The Protostellar Jet Model of Chondrule Formation

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ABSTRACT

A chondrule formation theory is presented where the chondrule formation zone is located within 0.1 AU of the protosun. This hot, optically thick, inner zone of the solar accretion disk is coincident with the formation region of the protosolar jet. It is suggested, that chondrules are ablation droplets produced by the interaction of the jet wind with macroscopic bodies that stray into the jet formation region.

Provided these droplets are small enough, they will be swept up by the jet wind and subsequently cool at approximately the same rate as the expanding gas in the jet. There is a critical gas density ($\sim 10^{-11} \text{ g cm}^{-3}$) below which chondrules will undergo large, damped oscillations in altitude and thereby suffer reheating.

We claim that it is in the cooler, high altitude regions above the mid plane of the inner accretion disk that compound chondrules are formed, and the collisional fragmentation of chondrules take place. Since these processes take place in the jet flow, one can make a prediction for the expected structure of triple compound-chondrules. For such chondrules, it is suggested that two “relatively large” secondary chondrules will avoid each other. This prediction is valid only if the gas-flow is sufficiently laminar or if the “spin-down time” for a double compound chondrule is less than the inter-chondrule collision time.

The model assumes that particles, ranging in diameter from 1 \(\mu\text{m}\) to 1 cm, can be ejected from the inner-accretion disk by the jet flow, and that the angular momentum of this material is sufficient to eject it from the jet flow. Given these assumptions, any material so ejected, will fly across the face of the accretion disk at speeds greater than the escape velocity of the system. This material can only be recaptured through the action of gas drag. Such a capture process naturally produces aerodynamic size sorting of chondrules and chondrule fragments, while the ejection of refractory dust provides a possible explanation for the observed complementarity between matrix and chondrules.

This transfer of material will result in the loss of angular momentum from the upper atmosphere of the outer accretion disk and thereby facilitate the accretion of matter onto the protosun.

PREFACE

In this paper we discuss a relatively new chondrule formation model, which uses a hypothetical protosolar jet as the main chondrule formation mechanism. As this model is in the earliest stages of development, this particular monograph should not be treated as the forever, enduring, last word on the subject. Instead, we hope that our discussion will demonstrate the explanatory and predictive power of

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this model, and possibly prompt our colleagues to view chondrule formation from a different point of view as compared to other theories.

Some readers would argue that fewer, not more, chondrule formation theories are required. Certainly with nearly twenty different published theories, how can one determine which theory, if any, is the correct one? The answer, of course, is to embrace the theory that can not only explain many of the observed data, but also predict new facts. So in this study, we will attempt to explain:

1. the chondrule formation process;
2. the chondrule cooling rate;
3. the collisional fragmentation of chondrules and the formation of adhering compound chondrules, where the secondary (and smaller) chondrule was plastic at the time of collision;
4. chondrule reheating;
5. the chondrule upper size limit;
6. chondrule size sorting,

and
7. the complementary chemical composition of matrix and chondrules.

We will also discuss a possible link between protostellar jets and disk accretion, where a protostellar jet may inject material with low angular momentum into the upper atmosphere of an accretion disk, thereby enhancing stratified accretional flow, i.e., the upper layers of the disk accrete onto the protostar more readily than do the layers adjacent to the mid plane of the accretion disk.

As for predictions, we will describe the expected structure of triple compound chondrules, i.e., where two small secondaries are adhering to a large primary. We suggest that, for such a compound chondrule, the two secondary chondrules will tend to “avoid” each other. We give the mathematical expression for the minimum avoidance angle between the two secondaries and discuss under what circumstances this prediction will break down. We also show that the experimental data collected so far (Unfortunately, for only four such chondrules) have structures that are consistent with the predicted structure. It is our fond hope, that we will inspire some of our colleagues to collect much more data, so that the validity or otherwise of this prediction can be confidently determined.

1. JET ABLATION

In the late 1970’s, it was found that protostellar systems formed bipolar outflows of material, i.e., the gas flows travelling in two opposing directions, approximately perpendicular to the plane of the accretion disk (for a review, see Beckwith and Sargent 1993). These flows were detected via the CO rotation lines and consisted mostly of molecular material. The flow speeds of the bipolar outflows were found to be around 20 km s$^{-1}$, and by dividing their length by the observed flow speed, one could deduce a “dynamic” lifetime for these flows of around $10^5$ years. Later, unbiased surveys of protostellar systems suggests that all protostellar systems undergo some form of bipolar outflow stage (Fukui et al. 1993).
Often, but not always, one can find within a bipolar outflow a faster, more collimated flow known as a protostellar jet. Protostellar jets are usually detected in the Si I, OI and Hα lines and are consequently referred to as optical jets. They have observed wind speeds in the range of 100 - 400 km s\(^{-1}\), they eject a large amount of gas (total mass \(10^{-3} - 10^{-1} \, M_\odot\)), are quite energetic (total kinetic energy \(10^{44} - 10^{46}\) erg), long-lived phenomena (\(10^6 - 10^7\) years) that exist at the very earliest stages of star formation (Cabrit et al. 1990, Edwards, Ray and Mundt 1993). They occur not only in the massively active FU-Ori stages of star formation, but they also are to be found in the more quiescent, Classical T-Tauri Star (CTTS) stages. They appear to be generated by the interaction between protostars and accretion disks, within 10 stellar radii (\(R_*\)) of the protostar (Hartmann 1992).

What is the connection between bipolar outflows and protostellar jets? There is a growing consensus that bipolar outflows are a byproduct of protostellar jets (Snell et al. 1980, Shu et al. 1993), where the protostellar jet sweeps up ambient molecular cloud material into two thin shells, which manifest themselves as the observed bipolar lobes of CO emission. Once the molecular cloud material has been swept away (on a timescale of \(10^5\) years), the bipolar outflow disappears, leaving the protostellar jet to erratically fire away for a further \(10^6 - 10^7\) years.

Our interest in protostellar jets is sparked by the following mass processing argument. Suppose the solar nebula formed a protosolar jet. Such a jet would have ejected \(10^{-5} - 10^{-3} \, M_\odot\) of “rocky” material (i.e. all elements excluding H and He) from the solar nebula. If only 10% of this material were to fall back to the solar nebula then we would have \(10^{-6} - 10^{-4} \, M_\odot\) of rocky (possibly refractory) material being contributed to the solar nebula over a \(10^6 - 10^7\) year period. Given that the “rock” mass of the planets is of order \(10^{-4} \, M_\odot\), a protosolar jet may have made a significant contribution to the chemical structure of the solar nebula. Indeed, it is possible that ejecta from the jet flow may have been incorporated into the best preserved samples of the solar nebula: the chondritic meteorites.

Of course, such an argument does not prove that protostellar jets formed chondrules, but only provides a plausible basis for constructing a theory of chondrule formation. To create such a theory, our first task is to investigate the thermal environment of the jet formation region. This is a function of the formation distance of a protosolar jet from the protosun. A distance which is uncertain, but model fits to the Li 6707 Å, Fe I 4957 Å, Fe II 5018 Å lines in FU Ori, plus the lack of extinction and infrared excess in the wind, suggests a formation distance of \(< 10R_*\) (\(\approx 0.1\) AU for the solar nebula), where \(R_*\) is the radius of the protostar (Hartmann 1992). The temperature of the disk surface \((T_s)\) at such distances, is given (approximately) by the formula (Frank, King and Raine 1992)

\[
T_s(r) = 850 \, \text{K} \left( \frac{(M/M_\odot) \left( \frac{\dot{M}}{10^{-7} M_\odot \, \text{yr}^{-1}} \right)}{(r/0.1 \, \text{AU})^3} \right)^{1/4} \left(1 - \left(\frac{R_*}{r}\right)^{1/2}\right)^{1/4},
\]

where \(M\) is the mass of the protostar, \(\dot{M}\) is the mass accretion rate, and \(r\) is the radial distance from the protostar’s centre. The maximum temperature of the midplane of the accretion disk \((T_m)\) can be approximately determined from the formula

\[
T_m \approx T_s(\eta \tau)^{1/4},
\]

where \(\eta\) is some number of order 1, and \(\tau\) is the optical depth (ibid., and Cassen 1993). Since \(\tau\) may have values as high as \(10^4\), it is probable that at distances of order 0.05 AU from the protostar, the temperature at the midplane maybe high enough to melt or vaporise any rocky bodies that happen to
be in that section of the accretion disk. This raises the possibility that the protostellar jet winds may ablate droplets from the surface of rather warm “rock” bodies and if the droplets are small enough, they may be ejected from the inner accretion disk by the drag force of the jet flow.

It may be asked, however, why embrace the idea that chondrules are ablation droplets? We are encouraged to advance such a hypothesis, because chondrules appear to have been extruded or drawn from one or more extended magma bodies (Dodd and Teleky 1967). Also, an ablative process readily forms chondrule-like spheres from meteors (Brownlee et al. 1983). These “ablation-spheres” are chondrule-like in size and shape. They also share some physical similarities, e.g., they often contain relict grains. Despite this supportive evidence, there are, however, some significant problems to overcome before one can believe the ablation hypothesis.

First, we require the appropriately sized rocky-bodies to enter into the jet formation region. Protostellar systems generally have accretion disks and large bodies could be caught up in the inward accretional flow. However, the existence of chondrules in single ordinary chondrites that have (admittedly uncertain) $^{129}$Xe$^*/^{127}$Xe ages ranging over $10^7$ years (Swindle et al. 1991, Swindle and Podosek 1988), suggests that the accretional flow in the solar nebula was small enough to allow the survival of the chondrite precursors. Indeed, it is possible that such accretional inflow may have been dependent on height within the accretion disk, where only the gas at higher altitudes took part in the accretional flow (see §5).

