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A critical analysis of shock models for chondrule formation

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ABSTRACT

In recent years many models of chondrule formation have been proposed. One of those models is the processing of dust in shock waves in protoplanetary disks. In this model, the dust and the chondrule precursors are overrun by shock waves, which heat them up by frictional heating and thermal exchange with the gas.

In this paper we reanalyze the nebular shock model of chondrule formation and focus on the downstream boundary condition. We show that for large-scale plane-parallel chondrule-melting shocks the postshock equilibrium temperature is too high to avoid volatile loss. Even if we include radiative cooling in lateral directions out of the disk plane into our model (thereby breaking strict plane-parallel geometry) we find that for a realistic vertical extent of the solar nebula disk the temperature decline is not fast enough. On the other hand, if we assume that the shock is entirely optically thin so that particles can radiate freely, the cooling rates are too high to produce the observed chondrules textures. Global nebular shocks are therefore problematic as the primary sources of chondrules.

1. Introduction

The origin of chondrules is one of the biggest mysteries in meteoritics. These 0.1–1 mm sized silicate once molten droplets, abundantly found in chondritic meteorites, must have cooled and solidified within a matter hours (e.g. Hewins et al., 2005). From short-lived radionuclide chronology data (e.g. Villeneuve et al., 2009) it is known that this must have taken place during the first few million years after the start of the Solar System, during the phase when the Sun was still likely surrounded by a gaseous disk (the "solar nebula"). What makes chondrule formation mysterious is that this few-hour cooling time is orders of magnitude shorter than the typical few-million-year time scale of evolution of the solar nebula. Chondrules can thus not be a natural product of the gradual cooling-down of the nebula. Instead, chondrules must have formed during "flash heating events" of some kind – but which kind is not yet known. There exists a multitude of theories as to what these flash heating events could have been. Boss (1996) and Ciesla (2005) give nice overviews of these theories and their pros and cons. So far none of these theories has been universally accepted.

One of the most popular theories is that nebular shock waves can melt small dust aggregates in the nebula, causing them to become melt droplets and allowing them to cool again and solidify (Hood and Horanyi, 1991). The origin of such shocks could, for instance, be gravitational instabilities in the disk (e.g. Boley and Durisen, 2008) or the effect of a gas giant planet (e.g. Kley and Nelson, 2012). Detailed 1-D models of the structure of such radiative shocks, and the formation of chondrules in them, were presented by Iida et al. (2001), Desch and Connolly (2002), Ciesla and Hood (2002) or Morris and Desch (2010). These models show that such a shock, in an optically thick solar nebula, would lead to a temperature spike in the gas and the dust that lasts for only a few seconds to minutes with cooling rates of up to 10^3 K/h, followed by a more gradual cooling lasting several hours, with cooling rates of the order of 50 K/h. These appear to be the right conditions for chondrule formation, which is one of the reasons why this model is one of the favored models of chondrule formation nowadays.

In this paper we revisit this shock-induced chondrule formation model. Our aim is to investigate the role of the downstream boundary condition and the role of sideways radiative cooling.

If the shock is a large scale shock, e.g. due to a global gravitational instability, then on a small scale (the scale of several radiative mean free paths λ_{free}) the shock can be modeled as a infinite 1-D radiation hydrodynamic flow problem. The "infinite 1-D" in this context means that the 3-D geometry of the problem only becomes important on scales $\gg \lambda_{free}$, so that on scales $\sim \lambda_{free}$ a

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1-D geometry can be safely assumed where the pre-shock boundary is set at $x = -\infty$ and the post-shock boundary is set at $x = +\infty$. With $\lambda_{\text{free}} = (\kappa\rho)^{-1}$ the pressure scale height of a typical minimum mass solar nebula is several hundred times larger than the optical mean free path. The shock is assumed to be stationary in the x-coordinate system, i.e. the coordinates move along with the shock and the shock is always at $x = 0$. In such an infinite 1-D system the full shock structure can be reconstructed when the physical variables at the upstream boundary $x = -\infty$ are all given. No downstream boundaries at $x = +\infty$ need to be given. In fact, the physical variables at $x = +\infty$ follow uniquely by demanding that the mass-, momentum- and energy-flux at $x = +\infty$ equals those of $x = -\infty$, but with subsonic gas velocity at $x = +\infty$. This gives a *global* Rankine–Hugoniot condition including all the radiative and dust physics. Note that right at the shock at $x = 0$ the jump in the gas variables is given by a *local* Rankine–Hugoniot condition in which only the gas fluxes on both sides are set equal. We will work out these shock models in Section 2 and show that after the temperature spike they lack a slower few-hour cooling phase. Instead, they stay at a constant temperature, i.e. the temperature in accordance with the global Rankine–Hugoniot condition. We will investigate in Section 3.1 whether the chondrule peak temperature can be made high enough for melting while keeping the post-spike temperature low enough to retain volatile elements. According to Fedkin et al. (2012) and Yu et al. (1996) the high-temperature phase should be of the order tens of minutes rather than of hours. From textural constraints, Hewins et al. (2005) and Desch et al. (2012) conclude that the cooling rates have to be of the order of 10^1–10^3 K/h.

