throughout the CA for simplicity. However, the gradient and fluctuation of the density of the matrix regions of a CA might exist to some extent. Considering the accretion process in Stage II, the initial growth of chondrules/CA is driven by the accretion of small and dense MAs. If the dense dust rims around chondrules can be maintained when CAs collide at a large velocity, dust-rimmed chondrules in the fluffy CA are subject to a large drag force $F_{\text{pen}}$ and the penetration length $L_{\text{pen}}$ of dust-rimmed chondrules might be smaller than the penetration length of naked chondrules. Although it is still unclear whether chondrules can maintain the dense dust rims or not, we should consider the formation and influence of the dust rims around chondrules.

### 5.4. Deceleration by Collisions of Chondrules

In this paper, we only consider the drag force for penetration $F_{\text{pen}}$ as a decelerating mechanism. Even though the narrow size distribution of chondrules (e.g., Nelson & Rubin 2002) might prevent chondrules from colliding unless there is a large variation in the initial velocity of penetration, chondrules in a large CA are likely affected by chondrule collisions when they move in the CA. In addition, the recapture of chondrules by large CAs is likely affected by this process because there is a large velocity difference between intruding chondrules and chondrules contained in CAs. We will examine this effect in the future.

### 5.5. Alternative Scenarios

As shown in Section 3, CAs would lose their chondrules if the bulk density of these aggregates were determined by static compression processes. Therefore, we need to consider alternative scenarios for chondritic planetesimal formation. Here, we propose two scenarios: changing the density evolution pathway of CAs to retain chondrules in CAs and...
escaping from the ejection barrier of planetesimal formation via the streaming instability before CAs lose their chondrules.

5.5.1. Modification of the Density Evolution Pathway

If chondrite parent bodies form via direct aggregation of CAs, then these aggregates need to avoid both the ejection barrier and the radial drift barrier. Here, we can recognize the “channel” of the two barriers in Figures 3 and 4. If the pathway of density evolution passes the channel, then the chondrite parent bodies might form via direct aggregation. Note that the density determined by static compression is the lower limit of the density evolution and that the bulk density of the CAs could increase if we nullify the assumption that CAs grow via collisions between similar-sized aggregates. Dominik (2009) showed that the hierarchical growth of large aggregates via the accretion of small aggregates leads to aggregates with a relatively dense structure. In addition, aggregates formed via hierarchical growth can explain the small tensile strength of porous matrices. In this study, we introduced a density evolution model of chondrule-dust compound aggregates and calculated the evolutionary track of porous aggregates. In addition, we estimated the stopping length of the chondrules in the matrices. For simplicity, we assumed that the chondrules are only decelerated by the drag force for penetration and the initial penetration velocity of the chondrules was parameterized due to the uncertainty in the collision behavior.

The main finding of this study is that fluffy aggregates likely lose their chondrules before they grow to meter-sized bodies, even if we evaluate the stopping length using lenient assessments (Figures 3 and 4). This “ejection barrier” might also prevent chondrule-dust compound aggregates from growing into planetesimals via the streaming instability. Therefore, a change in the density evolution pathway is necessary to form chondrite parent bodies in our solar system.

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5.5.2. Jumping over the Ejection Barrier

It is expected that chondrite parent bodies can be formed via gravitationally bound clumps of CAs and that the formation of dust clumps is triggered by the streaming instability if the onset of the instability is faster than the onset of critical ejection. As shown in Figures 3 and 4, the critical ejection of chondrules from CAs starts when the Stokes number of the CAs $St_{CA}$ reaches the order of $10^{-5}$–$10^{-3}$, depending on the disk parameters and the evaluation of the initial penetration velocity. However, Carrera et al. (2015) showed that whether the streaming instability occurs strongly depends on the Stokes number and that for a case where the dust-to-total mass ratio $Z$ is less than a few $10^{-2}$, the streaming instability can drive only when the Stokes number is within the range of $10^{-2} \leq St_{CA} \leq 1$. We colored this region in the lower panels in Figures 3 and 4 (the yellow regions). Therefore, forming chondrite parent bodies via the streaming instability might be improbable.

Note, however, that the critical Stokes number for driving the streaming instability is not yet fully understood, especially in the small Stokes number region (Yang et al. 2016). Moreover, there are several alternative models for making dust clumps, e.g., dust trapping at the local pressure maxima (e.g., Drążkowska et al. 2013) and/or vortices generated by hydrodynamical instabilities (e.g., Meheut et al. 2012; Gibbons et al. 2015). Therefore, studies of these mechanisms taking fluffy growth and the reainment of chondrules into account are necessary.

6. Summary

Fluffy aggregation helps dust aggregates grow into planetesimals via direct aggregation. Therefore, the chondrite parent bodies might have also formed via the fluffy aggregation of millimeter-sized chondrules and submicron-sized matrix grains (Figure 1). However, chondrules can be ejected from the highly porous matrices when two chondrule-dust compound aggregates collide (Figure 2).

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