Another mechanism for radial migration is that of gas drag due to the velocity difference between nebula gas and the near Keplerian motion of large bodies around the proto-sun. The infall velocity for material subject to this gas drag is (Whipple 1973, Weidenschilling 1977)

$$\frac{dr}{dt} \approx - \frac{2r\Delta V}{V_{Kep} t_D},$$

where $V_{Kep}$ is the Keplerian velocity, $\Delta V$ is ($\approx 10^{-3}V_{Kep}$) the difference between the gas angular velocity and $V_{Kep}$, while $t_D$ is the time scale for gas drag to influence the motion of the body. In the Stokes drag regime,

$$t_D = \frac{2r_{fod}^2 \rho_b}{9 \eta_g},$$

where $\rho_b$ is the mass density of the body and $\eta_g$ is the viscosity of the gas, and $r_{fod}$ is the radius of the “chondrule-fodder” body. Using the appropriate formula for the gas viscosity (e.g. Eq. (2.3.5) of Liffman 1992), one can obtain, from Eq. (3), the timescale for orbital decay ($\tau_{decay}$)

$$\tau_{decay} \approx 10^5 \left( \frac{r_{fod}/10^3 \text{ cm}}{(\rho_{fod}/1 \text{ g cm}^{-3})} \right) \left( \frac{(T_g/100)^{1/2} \text{ K}}{((T_g/100)^{1/2} \text{ K})} \right) \text{ yrs},$$

with $T_g$ being the temperature of the gas.

So, objects with a radius of around 10 m would fall into the Sun on a timescale of about $10^5$ years. This analysis, however, ignores the change in the size of the body due to accretion of dust from the solar nebula. If the dust to gas mass ratio is of $\sim 1$, accretion of material will increase the size of the body and eventually stop its infall into the Sun (Weidenschilling 1988). On the other hand fragmentation due to inter-body collisions would have been a source of smaller material which may have eventually reached the boundary layer between the proto-sun and the solar nebula.
There may have been other mechanisms that brought macroscopic material into the inner regions of the solar nebula. We will simply note that it is a plausible assumption and one we require for our model.

Once we have these large bodies in the jet-formation region, we are faced with the following scenario: the high speed jet-flow will probably occur at $z \geq$ the scale height of the accretion disk, $H$, since the outflow jet must be governed by the conservation of mass equation, and so high velocities will only occur when $\rho_g(z) \ll \rho_g(0)$. On the other hand, our “chondrule-fodder” bodies will be located on or near the midplane of the nebula ($z = 0$). This separation between the critical points in the wind and the chondrule-fodder bodies would appear to be the death-knell for our ablation hypothesis - for how else can one produce ablation droplets if the hypothetical windflow is nowhere near our chondrule-producing planetesimals?

To answer this, we examine protostellar jet theory.

Most protostellar outflow models assume that magnetic fields provide the main coupling mechanism between the accretion disk and the outflow (for recent reviews, see Bicknell 1992, Königl and Ruden 1993 plus Shu et al. 1993). In the most recent models, the magnetic driving mechanism involves the interaction between the dipole field of the protosun and the inner accretion disk (e.g. see Lovelace et al. 1991, and Shu et al. 1994). Studies of this protosun-nebula interaction can be traced back to the work of Freeman (1977), who assumed that the dipole field of the protosun was able to thread the inner (partially ionized) disk as shown in Fig. 1. In effect, the magnetic field of the protosun is “tied” to its surrounding accretion disk.

In Freeman’s model, the dipole field of the protosun rotates (approximately) with a rigid body velocity, so there exists a point from the protosun ($\sim 0.04$ AU), where the speed of the field sweeping over the disk equals the (approximately) Keplerian velocity of the disk. In Fig. 1, this position is denoted as the “synchronous orbit”. For distances $r$ greater than this synchronous distance, the magnetic field of the protosun has a greater angular velocity than the the accretion disk and so the magnetic field becomes “wrapped up”, i.e, the purely poloidal magnetic field of the protosun is converted into a toroidal field in the disk.

In Fig. 1, we show a side view of the resulting field structure. The toroidal field becomes the dominant field in the disk, and its direction reverses as one traverses the central plane. For such a configuration, the central plane of the nebula can become a “current sheet” where magnetic fields reconnect and the magnetic energy so released is converted into particle energy.

As discussed in Priest (1994) the particle flow velocity ($v_f$) obtained from the merging of magnetic fields ($B$) is close to the characteristic speed of a magnetic medium (i.e., the Alfvén speed $C_A$), which in mathematical notation has the form

$$v_f \approx C_A = \frac{B}{\sqrt{\mu_0 \rho_g}} = 3 \left( \frac{B}{100 \text{ G}} \right) \left( \frac{10^{-8} \text{ g cm}^{-3}}{\rho_g} \right)^{1/2} \text{ km s}^{-1} ,$$  \hspace{1cm} (6)

where $\rho_g$ is the mass density of the gas and $\mu_0$ is the permeability of free space. In the accretion disk, the direction of the reconnection flow would be roughly parallel to the central plane, with a small component in the $z$ direction.

To obtain ablative behaviour from gas flows with gas densities in the range of $10^{-6}$ to $10^{-8}$ g cm$^{-3}$, the required gas speeds range from 1 to 20 km s$^{-1}$ (Liffman 1992). Comparing this to Eq. (6) suggests
that these “reconnection” flows may have the required densities and speeds to ablate the planetesimals. Some of the ablation droplets so produced would then, presumably, be caught up in the flow field that eventually becomes the main protostellar jet flow.

It may turn out, however, that outflows are not powered by “wrapped-up” toroidal fields. Our purpose in describing this particular model is to demonstrate that, in the outflow region, the central plane of the accretion disk may be a highly active region, where conditions are conducive to chondrule formation by ablation.

2. CHONDRULE COOLING

A long standing problem in chondrule formation is how to explain the slow rate of chondrule cooling. Experimental simulation of chondrule formation suggests that chondrules cooled at a rate of 5 to 200 °C/hour (Hewins 1988). Such cooling rates are 3 to 5 orders of magnitude smaller than those expected for an isolated black body radiating heat directly into space. It would appear that chondrules were formed in a hot optically-thick medium or were produced in close proximity to other chondrules so that mutual radiation could damp the cooling rate.

In our model, chondrules are produced in the hot, optically-thick midplane of the inner accretion disk around a protostar. They are produced by the ablative interaction between the initial stage of a protostellar jet wind and rocky chondrule “fodder” bodies that happen to stray into the jet formation region. If the subsequent droplets are small enough, they will be swept up by the wind and begin to move with the jet flow.

As the jet flow moves away from the midplane of the accretion disk, it is likely that the gas in the flow will cool, since work is being done to expand the gas. Consequently, if the molten droplets are close to thermodynamic equilibrium with the surrounding gas flow, they must also cool at just about the same rate as the expanding gas flow.

Naturally, to model this process we require a model of a protostellar jet. Unfortunately, a comprehensive theory of protostellar jets is unavailable at this time. To partially circumvent this difficulty, we adopt a simple parameterised model of an outflow. First, we describe our system with cylindrical coordinates, the plane $z = 0$ being the midplane of the accretion disk with the protostar residing at the origin (see Fig. 11). With this coordinate system, the steady state form of the continuity or mass conservation equation ($\nabla \cdot (\rho_g v_g) = 0$) is easily solved, for an axisymmetric flow, to give the equation

$$2\bar{\rho}_g(z)\bar{v}_{gz}(z)\pi r(z)^2 = \text{constant} = \dot{M}_o,$$

where $\bar{\rho}_g(z)$ is the “r-averaged” gas mass-density as a function of $z$, $r(z)$ is the cylindrical radius of the outflow, $\bar{v}_{gz}(z)$ is the $z$ component of the “r-averaged” gas velocity, and $\dot{M}_o$ is the mass-loss rate of the outflow. The factor of two in Eq. (7) arises, because the protostellar jet is produced from both sides of the accretion disk.

To model the gas density and gas velocity profile within the disk, we assume that the disk constrains the radial size of the protostellar jet. If we now let

$$\bar{\rho}_g(z) \approx \bar{\rho}_g(z_1) \left(\frac{z_1}{z}\right)^m, \quad z \lesssim z_1, \quad m > 0,$$

where $\bar{\rho}_g(z_1)$ is the gas mass-density at the midplane, $z_1$ is the distance from the protostar to the midplane, and $m$ is a parameter that determines the rate of change of the gas density with radius.
where \( z_1 \) is the distance above the disk midplane where we can best define the flow variables. We can substitute Eq.(8) into Eq.(7) and obtain

\[
\bar{v}_{gz}(z) \approx \bar{v}_{gz}(z_1) \left( \frac{z}{z_1} \right)^m, \quad z \lesssim z_1, \quad m > 0. \tag{9}
\]

These solutions break down as \( z \to 0 \), this being the price we pay for not solving the momentum and energy equations. Equations (8) and (9) are only a first order fit to the mass conservation equations. However, they do, approximately, satisfy the density-velocity relationship observed in thermally driven winds, i.e., as \( \rho_g \) decreases, \( v_{gz} \) increases.

If the gas pressure \((P_g)\) can be globally modelled as a polytropic gas, then

\[
P_g = \kappa \rho_g^\gamma, \tag{10}
\]

where \( \kappa \) and \( \gamma \) are constants (\( \gamma \) is the ratio of the specific heats if the jet is adiabatic). Using the ideal gas law with Eqns (8) and (8), we can obtain an expression for the gas (and chondrule) temperature \((T_g)\) as a function of \( z \).

\[
T_g(z) \approx T_g(z_1) \left( \frac{z_1}{z} \right)^{m(\gamma-1)}. \tag{11}
\]

The gas flow and entrained chondrules decrease in temperature as they travel away from the central plane of the accretion disk.

If we assume that the particles are entrained with the flow, and they start from \( z = 0 \), then we can use Eqns (9) and (11) to give the gas and chondrule temperature as a function of time \((t)\). Solving for \( t \) in Eq. 9 gives

\[
\frac{z}{z_1} = \left[ \frac{t}{\tau_1} \right]^{1/(1-m)} \quad 0 < m < 1, \tag{12}
\]

where \( \tau_1 = z_1/((1-m)\bar{v}_{gz}(z_1)) \). Substituting the above equation into Eq.(11) gives,

\[
T_g(t) = T_g(z_1) \left[ \frac{\tau_1}{t} \right]^{m(\gamma-1)/(1-m)} \quad 0 < m < 1. \tag{13}
\]

The gas and chondrule cooling timescale is determined by the value of \( \tau_1 \) and \( T_g(z_1) \). As we shall show in the next section, for \( T_g(z_1) \sim 1200 \) K, \( z_1 \sim 10^4 - 10^5 \) km and \( \bar{v}_{gz}(z_1) \sim 0.1 - 1 \) km \( s^{-1} \). As a consequence, a minimum parameterization for \( \tau_1 \) is

\[
\tau_1 = 10^4 s \frac{(z_1/10^4 \) km\)}{(1 - m)(\bar{v}_{gz}(z_1)/1 \) km \( s^{-1}). \tag{14}
\]

The general decrease in gas temperature, and chondrule temperature, as a function of \( z \) within or near the accretion disk has important consequences for the collisional interactions between chondrules. Indeed, a simple analysis of this phenomenon leads to a model which can possibly explain the observed structure of compound chondrules and leads to a prediction for the physical structure of triple compound chondrules.
3. CHONDRULE COLLISIONS

3.1. Hover particles

Let us consider a chondrule that has just been ablated from its parent body. We assume, that the chondrule will be accelerated by the gas jet so that its motion is, initially, in the \(z\) direction. The equation of motion for the particle, parallel to the \(z\) axis, is given by

\[
m_p \ddot{z} = \frac{C_D}{2} \rho_g (v_{gz} - \dot{z})^2 \pi a_p^2 - \frac{G M m_p}{(z^2 + r^2)^{3/2}},
\]

(15)

where \(C_D\) is the coefficient of gas drag, while \(m_p\) and \(a_p\) are the mass and radius of the particle, respectively.

Suppose the particle reaches a state where the gas drag is balanced by gravity, so that \(\ddot{z} = 0\), and \(\dot{z} = 0\). Our equation of motion becomes (for \(z \ll r\))

\[
0 \approx \frac{C_D}{2} \rho_g v_{gz}^2 \pi a_p^2 - \frac{G M m_z}{r^3},
\]

(16)

where \(z_h\) is the value of \(z\) for the hovering particle. For \(z\) close to the midplane, we should expect that \(v_{gz}\) will be less than the sound speed, so if the mean free path of the gas (\(l\)) satisfies the relation

\[
l/2a_p > 10
\]

(17)

then \(C_D\) has the Epstein (1924) form

\[
C_D \approx \frac{8}{3v_{gz}} \sqrt{\frac{8kT_g}{\pi m_g}},
\]

(18)

where \(k\) is the Boltzmann constant and \(m_g\) is the mass of a gas particle (in this case we assume monatomic hydrogen). Using Eqns (16), (18) and (7) we find

\[
z_h \approx \left( \frac{8kT_g}{\pi m_g} \right)^{1/2} \frac{\dot{M}r}{2\pi GM \rho_p a_p}
\]

(19)

or

\[
z_h \approx 5 \times 10^4 \frac{(\dot{M}/10^{-8} \text{ M}_\odot/\text{ yr})(r/0.1 \text{ AU})(T_g/10^3 \text{ K})^{1/2}}{(M/\text{M}_\odot)(\rho_p/1 \text{ gm cm}^{-3})(a_p/0.1 \text{ cm})} \text{ km}.
\]

(20)

Thus, \(z_h \sim 10^{-4} \text{ AU} \ll 0.1 \text{ AU} \sim r\) as is required from the assumption leading to Eq. (16). Now, \(z_h\) is inversely proportional to the radius of the particle \((a_p)\), so smaller particles will have higher hover heights than larger particles of the same density. To reach their respective hover altitudes, the particles will move relative to each other and may undergo collision interaction. If the collisional velocities are small enough and the temperatures high enough, we may have chondrules fusing together to form compound chondrules.
3.2. Compound Chondrule Formation

Studies of compound chondrules have classified them into two general types: \textit{enveloping}, where one chondrule envelopes the other and \textit{adhering}, where one chondrule forms a “bump” on the other (Wasson 1993). In this paper, we restrict our attention to adhering compound chondrules.

Most adhering compound chondrules consist of a smaller chondrule stuck to a larger particle, where the smaller chondrule was plastic at the time of collision (Wasson \textit{et al.} 1995). This fact immediately poses a major problem for most chondrule formation theories which use simple “flash” heating scenarios. Since, if the chondrules were all formed in the one flash heating event, then we should expect the larger chondrules to be plastic at the time of collision.

In the Jet model, chondrules are continuously formed by an ablative process as chondrule fodder bodies stray into the jet formation region. As discussed in the derivation of Eq. (20), it is reasonable to assume that 0.1 cm particles will be supported by the jet flow at a $z$ distance of $10^{-4}$ to $10^{-3}$ AU above the midplane of the accretion disk. Let us suppose that at such distances, the temperature of the gas in the Jet stream is around 1200 K - the approximate solidus temperature of a chondrule.

Assuming a fairly steady jet flow, a hover particle in such a position should have sufficient time to equilibrate with the gas temperature and be relatively non-plastic. The hover particle will only see smaller particles flying past it in the jet flow, since larger particles of approximately the same mass density will have a lower hover altitude.

As these smaller particles fly past the hover particle, the thermal inertia of these particles will give them higher temperatures, and therefore greater plasticity than the hover particle. If the hover particle collides and fuses with these particles, then the observed structure of adhering compound chondrules will be obtained (for a schematic depiction of this process, see Fig. 2).

To give this idea a quantitative context, we note that the equation for the rate of temperature change of a spherical chondrule has the form

$$\frac{4}{3} \pi a_p^3 \rho_p C_V \frac{dT_S}{dt} = \Lambda 4 \pi a_p^2 q + 4 \pi a_p^2 \epsilon_a \sigma T_e^4 - 4 \pi a_p^2 \epsilon_e \sigma T_S^4,$$  \hspace{1cm} (21)

(Liffman 1992 and references therein).

The left hand side of the equation describes the time rate of change of the heat energy of the body, where $a_p$ is the radius of the particle, $\rho_p$ is the density of the particle, $C_V$ is the specific heat per mass of the body ($\sim 10^7$ erg g$^{-1}$ K$^{-1}$ for chondrules), $T_S$ is the surface temperature of the body, and $t$ the time. The first term on the right hand side describes the energy added to the body by the gas/body interaction, where $\Lambda$ is the heat transfer coefficient and $q$ is the gas/body heat transfer rate per unit surface area. The last two terms are radiation terms, where $T_e$ is the radiation temperature of the surrounding environment, $\epsilon_a/e$ is the absorptivity/emissivity of the surface (absorptivity and emissivity are assumed to be the same in this case), and $\sigma$ is the Stefan-Boltzmann constant. In the following discussion, we will assume local thermodynamic equilibrium applies and that $T_e = T_g$.

For chondrule formation, the size of the chondrules ($< 1$ cm) and the expected low gas density ($< 10^{-6}$ g/cc) indicate that the mean-free path of the gas is large relative to the size of the chondrules. For this “free molecular” flow regime $\Lambda \approx 1$ and $q$ has the form

$$q = \rho_g |v_g - v_p|(T_{rec} - T_g)C_H,$$  \hspace{1cm} (22)
Fig. 1.— The dipole field of the protosun interacts with the accretion disk, such that for distances greater than a critical distance, (≈ 0.04 AU) from the protosun, the poloidal field of the Sun is “wrapped-up” by the accretion disk. This process may cause the central plane of the inner accretion disk to become a “current sheet” produced by the merging magnetic field lines. This merging will be due to magnetic diffusion and the switch in the sign of the toroidal field as one traverses the central plane.

Fig. 2.— Schematic depiction of compound chondrule formation in a jet flow. The gas flow is depicted within the accretion disk, where the gas temperature decreases as \( z \) (the distance from the midplane of the disk) increases. A “large” hovering chondrule, with a temperature near or at the solidus temperature of 1200 K, collides with smaller chondrules that are entrained in the gas flow. Due to thermal inertia, the smaller chondrules are warmer and more plastic than the hover chondrule. In this way, one can obtain the observed phenomena that smaller chondrules were plastic at the time the compound chondrule was formed.
Thus, \( q \rho \) is the mass density of the gas, \( v_g \) is the gas velocity, \( v_p \) is the particle velocity, \( T_{\text{rec}} \) is the recovery temperature, and \( C_H \) is the heat transfer function for free molecular flow. \( T_{\text{rec}} \) and \( C_H \) have the forms

\[
T_{\text{rec}} = \frac{T_g}{\gamma + 1} \left[ 2\gamma + 2(\gamma - 1) s_{gp}^2 - \frac{\gamma - 1}{0.5 + s_{gp}^2 + s_{gp} \pi^{-0.5} \exp(-s_{gp}^2) / \text{erf}(s_{gp})} \right],
\]

(23)

and

\[
C_H = \left( \frac{\gamma + 1}{\gamma - 1} \right) \frac{k}{8 m_g s_{gp}^2} [\pi^{-0.5} s_{gp} \exp(-s_{gp}^2) + (0.5 + s_{gp}^2) \text{erf}(s_{gp})],
\]

(24)

\( \text{ibid.} \) where \( T_g \) is the gas temperature, \( m_g \) the mean gas particle mass \((\sim 1.66 \times 10^{-24} \text{g})\), \( \gamma \) is the ratio of specific heats, \( k \) is Boltzmann’s constant, \( \text{erf}(s) \) is the error function \( (\frac{2}{\sqrt{\pi}} \int_0^s \exp(-t^2) \text{d}t) \), and \( s_{gp} \) is the ratio of the relative streaming gas speed and the most probable Maxwellian gas speed,

\[
s_{gp} = \frac{|v_g - v_p|}{\sqrt{2kT_g/m_g}}.
\]

(25)

We are considering flows at the very base of the protostellar jet, so we should expect that \( s_{gp} \ll 1 \). Thus, \( q \) becomes

\[
q \approx \rho_g (T_g - T_s) \left( \frac{\gamma + 1}{\gamma - 1} \right) \sqrt{\frac{k^3 T_g}{8\pi m_g^3}}.
\]

(26)

or

\[
q \approx 5 \times 10^7 \left( \frac{\rho_g}{10^{-8} \text{ g cm}^{-3}} \right) \left( \frac{T_g}{10^3 \text{ K}} \right) \left( \frac{T_s}{10^3 \text{ K}} - \left( \frac{T_s}{10^3 \text{ K}} \right) \left( \frac{\gamma + 1}{\gamma - 1} \right) \sqrt{\frac{k}{10^3 \text{ K}}} \text{ergs}^{-1} \text{ cm}^{-2} \text{ s}^{-1}.
\]

(27)

When \( \rho_g \gtrsim 10^{-7} \text{ g cm}^{-3}, \gamma = 5/3 \) (monatomic gas), and \( T_g \approx 1200 \text{ K}, \) then the \( q \) term in Eq. (21) dominates the radiative terms, since

\[
\sigma T_g^4 \approx 5.7 \times 10^7 \left( \frac{T_g}{10^3 \text{ K}} \right)^4 \text{ergs}^{-1} \text{ cm}^{-2} \text{ s}^{-1}.
\]

(28)

So Eq.(21) becomes a first-order linear differential equation in \( T_s \), and has the solution

\[
T_s(t) \approx T_g + (T_s(0) - T_g) \exp(-t/\tau),
\]

(29)

where \( \tau \) has the form

\[
\tau = \frac{a_p \rho_p C_V}{3\Lambda \rho_g} \left( \frac{\gamma - 1}{\gamma + 1} \right) \sqrt{\frac{8\pi m_g^3}{k T_g}}.
\]

(30)

or

\[
\tau \approx 5 \left( \frac{a_p/0.1 \text{ cm}}{(\rho_p/3 \text{ g cm}^{-3}) (C_V/10^7 \text{ erg g}^{-1} \text{ K}^{-1}) (T_g/10^3 \text{ K})} \right) \left( \frac{T_g}{10^3 \text{ K}} \right) \text{s}.
\]

(31)

So, if we have \( \rho_g \gtrsim 10^{-7} \text{ g cm}^{-3} \) then particles with a radius between 0.01 cm and 0.1 cm will have a minimum temperature equilibration timescale of 0.5 to 5 s. Similar timescales are obtained for the low density case \((\rho_g \ll 10^{-7} \text{ g cm}^{-3})\), where the radiative terms in Eq.(21) dominate the “\( q \)” convective
term. In this case, Eq.(21) can only be solved numerically. The results from this calculation are shown in Fig. 3.

As can be seen, the e-folding timescales in the radiative case turn out to be about the same as those for the pure conduction case. One should note that these are minimum timescales, because we have neglected the latent heat of fusion and the diffusion time for heat to travel from the surface of the particle to its centre.

Despite these caveats, it is probable that the timescales for thermal equilibration are not much greater than as stated in Eq.(31). In such circumstances, we can only expect the smaller chondrule to be plastic at the time of collision if, in the collision zone, there is a sharp decline in gas temperature. We suggest that a temperature drop of 10 to 100 K s$^{-1}$ is required for compound chondrules to form. We also expect that the collision zone for compound chondrule formation must be quite thin, since the maximum time that the smaller secondary can remain plastic is $\sim 20$ s and the maximum collision velocity is $\sim 0.1$ km s$^{-1}$ (a higher collision velocity would fragment the chondrules, see Vedder and Gault (1974)). These numbers give a maximum length scale for the thickness of the compound chondrule formation region of around 1 km.

Now, suppose we have a chondrule hover-particle in a region of the flow where the ambient temperature of the gas flow is much less than the solidus temperature of the particle. If this solid hover-particle collides with smaller particles which are entrained in the jet flow, and if the relative speed of the two particles is greater than 0.1 km s$^{-1}$ then it is likely that one or both of the particles will undergo some damage from the collision. Depending on the relative speed of collision, this damage can range from slight chipping to complete fragmentation.

It is well known that many chondrules have undergone chipping and fragmentation. In the Jet model, one can provide a cause for this damage: collisions between chondrules. One can also give a site for where this damage will take place: in the jet flow at higher (and hence colder) altitudes than the compound chondrule formation region (see Fig. 11). In this way, the jet model can provide a unified scheme linking both compound chondrule formation and chondrule fragmentation.

Of course, like all theoretical mutterings in this field, the above results and ideas should be (and will be) treated with a reserved caution. A physical theory is of little use, unless it can be used to predict as well as to explain. So far we have attempted the latter, in the next section we will try the

![Fig. 3.— The radiative cooling of 3 particles of different radii from 1250 K to 1200 K.](image-url)
former. Let us throw caution to the winds (no pun intended) and predict the physical structure of triple compound chondrules.

### 3.3. The Weather Vane Effect

If compound chondrules were formed in protostellar jets, then gas drag would orient the compound chondrules immersed in the jet flow. This is simply because the spherical symmetry of the simple chondrule is destroyed. A secondary chondrule will act like the tail on a weather vane, orienting a binary compound chondrule such that the smaller secondary will be on the “downstream” side of the primary. A subsequent collision and fusion with another chondrule will produce a triple compound chondrule where the secondary chondrules are separated by a minimum avoidance angle. This situation is shown schematically in Fig 4.

![Fig. 4.— The “Weathervane” Effect. Secondary chondrules in a streaming gas flow, will tend to orient the compound chondrule such that the secondary is pointing in the direction of the flow. Any incoming chondrule will hit the rear end of the compound chondrule. Hence the prediction that secondary chondrules in a triple compound chondrule will tend to avoid each other.](image-url)

This simple scenario does have some initial complications, the major being that if the impact parameter (the offset distance between the centres of the particles, see Fig 4) between the colliding chondrules is greater than zero, then the binary compound chondrule will be set spinning. This rotational
motion will, however, be damped, because the secondary chondrule will be travelling half the time with the gas jet, and half the time against the jet. If the secondary chondrule has a cross-sectional area $A$, then the rate of change of the rotation rate, $\omega$, is given by

$$I \dot{\omega} \approx \frac{C_D}{2} \rho_g a_{ps} (v_{gz} - a_{ps} \omega)^2 A - \frac{C_D}{2} \rho_g a_{ps} (v_{gz} + a_{ps} \omega)^2 A,$$  \hspace{1cm} (32)$$

where $I$ is the moment of inertia of the system, and $a_{ps}$ is the distance between the primary and secondary chondrules. If we substitute the Epstein solution for $C_D$ (see Eq. 18), we obtain

$$\dot{\omega} \approx - \frac{8}{3} \sqrt{\frac{8kT_g}{\pi m_g}} \left( \frac{\rho_g a_{ps}^2 A}{I} \right) \omega,$$ \hspace{1cm} (33)$$

which has the solution

$$\omega \approx \omega_0 \exp(-t/\tau_{sd}),$$ \hspace{1cm} (34)$$

where $\omega_0$ is the initial rate of rotation, and $\tau_{sd}$ - the e-folding “spin-down time” - has the form

$$\tau_{sd} = \frac{3I}{8\rho_g Aa_{ps}^2 \sqrt{\pi m_g}} \sqrt{\frac{8kT_g}{8kT_g}}.$$ \hspace{1cm} (35)$$

To obtain an estimate for $\tau_{sd}$, we note that

$$I \sim (m_p + m_s)a_{ps}^2 \approx 2m_p a_{ps}^2,$$ \hspace{1cm} (36)$$

and

$$A \sim \pi a_s^2,$$ \hspace{1cm} (37)$$

where $m_p$ ($a_p$) and $m_s$ ($a_s$) are the masses (radii) of the primary and secondary chondrules, respectively. Combining Eqns (36) and (37) with the observation that $a_s \gtrsim 0.1a_p$ and putting it all into Eq.(35) gives

$$\tau_{sd} \lesssim 6.5 \times 10^3 \left( \frac{\rho_p/3 \text{ g cm}^{-3}}{\rho_g/10^{-8} \text{ g cm}^{-3}} \right) \left( \frac{a_p/0.1 \text{ cm}}{a_p} \right) \sqrt{\frac{10^3 \text{ K}}{T_g}} \text{ s}.$$ \hspace{1cm} (38)$$

So for a range in gas densities of $10^{-6}$ to $10^{-10}$ g cm$^{-3}$ the spin-down timescales for a binary compound chondrule range from around a minute to about a week.

Once a binary compound chondrule has stopped spinning, then it will take on a specific orientation, where the smaller secondary will be on the downstream side of the primary chondrule. As shown in Fig 5, the system can still oscillate between two extreme positions, but there will always be a section of the primary chondrule which will be shielded from colliding with particles that are entrained in the gas flow.

From Fig 5, one can deduce a minimum avoidance angle $\theta$ that two secondary chondrules on a triple compound chondrule should obey if the initial binary compound chondrule has stopped spinning. Given that the distance between the centres of the primary and secondary chondrules is $a_{ps} = a_p + f a_s$, where $f$ is the deformation factor of the secondary chondrule, then

$$\cos(\theta) = \frac{a_p - a_s}{a_p + fa_s} = \frac{1 - a_s/a_p}{1 + f a_s/a_p}.$$ \hspace{1cm} (39)$$
A plot of a solution to this equation is shown in Fig 6. If the weather vane effect did not occur then we would expect a uniform distribution of angles between secondaries for each value of \( a_s/a_p \). If the weather vane effect does occur, then the distribution should be skewed towards the 'allowed angles' shown in Fig 6. This is what we see in the few data we have been able to scour from the literature.

Again, we should make clear that the weather vane effect will only be observed if the timescale for a binary compound chondrule to stop spinning is less than the chondrule-chondrule collisional timescales. It will also only be observed in triple compound chondrules (i.e. one primary and two adhering secondaries). Compound chondrules that have two or more adhering secondaries will have unpredictable orientations and so an additional secondary may land anywhere on such a compound chondrule.

Finally, we should point out that we have ignored gas flow fluctuations that are parallel to the central plane of the accretion disk and that the derivation of Eq.(33) assumes \( v_{gz} \) remains constant during one revolution of the spinning compound chondrule. This latter assumption is probably true just after the formation of the compound chondrule, but may not be true as the compound chondrule stops spinning.

To see this, we note that a compound chondrule will stop spinning when it receives enough torque from the gas jet to remove all the angular momentum and reverse the chondrule’s initial spinning direction. The length of time the secondary will be moving against the gas jet is \( \pi/\omega \). The torque is approximately \( C_D \rho_g v_{gz}^2 A a_p/2 \). So the minimum rotation rate \( \omega_{min} \), is given by

\[
I \omega_{min} \simeq \frac{C_D}{2} \rho_g v_{gz}^2 A a_p \frac{\pi}{\omega_{min}}.
\]

Thus

\[
\omega_{min} \simeq \left( \frac{\pi C_D \rho_g v_{gz}^2 A a_p}{2I} \right)^{1/2}.
\]

Using the Epstein drag law (Eq.(18)) plus the approximate values for \( I \) and \( A \) (Eqs (36) and (37), respectively). One can show that

\[
\omega_{min} \approx 0.015 \left( \frac{a_s/0.1 \text{ cm}}{a_p/0.1 \text{ cm}} \right)^2 \left( \frac{v_{gz}/0.1 \text{ km s}^{-1}}{\rho_g/10^{-11} \text{ g cm}^{-3}} \right) \left( \frac{\rho_p/3 \text{ g cm}^{-3}}{m_g/m_H} \right)^{1/2}, \text{ rad Hz}
\]

where we have assumed that \( T_g = 1000 \text{ K} \) and that \( m_g = m_H \). So the period of rotation \( (2\pi/\omega) \) of this slowly spinning compound chondrule is approximately 400 seconds. In a 0.1 km s\(^{-1}\) gas flow, this gives a scale length \( (L) \) of order 10 to 100 km.

So the relevant Reynolds number \( (Re) \) of the flow (see Liffman 1992) is given by

\[
Re \approx 8.7 \times 10^2 \left( L/10^3 \text{ km} \right) \left( v_{gz}/0.1 \text{ km s}^{-1} \right) \left( \rho_g/10^{-11} \text{ g cm}^{-3} \right) \sqrt{(T_g/10^3 \text{ K})(m_H/m_g)},
\]

so, with \( Re \) ranging from 100 to 1,000, \( v_{gz} \) may undergo fluctuations which may modify the spin-down time of Eq.(35).

### 3.4. Chondrule Reheating

If we combine the theory that we have acquired from §2 & §3, we can present a scenario for the phenomenon of chondrule reheating, where high temperature rims are observed around chondrule cores.
Fig. 5.— The extreme positions of the secondary chondrule while oscillating from side to side.

Fig. 6.— Comparison of theory vs observations for triple compound chondrules. We plot the angle between the centres of secondary chondrules vs the radius ratio of the largest secondary to the primary. The largest secondaries were used for $a_s$ values, as they would show ‘disallowed’ chondrules, if any were present. As can be seen, all the observed angles lie in the ‘allowed’ zone. The line separating the ‘forbidden’ and ‘allowed’ zones is given by Eq.(39).
(Kring 1991). We suggest that reheating may be a consequence of a chondrule overshooting its hover altitude and then oscillating, in a damped manner, around the hover altitude. Each time the chondrule decreases its altitude, it will undergo an increase in temperature. If the altitude decrease is large enough, the corresponding increase in temperature will be sufficient to remelt the chondrule.

The actual mechanism whereby these high temperature rims are formed will not be discussed in any depth here. We only suggest that as the chondrule oscillates around its hover point, it accretes dust-grains/molten droplets and that this material forms the foundation of the rim once the chondrule is reheated.

To understand how these oscillations arise, we note that (see also §3.1) the equation of motion for a particle, parallel to the $z$ axis, is given by

$$m_p \ddot{z} = \frac{C_D}{2} \rho_g (v_{gz} - \dot{z})^2 \pi a_p^2 - \frac{GM m_p \dot{z}}{(z^2 + r^2)^{3/2}}. \quad (44)$$

As noted in §3.1, the drag coefficient ($C_D$) is likely to have the Epstein form

$$C_D \approx \frac{8}{3 v_{gz}} \sqrt{\frac{8 k T_g}{\pi m_g}}. \quad (45)$$

Combining Eqs (44) and (45) and assuming $z \ll r$, we have:

$$m_p \ddot{z} + \frac{4}{3} \left( \frac{8 \pi k T_g}{m_g} \right)^{1/2} \rho_g \pi a_p^2 \dot{z} + \frac{GM m_p \dot{z}}{r^3} \approx \frac{GM m_p z_h}{r^3}, \quad (46)$$

where $z_h$ is the hover height given by Eq.(19).

Let $x = z - z_h$, then Eq.(46) becomes

$$m_p \ddot{x} + c \dot{x} + k x = 0, \quad (47)$$

where

$$c = \frac{4}{3} \pi a_p^2 \rho_g \left( \frac{8 \pi k T_g}{m_g} \right)^{1/2}, \quad (48)$$

and

$$k = \frac{GM m_p}{r^3}. \quad (49)$$

If we assume that $c$ and $k$ are approximately constant then $c^2 < 4 km_p$ gives solutions of Eq.(49) that oscillate, with decreasing amplitude, as a function of time ($t$). If $c^2 = 4 km_p$ or $c^2 > 4 km_p$ then the solutions of Eq. (49) are critically or strongly damped and the particle approaches its hover height without oscillation.

The case $c^2 = 4 km_p$ translates into an equation which gives the critical density for particle oscillation ($\rho_{gc}$):

$$\rho_{gc} = 2.7 \times 10^{-12} \frac{(a_p/0.1 \text{ cm}) (\rho_p/1 \text{ gm cm}^{-3}) (M/M\odot)^{1/2} (m_g/m_H)^{1/2}}{(r/0.1 \text{ AU})^{3/2} (T/10^3 \text{ K})^{1/2}} \text{ gcm}^{-3}. \quad (50)$$
If $\rho_g < \rho_{gc}$ then $c^2 < 4km_p$ and the particle will oscillate around the hover point $z_h$. It is our contention that it is this oscillation process that causes chondrule reheating. Since the temperature of the jet flow increases as $z$ decreases, thus as the particle travels to lower altitudes its temperature must increase.

To be specific, when $\rho_g < \rho_{gc}$ then Eq.(47) admits the solution

$$z = z_h + \frac{\dot{z}(z_h)}{\mu} e^{-bt} \sin \mu t,$$

(51)

where $\dot{z}(z_h)$ is the $z$ speed of the chondrule when it first reaches $z = z_h$,

$$\mu^2 = \frac{4m_p k - c^2}{4m_p^2},$$

(52)

and

$$b = \frac{c}{2m_p}.$$  

(53)

Thus the period of the oscillation is

$$\frac{2\pi}{\mu} \approx 11.6 \left(\frac{r/0.1 \text{ AU}}{(M_{\odot})^{1/2}}\right)^{3/2}, \text{ days}$$

(54)

while the damping timescale is given by

$$\frac{1}{b} = 5 \left(\frac{a_p/0.1 \text{ cm}}{(\rho_g/10^{-12} \text{ gm cm}^{-3})} \left(\frac{m_g/m_H}{(T_g/10^3 \text{ K})^{1/2}}\right)\left(\frac{\bar{v}_g(z)/0.1 \text{ km s}^{-1}}{(r/0.1 \text{ AU})^{2/3}}\right)\right), \text{ days}$$

(55)

and the amplitude of the oscillation has the form

$$\frac{\dot{z}(z_h)}{\mu} = 1.6 \times 10^4 \left(\frac{\dot{z}(z_h)/0.1 \text{ km s}^{-1}}{(r/0.1 \text{ AU})^{3/2}}\right) \left(\frac{(\bar{v}_g(z)/0.1 \text{ km s}^{-1})}{(M_{\odot})^{1/2}}\right), \text{ km}$$

(56)

Comparing Eqs (56) and (20) shows that if $\dot{z}(z_h)$ is comparable to the gas flow speed, then the amplitude of the oscillations can be a significant fraction of the hover height. The corresponding temperature variations may also be significant.

Of course, from Eq.(48), $c$ is a function of $\rho_g$ and $T_g$ and cannot be considered a constant, so the above analysis should only be treated as an informative approximation to the complete system. A slightly more realistic analysis can be obtained by constructing a computer simulation of the flow. We were able to accomplish this by assuming a $z$ velocity for the outflow wind of the form

$$\bar{v}_g(z) \approx 0.1 \left(\frac{z}{z_h}\right)^m \text{ km s}^{-1}, \quad z \approx z_h,$$

(57)

with $m = 0.1$, with the orbital (angular) velocity of the gas assumed to be Keplerian. The density structure of the gas flow could then be easily deduced from Eqs (7) and (57), and had the form

$$\bar{\rho}_g(z) \approx 4.5 \times 10^{-12} \left(\frac{M/10^{-8} \text{ M}_\odot/ \text{ yr}}{(\bar{v}_g(z)/0.1 \text{ km s}^{-1})(r/0.1 \text{ AU})^2}\right) \text{ g cm}^{-3},$$

(58)
while the temperature of the gas was given by

\[ T_g(z) \approx 1200 \left( \frac{zh}{z} \right)^{m(\gamma - 1)} \text{K}, \]

where we set \( \gamma = 5/3 \).

As can be seen from the above formula, \( T_g(z) \) will be greater than 1900 K when \( z \lesssim 0.001zh \). In such cases, we simply set \( T_g(z) = 1900 \text{K} \). Finally, to compute the temperature of the particle, we did a full time integration of Eq.(21).

In Fig. 7a we show the motion of a chondrule-like particle \((a_p = 0.1 \text{ cm}, \rho_p = 3.8 \text{ g cm}^{-3})\) released from very close to the midplane \((z = 100 \text{ cm})\) into a jet flow with a total mass loss rate of \(10^{-9} \text{M}_\odot \text{ yr}^{-1}\).

The particle was released at a distance \((r)\) of 0.05 AU from the center of a solar mass protostar, where 0.05 AU is also the assumed radius of the protostellar jet. The particle underwent a series of damped oscillations the period of which was slightly less than 6 days, while the damping timescale was approximately 30 days. The corresponding temperature of the particle is shown in Fig. 7b. The particle temperature has a damped “saw-tooth” pattern. The maximum temperature of the resulting temperature peaks being approximately 1420 K. Of course higher temperature values can be obtained if the temperature gradient is steeper than what we have assumed \((e.g., \text{ if } m > 0.1)\).

While this oscillating motion is of interest, the particles can also undergo a critically damped trajectory. In particular, if \( \rho_g > \rho_{gc} \) then Eq. (47) gives the solution

\[ z = z_h + \frac{\dot{z}(z_h)}{2\sqrt{-\mu^2}} (e^{(-b+\sqrt{-\mu^2})t} - e^{(-b-\sqrt{-\mu^2})t}), \]

The trajectory and temperature of such a particle are shown in Figs 8a & b, where now the mass loss rate of jet is assumed to be \(2 \times 10^{-8} \text{M}_\odot \text{ yr}^{-1}\).

From this simple model, we would suggest that chondrule reheating occurs when the Jet flow is declining in mass flux, \(i.e.,\) during the later stages of a CTTS’s evolution.

4. THE CHONDRULE SIZE LIMIT

Why do chondrules have an upper size limit? Possible solutions such as size-limited precursor dust balls to aerodynamic sorting have been suggested. In the Jet model, one has to eject particles from the inner accretion disk and ensure that they arrive in the outer parts of the accretion disk. Clearly, there must be some size limit to this process and this size limit must be dependent on the “strength” of the jet flow, which in turn is dependent on the gravitational potential of the inner accretion disk. To derive a quantitative relationship between chondrule and accretion disk, we must first describe the behaviour of a droplet in a flow.

Suppose we have a droplet of molten material that is subject to a streaming gas flow. The surface drag of the flow will tend to “rip” the droplet apart, while the surface tension of the melt will try to minimise the exposed surface area of the droplet and keep the droplet together. The balance between these two conflicting forces produces a stable droplet of maximum radius \(a_p\), the formula for which is

\[ a_p \approx \frac{\gamma We_0}{CD\rho_g v_g}, \]
Fig. 7.— (a) A chondrule-like particle is released into a protostellar jet flow at a distance of 0.05 AU from a protostar. The mass loss rate of the flow is set at $10^{-9} \, M_\odot \, yr^{-1}$. Such a low mass loss rate ensures the gas density is lower than the critical density given in Eq. (50). In such a circumstance the particle will oscillate, in a damped manner, around the hover point (see Eq. (51)). (b) The temperature of a particle as it undergoes the trajectory shown in Fig. 7a.

Fig. 8.— (a) A chondrule-like particle is released into a protostellar jet flow at a distance of 0.05 AU from a protostar. The mass loss rate of the flow is set at $2 \times 10^{-8} \, M_\odot \, yr^{-1}$. Such a low mass loss rate ensures the gas density is higher than the critical density given in Eq. (50). In such a circumstance the particle will approach the hover point in a critically damped manner (see Eq. (60)). (b) The temperature of a particle as it undergoes the trajectory shown in Fig. 8a.
where $\gamma$ is the surface tension of the molten material, and $We_0$ is a dimensionless factor called the “critical Weber number”. $We_0$ accounts for the non-uniformity of the gas drag pressure over the surface of the droplet. Typically, $We_0 \approx 10$ (Bronshten 1983).

The streaming gas flow will not only determine the stable size of the droplet, but it will also subject the particle to a drag force given by $C_D \rho_g v_{gp}^2 A_p$, where $A_p$ is the cross sectional area of the particle that is facing the gas flow.

One can show (Liffman and Brown 1995) that the work done, $W$, by the wind - flow in ejecting a particle is

$$W = \frac{C_D^*}{2} A_p < \rho_g v_{gp}^2 > L, \quad (62)$$

where $C_D^*$ is a “representative” value of the drag coefficient during the propulsion phase, $< \rho_g v_{gp}^2 >$ is the mean value of $\rho_g v_{gp}^2$, and $L$ represents the actual length of the propulsion stage.

For a particle to escape the disk, we require that

$$< \rho_g v_{gp}^2 > = \frac{kGMm_p}{RC_D^* A_p L}, \quad (63)$$

where $k$ is $\geq 1$, $G$ is the gravitational constant, $m_p$ is the mass of the particle, and $R$ is the initial radial (or semi-major axis) distance of the particle from the protostar. Substituting Eq.(63) into Eq.(61) gives

$$a_p \approx \frac{We_0 C_D^* \gamma}{kC_D(0)GMm_p} R A_p L, \quad (64)$$

where $C_D(0)$ is the value of the gas drag coefficient at the position where the particle is formed. If we assume that our particle is spherical (as are most unfragmented chondrules) and that our propulsion distance $L$ is proportional to the height of the disk at a distance $R$ away from the protostar, i.e. $L = \lambda R^\beta$, then

$$a_p \approx \frac{3\lambda We_0 C_D^* \gamma R^{\lambda+\beta}}{4C_D(0)kGM\rho_p} \frac{1}{1/2}. \quad (65)$$

Now $C_D$, generally, decreases with increasing gas-flow speed. So, we should expect that $C_D^*/C_D(0) \lesssim 1$, since $C_D^*$ is an ‘average’ gas drag coefficient over the entire propulsion stage, while $C_D(0)$ samples the gas flow at the beginning of the propulsion stage, where the gas flow is probably at its slowest. Setting $\lambda \approx 0.01, \beta \approx 1$, (typical, approximate values for the disk height), $k \approx 1, M \approx M_\odot$, and $We_0 \approx 10$, we obtain

$$a_p (\text{cm}) \lesssim 0.4 \left[ \frac{\gamma}{\rho_p} \right]^{1/2} R (\text{AU}). \quad (66)$$

The values for the surface tension, $\gamma$, and the density, $\rho_p$, are material dependent. We are interested in Fe-Ni and silicate chondrules, so we shall use the surface tensions of meteoric iron and stone, which are $\gamma_{\text{iron}} = 1,200$ and $\gamma_{\text{stone}} = 360$ g s$^{-2}$ (Allen et al. 1965) with corresponding mass densities of $\rho_{\text{iron}} = 7.8$ and $\rho_{\text{stone}} = 3.4$ g cm$^{-3}$. Substituting these values into Eq.(66) gives us the approximate radius of iron and silicate droplets as a function of distance from the protostar:

$$r_{\text{iron}} (\text{cm}) \lesssim 5R (\text{AU}) \quad \text{and} \quad r_{\text{stone}} (\text{cm}) \lesssim 4R (\text{AU}). \quad (67)$$
If chondrules were formed from this ablative process, then to obtain the observed maximum chondrule sizes we require that $R \lesssim 0.1 \text{ AU}$ (and $k \geq 1$). This is consistent with the conclusion that protostellar jets are formed in an accretion disk at a distance of 0.05 to 0.1 AU from a solar-type protostar (Camenzind 1990, Hartmann 1992).

Although this result is encouraging, there is a contradiction between the ideas presented in §3 and §4. In §3, we require the chondrules to remain stationary in the flow so to allow the formation of compound chondrules, while in this section we have shown that the chondrule size limit is determined by the “power” of the jet flow to eject particles from the inner accretion disk. This contradiction would appear to severely limit, perhaps destroy, our theory. After all, how can one require particles to be stationary relative to the accretion disk and also expect them to be ejected at speeds $\geq$ the escape speed of the protostellar system?

The only way out of this contradiction is for the jet flow to be highly variable in both density and velocity. Compound chondrules would presumably form when the flow is relatively quiescent, while chondrule ejection would perhaps occur when there is a major increase in the density and/or velocity of the flow. Observations do suggest that jet flows are highly variable in their behaviour (Edwards et al. 1993). For example, Mundt (1984) obtained observational evidence of variations in young stellar winds that vary on time scales of months. Despite this observational support, the resolution of this problem must await a coherent theory of protostellar jet formation, which in turn must await high resolution observations of protostellar jets.

Besides determining a maximum size limit for chondrules, the ejection of chondrules by a protostellar jet also, indirectly, causes chondrules and chondrule fragments to be size sorted.

5. SIZE SORTING AND DISK ACCRETION

It has been known for many years (Dodd 1976) that chondrules, both silicate and metal, are size-sorted in meteorites, i.e., particles in a particular chondrite satisfy the relation (Skinner and Leenhouts 1993)

$$\rho g a_p \approx \text{constant}.$$  \hspace{1cm} (68)

It has recently become apparent that size-sorting also applies to fragments of chondrules (Skinner and Leenhouts 1991). This shows that chondrites act as “size-bins”. Chondrules and their fragments were formed, mixed and later sorted by size within the chondrite forming regions of the solar nebula.

The usual explanation for this phenomenon (e.g. Dodd 1976 and references therein), is some form of aerodynamic drag effect, where larger and/or denser particles can travel further into a resistive medium than can smaller and/or less dense particles. We too shall use this idea, for it arises naturally from the Jet model.

To see this, suppose that a particle is ejected from the inner accretion disk at speeds close to or exceeding the escape velocity. Suppose further that the angular momentum of the particle is high enough to eject the particle from the gas flow and allow it to move across the face of the accretion disk. If the particle is not subject to gas drag, from the upper-atmosphere of the accretion disk, it will simply move out of the system and into interstellar space. If the particle is subject to gas drag then, given sufficient drag, the particle will be recaptured by the protostellar system and may fall into the outer parts of the accretion disk.
A stream of such particles will be size-sorted. Since the smaller, less dense particles will fall close to the protostar, while the larger, more dense particles will fall further away from the protostar. Thus, we should expect that meteorites from the outer parts of the solar system will contain larger particles than meteorites from the inner solar system, and indeed, chondrules from the carbonaceous chondrites are (usually) larger than chondrules from ordinary chondrites.

There is an additional consequence of this model which has to do with angular momentum transfer. A particle in Keplerian orbit around a star has an angular momentum that is proportional to the square root of the distance between the particle and the star. Thus, a particle that is ejected from the inner accretion disk, and subsequently stopped by gas drag, will simply fall back into its original orbit unless angular momentum is transferred from the disk to the particle. If gas drag allows the disk to transfer angular momentum to the particle, then the particle can fall to the outer parts of the disk.

Clearly, if the disk gas gives up angular momentum, then it must move in towards the protostar. Since protostellar jets are fueled by disk accretion, one obtains the schematic picture of a jet flinging out material into the disk such that the disk will accrete onto the protostar and refuel the jet. Heuristically, one can think of chondrules and other associated jet ejecta as delayed protostellar jet fuel.

To turn this qualitative speculation into quantitative speculation, we have to model the particle ejection process. Unfortunately, this is a difficult thing to do, since there is no consensus on how protostellar jets work. So, we make the following assumptions:

1. We suppose that our chondrule particles are, initially, in a circular Keplerian orbit of radius $R \leq 0.1$ AU from the protostar.

2. We assume that the protostellar jet gives the particles an initial “boost” velocity $\dot{z}(0)$ that is comparable to the protostar’s escape velocity at that point. Our tentative justification for this assumption is that protostellar jets are observed to have speeds comparable to the maximum escape velocity of a protostellar system. So, we presume that particles initially entrained in such a flow may also obtain similar speeds.

3. Finally we assume that the particles have, initially, no radial velocity, and that the self-gravity of the disk is negligible compared to the gravity of the protostar. The former of these two assumptions comes from the observation that jet flows tend to be perpendicular to their respective accretion disks. Protostellar jets do have a nonzero radial velocity, but we have, for simplicity, ignored this component. The validity or otherwise of these ideas is discussed at some length in Liffman and Brown (1995).

Given these assumptions, the equations of motion for a particle become

\[ \ddot{r} = \frac{h^2}{r^3} - \frac{GMr}{[r^2 + z^2]^{3/2}} , \]

(69)

\[ v_\theta = r\dot{\theta} = \frac{h}{r} , \]

(70)

and

\[ \ddot{z} = -\frac{GMz}{[r^2 + z^2]^{3/2}} , \]

(71)

where $h$ is the specific angular momentum of the particle and has the value

\[ h = \sqrt{GMR} . \]

(72)
The value of $h$ is a constant, since it is assumed that there are no external torques, parallel to the $z$ axis, acting on the system.

These equations are not difficult to model numerically and can be solved by standard techniques. Although a brief analysis of the above equations shows that if the particles are travelling at speed greater than the escape velocity of the system, then the initially vertical path of the projectile will quickly turn into a “horizontal” path across the face of the accretion disk. A computer simulation of this phenomenon is given in Fig. 9 (see also Liffman and Brown (1995)).

In Fig. 9, particles are ejected at $r = 0.04$ AU with different vertical velocities ranging from 149 to 174 km s$^{-1}$. The particles are assumed to move out of the gas outflow at $r = 0.1$ AU, whereupon they encounter the gas halo of the accretion disk and their subsequent motion is governed by gas drag plus the gravitational force from the protostar.

As the particles move through the halo gas of the accretion disk, they will acquire, by gas drag, angular momentum from the disk. If we assume that the centrifugal force of the halo gas balances the radial component of the protostar's gravity, then the angular speed of the halo gas, $v_{\theta,\text{gas}}$, will be given by

$$v_{\theta,\text{gas}} = \frac{r \sqrt{GM}}{(r^2 + z^2)^{3/4}}$$

(73)

Once the particle has come to rest, relative to the halo gas, it will have the angular velocity given by Eq.(73) and a specific angular momentum, $h$, given by $h = rv_{\theta,\text{gas}}$. As can be deduced from Eq.(69), such a specific angular momentum implies that $\ddot{r} = 0$, and the only force acting on the particle will be the $z$ component of the gravitational force, which will point towards the accretion disk. As the particle moves towards the accretion disk, the gas density and angular velocity of the gas will increase, thereby keeping $\ddot{r} \approx 0$. As can be seen from Fig. 9, the subsequent path of the particle is roughly parallel to the $z$ axis.

The paths of the captured particles, shown in Fig. 9, can be approximated to that shown in Fig. 10, where the ascending path length, $l$, is given by $l = K\chi$, with $K$ being a number in the range of $4 \pm 1$, (Liffman and Brown 1995) and $\chi$ is the “stopping distance” as defined by the equation,

$$\chi = \frac{4a_p\rho_p}{3\rho_g}.$$  

(74)

As a justification for the above equation, we note that a macroscopic particle will come to rest when it has encountered a total gas mass approximately equal to its own mass. The “stopping distance”, so defined by this prescription, is easily shown to be the $\chi$ length scale as given in the above equation. So, for a constant gas density, the larger or more dense a particle is, the further it will be able to travel before it comes to rest.

The range of a recaptured projectile is simply given by the formula

$$r \approx l \cos(\theta),$$

(75)

where $\theta$ is the angle between the ascending path of the particle and the midplane of the accretion disk.

The simplicity of Eq.(75) suggests that aerodynamic size sorting may occur to particles that are recaptured by the protostellar system. To see this, we suppose that we have two particles with different
Fig. 9.— Simulation data showing chondrule-like particles travelling above the disk and falling back into it. A solar-mass protostar is located at \( r = 0, z = 0 \). Surrounding the protostar is an accretion disk, the scale height of which is shown in profile. Chondrule-like particles (radius = 0.1 cm, density = 3.5 g cm\(^{-3}\)) are given a velocity boost in the \( z \) direction, the magnitudes of which (in km s\(^{-1}\)) are shown next to the trajectories. Particles that are subject to sufficiently high gas drag are later recaptured.

Fig. 10.— An approximation of chondrule trajectories as a triangular path.
mass densities: \( \rho_1 \) and \( \rho_2 \). Suppose further that these two particles fall back to the accretion disk at the same distance from the protostar. Then we can write the radius ratio of particle 1 to particle 2 as

\[
\frac{a_1}{a_2} \approx \frac{\rho_2}{\rho_1} Q(r),
\]

(76)

where

\[
Q(r) = \frac{K_2 \rho g_1 \cos(\theta_2)}{K_1 \rho g_2 \cos(\theta_1)}.
\]

(77)

The factor \( Q(r) \) is dependent on the initial ejection speeds of the particles (through \( K \) and \( \theta \)), the initial distance of the particles from the protostar (again \( K \) and \( \theta \)) and the scale height of the accretion disk (through \( \rho_g \)). If the particles are created at about the same distance from the protostar then we would expect that \( K, \theta \) and \( \rho_g \) would be similar for both particles, since the average flight path of both particles would be about the same. This would imply that \( Q(r) \approx 1 \). So, from the collection of particles that fall to the accretion disk at \( r \), the denser particles should have smaller radii.

As has been discussed, size sorting is density dependent in that dense Fe-Ni chondrules are always smaller than less dense silicate chondrules. The mass densities of the two types of chondrules are (Skinner and Leenhouts 1993) \( \rho_{Si} \approx 3.8 \, \text{g cm}^{-3} \), and \( \rho_{Fe} \approx 7.8 \, \text{g cm}^{-3} \), which implies that \( \rho_{Si}/\rho_{Fe} \approx 0.5 \), and so if Eq.(69) is applicable to chondrules, we should expect that

\[
\frac{a_{Fe}}{a_{Si}} \approx 0.5 Q(r).
\]

(78)

Using published data for the meteorite types H, L, LL, (Dodd, 1976) and CR (Skinner and Leenhouts 1993), we can compute the average Fe to Si size ratio for all these different types of meteorites (excluding Bjurböle)

\[
\langle \frac{a_{Fe}}{a_{Si}} \rangle_{H,L,LL,CR} = 0.52 \pm 0.16.
\]

(79)

The corresponding approximate mean \( Q \) values is

\[
\langle Q_{H,L,LL,CR} \rangle = 1.04 \pm 0.32.
\]

(80)

Of course, the agreement between the above \( Q \) value and our theoretical model should be treated with caution, since we have only presented the bare beginnings of a quantitative model. Nonetheless, it does illustrate the potential of the “Jet” model to explain the phenomenon of size-sorting.

Finally, we return to our discussion of angular momentum transfer and disk accretion. Let us consider a ring of material in an accretion disk at a distance \( r \) from a protostar. The angular momentum of this material is

\[
L(r) = m(r) \sqrt{GMr},
\]

(81)

where \( m(r) \) is the mass of the ring of material at \( r \). Now suppose that material falls onto the accretion disk, and this infalling material has essentially zero angular momentum. In such a case, the angular momentum of the ring is conserved, and the mass of the ring becomes a function of time, i.e., \( m \equiv m(r,t) \). For such a case, we can differentiate Eq. (81) to obtain

\[
\frac{dr}{dt} = -\left(\frac{2r}{m}\right) \frac{dm}{dt}.
\]

(82)
Thus, if infalling material doubles the mass of the ring, it will move from its initial position $r(0)$ to $r(0)/4$. It is via this mechanism of mass and angular momentum transfer that disk accretion may be, in part, mediated. Indeed, because it is the halo gas of the accretion disk that will be transferring most of the angular momentum to the “Jet projectiles”, it is possible that the upper layers of the accretion disk are the ones that undergo most of the accretion, leaving the midplane relatively untouched.

This type of process may explain an implicit contradiction between observations and meteoritics. Observations suggest that protostars keep on accreting material from their disk for periods of up to $10^7$ years (Cabrit et al. 1990). Radiometric data from the decay of $^{129}$I and $^{26}$Mg suggest that meteorites accreted material for periods of order $10^6 - 10^7$ years. How could the meteoritic material have been preserved if a major portion of the solar nebula was accreted onto the protosun? The answer, we suggest, is that disk accretion was, in part, altitude dependent. Material flung from the protosolar jet mediated the angular momentum transfer and one component of this mass transfer was the chondrule.

### 6. CHONDRULE-MATRIX COMPLEMENTARITY

In this paper, we claim that chondrules have the physical characteristics expected of ablation droplets that have been formed and ejected by a protosolar jet, and then recaptured by the solar nebula through the action of gas drag.

Such a model, however, is immediately confronted with the complementary composition of matrix and non-matrix material in meteorites. For example, Wood (1985) discusses the case of Murchison, where the matrix has an Fe/Si ratio of 1.23, while the non-matrix material (chondrules, CAIs, isolated crystals) is $\sim 0.2$. These two dissimilar components combine to give an Fe/Si ratio of 0.81, which is close to the solar value of 0.9. This is unlikely to be accidental, and is clear evidence for the local formation of chondrules, and the refractory component of chondritic material.

Before we throw out our wind-transport model, however, one should note that chondrules, at least in this model, form in a relatively small region of the solar nebula, i.e., in or near the boundary layer of the protostellar system. If the chondrules have a low Fe/Si value then the surrounding material will, by mass balance, have a high Fe/Si ratio. If this latter material comes in a non-gaseous form (e.g. 10 $\mu$m dust, CAI &c) then it too will be ejected with the chondrules. One will obtain the desired mass balance if all this material lands back into the solar nebula, in a uniform manner, and over a long period of time.

For such a model to work, we require a fairly large component of matrix material to be made from dust that has been recycled through the protostellar jet. There are at least two consequences if this idea is correct. First, this recycled dust would have to be more refractory than CI material, since the jet formation region of the inner solar nebula would have been far warmer than the regions where most chondrites were formed. Second, if matrix material were formed from dust that had been lofted into the upper atmosphere of the solar nebula then this may be an observable phenomenon.

To understand this latter point, we need an estimate for the amount of dust that should be resident in the upper atmosphere of the solar nebula at any particular time.
Protostellar jets have average mass loss rates of order $10^{-8} \, M_\odot \, \text{yr}^{-1}$ and since the mass of dust to gas in the Interstellar Medium (ISM) is 1/100, this implies that $\lesssim 10^{-10} \, M_\odot \, \text{yr}^{-1}$ of dust is blown out by the protostellar jet. So, an upper limit for the amount of dust that is lobbed into the upper atmosphere of an accretion disk in a Classical T Tauri Star (CTTS) per year is $10^{-10} \, M_\odot$.

The dust settling timescale ($\tau_{\text{settle}}$) can be deduced from Eq. (15) with $v_{gz} = 0$, i.e.

$$m_p \ddot{z} = -\frac{CD}{2} \rho_g \dot{z}^2 A_p - \frac{GMm_p \dot{z}}{[r^2 + z^2]^{3/2}}. \quad (84)$$

The motion of the dust is subsonic, so the drag coefficient takes the Epstein form for $C_D$ (Eq. 18). To compute the velocity of the particle, we note that $\ddot{z} \approx 0$, which implies

$$\dot{z} \approx -\frac{\omega_K^2 (r, z) \rho_p a_p}{\rho_g \bar{v}} \, z, \quad (85)$$

where $\omega_K (r, z) = \sqrt{GM/[r^2 + z^2]^{3/4}}$ is the Keplerian angular velocity at the point $(r, z)$, and $\bar{v} = \sqrt{8kT_g/\pi m_g}$ is the mean Maxwellian speed of the gas particles.

If $z \ll r$ then $\omega_K (r, z) \approx \omega_K (r) = \sqrt{GM/r^3}$ and we can compute the settling time:

$$\tau_{\text{settle}} \approx 8,000 \left(\frac{\rho_g/10^{-11} \, \text{g cm}^{-3}}{(a_p/1 \, \mu \text{m}) (\rho_p/1 \, \text{g cm}^{-3})} \right) \left(\frac{\bar{v}/1 \, \text{km s}^{-1}}{(r/1 \, \text{AU})^3} \right) \text{yr}. \quad (86)$$

Thus, we have characteristic dust-settling timescales in the range $10^3 - 10^4$ years, which means that up to $10^{-7} - 10^{-6} \, M_\odot$ of dust will be in the upper atmosphere of a CTTS accretion disk at any one time.

As discussed in Natta (1993), $10^{-7} \, M_\odot$ of high-altitude dust may produce the observed “flat-temperature distributions” in CTTSs. These temperature distributions are a surprisingly common phenomenon in CTTSs. They arise when the temperature of the disk does not decrease as rapidly with distance from the protostar as one would predict from standard accretion disk theory. Natta suggested that a spherical halo of dust around a CTTS would reflect light from the protostar and into the accretion disk, thereby increasing the temperature of the outer disk.

Of course, the idea that protostellar jets can loft dust into the outer parts of the surrounding accretion disk requires a quantitative investigation to determine whether dust grains can be ejected from the jet flow. There will be a size limit where particles smaller than a certain size will simply be entrained in the jet flow and ejected from the system. We simply note, that we require dust ejection similar to that shown in §5 to account for chondrule-matrix complementarity.

7. CONCLUSIONS

Chondritic meteorites are typically an agglomeration of igneous rocks, i.e., chondrules and refractory inclusions (formation temperatures 1500-2000K), surrounded by sedimentary material that, in some cases, has never experienced temperatures greater than 500K. This unusual structure has prompted theorists to develop heating mechanisms (e.g. lightning) that can provide brief, intense impulses of energy in the otherwise cold outer regions of the solar nebula. These energy impulses are presumed
to have melted small dust aggregates into chondrules, which were then incorporated into larger dust aggregates that eventually formed meteorites.

Such energy impulse theories are not required in the ‘Jet’ model of chondrule formation. Chondrules are formed in the hot inner regions of the accretion disk adjacent to the protostar. There is no difficulty in obtaining the required temperatures, because in or near this region the accretion disk dumps around half of its gravitational energy. The protostellar jet ejects the chondrules from the hot inner disk and gas drag brings these particles back to the cooler outer regions, where they could be incorporated into growing aggregates of cool nebular material. The “plum pudding” structure of chondritic meteorites is a natural consequence of this model.

One of the many problems in chondrule formation is the deduced low cooling rate (∼1 - 1000 K/hour) for these particles. This cooling rate is many orders of magnitude smaller than that expected for a particle radiating directly into space. In our model, chondrules are formed in the optically thick regions of the inner accretion disk. They will, therefore, be in thermodynamic equilibrium with the gas and their temperature variations will be damped.

Of course if the chondrules were to remain in this environment, their temperatures would not decrease. Chondrules, however, are produced by the ablative interaction between a streaming gas flow (perhaps produced from the merging of magnetic field lines) and molten material. Particles, that are small enough, will be swept up with the flow. As the gas flow moves away from the midplane of the inner accretion disk, it will expand and therefore probably cool. Particles that are entrained in this gas flow will also cool at the same rate.

Particles that are moving with the gas flow may be ejected from the accretion disk, but it is possible that a particle may simply hover at some distance away from the midplane of the disk. If the gas density in the flow is below a critical gas density (∼10^{-11} g cm^{-3}) then the particles will undergo damped oscillations around their hover points. It is due to these oscillations, we suggest, that chondrules can undergo reheating. Smaller particles of the same mass density will still move past the hovering particles, and so may collide with these particles.

Such a scenario allows for the formation of adhering compound chondrules. These chondrules pose a major problem for chondrule formation theories as it is nearly always the smaller chondrules that were plastic at the time of collision. Smaller particles will lose heat more readily than larger particles, since they have a larger surface to area ratio. If chondrules were formed in a single flash heating event, we should expect the smaller particles to become solid before the larger particles. That we actually see the opposite behaviour (i.e., the large particles were solid, while the small particles were semi-molten), strongly suggests that single flash-heating models require some modification.

In the Jet model, adhering compound chondrules are formed when small chondrules collide with larger chondrules that are stationary in the flow. The smaller particles that are moving with the flow will be warmer than the larger hovering particles, because the latter particles have had time to equilibrate with the local temperature of the gas. If the temperature gradient in the flow is sufficiently steep (∼10 to 100 K/km) at this collision point, compound chondrules with the observed structure will be formed.

Once a compound chondrule has been formed, and has stopped spinning, the smaller secondary chondrule will orient the entire compound chondrule, just like a weathervane, such that the secondary chondrule is pointing in the direction of the flow. Any incoming secondary chondrule will hit the rear end of the primary chondrule. The resulting triple compound chondrule, will have two secondaries
chondrules that will tend to “avoid” each other. Of course, this prediction implicitly assumes a fairly steady gas flow, an assumption that may be incorrect. However, if this prediction is found to be valid, it would be major piece of evidence in favour of the Jet model, as it is difficult to produce such an effect with other chondrule formation theories.

The fragmentation of chondrules is a simple extension of compound chondrule formation, where instead of one particle being plastic at the time of collision, we now have two solid particles colliding at higher velocities. We expect the fragmentation zone to be ‘above’ the compound-chondrule formation zone, since the gas flow will be cooler and the flow speed higher as one moves further away from the midplane of the disk.

The chondrule size range is yet another aspect of chondrule formation which has not been satisfactorily explained. In the Jet model, the sizes of the ejected droplets are determined by the balance between gas drag and the surface tension. A jet flow with high energy density will tend to make molten droplets smaller, while molten materials with higher surface tensions will tend to form larger droplets. Observations suggest that protostellar jets are produced within 0.1 AU of solar-mass protostars. The minimum energy density of a wind that can eject Si and Fe droplets from such close proximity to the protostar is such that the radii of these droplets is $\leq 1$ cm. Protostellar winds that form larger droplets, will have a lower energy density and will not be able to eject them from the jet formation region and we will not see them in meteorites.

This ejection mechanism may also explain the complementary chemical structure of chondrules and their surrounding matrix. For this to occur, we require that refractory dust as well as chondrules are ejected by the jet flow so that mass balance will be obtained once the dust settles back to the solar nebula. This leads to a simple calculation, which suggests that up to $10^{-7} - 10^{-6}$ M$_\odot$ of dust will be in the upper atmosphere of a CTTS accretion disk at any one time. These figures appear to be consistent with those deduced from observation (Natta 1993).

The transfer of dust and chondrules from the inner to outer portions of the accretion disk, necessarily requires the removal of angular momentum from the upper atmosphere of the accretion disk. This will increase the viscosity of the disk and hasten the stratified infall of disk material, i.e., the upper atmosphere of the disk will accrete more readily than the central plane of the disk.

Finally, aerodynamic size sorting of the ejected particles is a natural consequence of the Jet model. Ejected particles that are captured by the accretion disk must suffer aerodynamic size sorting, since the particles will be subject to gas drag. Particles that are not subject to gas drag will simply leave the protostellar system, since their initial velocities were higher than the escape velocity of the system.

It is for all these reasons that we consider chondrules to be ablation droplets formed by a protosolar jet in the first $10^6 - 10^7$ years of the solar system.

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Fig. 11.— A pictorial summary of the ideas presented in this paper. An accretion disk surrounds a protostar. At a distance within 0.1 AU of the protostar, the disk produces a protostellar jet (only one side of which is shown). Infalling km-sized bodies move into this hot “jet zone” and are subsequently ablated by the jet wind. The resulting droplets, if small enough, move with the flow and away from the ablating body. It is our claim that these ablation melt-droplets eventually become chondrules. At higher altitudes, within the accretion disk, some of the larger particles may hover in the gas flow. Smaller particles will still move with the flow, however, and may collide with the hover particles. These chondrule-chondrule collisions produce compound chondrules and chondrule fragments. The fragmentation of chondrules occurs at higher altitudes relative to the compound-chondrule formation zone, because the flow cools as it increases in altitude and the two colliding chondrules will be solid at the time of collision. If the protostellar jet is sufficiently powerful, it will eject particles from the inner accretion disk, and if the orbital angular momentum of the ejected particles is large enough, the particles will move out of the jet flow and travel across the face of the accretion disk. With sufficient gas drag from the upper atmosphere of the accretion disk, these (now size-sorted) particles will be brought down to the outer accretion disk, where they will be incorporated into planetesimals.
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