Since the infinite 1-D shock solution is a geometric approximation we will implement effects of sideways cooling (i.e., from top and bottom of the disk) in Section 3.2 to improve the realism of the model. This will re-introduce the slow cooling phase after the temperature spike, but we find that this slow cooling phase is of the order of weeks/months/years rather than hours, because 2-D/3-D radiative diffusion will take place on x-scales of the same order as the vertical scale height of the disk.

Finally we discuss in Section 3.3 two versions of the shock-induced chondrule formation where the cooling time can be rapid. One is a locally induced one, for instance due to a supersonic planetesimal bow shock (Hood, 1998; Morris et al., 2012). The other is if the disk is fully optically thin, allowing the post-shock material to cool straight to infinity.

2. The model

Our one-dimensional shock model is built on the work by Desch and Connolly (2002). We used their approach, generalized it to arbitrarily many gas species, particle populations and chemical reactions and modified it where we think modifications or corrections are needed.

In this model we assume that all the parameters (densities, temperatures, velocities, etc.) only change along one direction, the x-direction. Therefore, the model consists of infinitely extended, plane-parallel layers of constant temperatures, velocities and densities.

2.1. Radiative transfer

Even though the model is one-dimensional, we must allow the photons to move into all three directions. But fortunately, because of the plane-parallel assumption the radiative transfer equations here can be vastly simplified.

In this case the optical depth τ is a monotonic increasing function of the distance from the post-shock boundary and therefore a measure of the position x inside the computational domain:

$$\tau(x) = \int_x^{x_{\text{post}}} \alpha \, \mathrm{d}x, \tag{1}$$

where x_{post} is the location of the post-shock boundary. This means the optical depth increases from $\tau(x_{\text{post}}) = 0$ at the post-shock boundary to a maximum value of $\tau(x_{\text{pre}}) = \int_{x_{\text{pre}}}^{x_{\text{post}}} \alpha \, \mathrm{d}x \equiv \tau_{\max}$ at the pre-shock boundary. α is the absorption coefficient, which is the sum of the contribution of the gas and the particles:

$$\alpha = \alpha_{\text{g}} + \alpha_{\text{p}} = \rho_{\text{g}}\kappa_{\text{P}} + n\pi a^2 \varepsilon. \tag{2}$$

The absorption coefficient of the gas is the product of the gas mass density ρ_{g} and the temperature-dependent Planck-mean opacity κ_{P}, which we took from Semenov et al. (2003) (see Fig. 1). The absorption coefficient of the particles is the product of their number densities n, their geometrical cross-section and their absorption efficiency ε, which we adopted from Desch and Connolly (2002) as

$$\varepsilon = 0.8 \times \min\left[1, \left(\frac{a}{2\ \mu\text{m}}\right)\right], \tag{3}$$

where a is the particle radius.

To calculate the thermal histories of the particles one has to calculate the mean intensity $J_{\text{rad}}(\tau)$ at every position τ our radiative transfer is gray, i.e. wavelength independent. The mean intensity is defined as the average intensity I per solid angle coming from all directions

$$J_{\text{rad}}(\tau) = \frac{1}{4\pi}\int_{\Omega} I(\tau, \mu)\mathrm{d}\Omega. \tag{4}$$

In the plane-parallel assumption the intensity depends only on the position τ and the angle θ between the incoming ray and the x-axis. μ is defined as $\mu = \cos\theta$. Therefore, the mean intensity can be simplified to

$$J_{\text{rad}}(\tau) = \frac{I_{\text{pre}}}{2}E_2(\tau_{\max} - \tau) + \frac{I_{\text{post}}}{2}E_2(\tau) + \frac{1}{2}\int_0^{\tau_{\max}} S(\tau')E_1(|\tau' - \tau|)\mathrm{d}\tau', \tag{5}$$

by using the exponential integrals (see Mihalas and Weibel-Mihalas, 1999)

$$E_n(x) = \int_1^{\infty} \frac{e^{-xt}}{t^n}\mathrm{d}t, \tag{6}$$

$I_{\text{pre}} = \frac{\sigma}{\pi}T_{\text{pre}}^4$ and $I_{\text{post}} = \frac{\sigma}{\pi}T_{\text{post}}^4$ are the incoming radiations from both boundaries. The source function $S(\tau)$ is defined as

$$S = \frac{\rho_{\text{g}}\kappa_{\text{P}}B(T_{\text{g}}) + n\pi a^2\varepsilon B(T)}{\rho_{\text{g}}\kappa_{\text{P}} + n\pi a^2\varepsilon}, \tag{7}$$

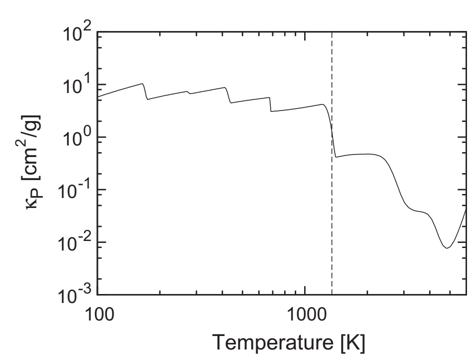

Fig. 1. Planck-mean opacity from Semenov et al. (2003). At ~1350 K (vertical line) the fine-grained dust associated with the gas gets evaporated. Therefore, the opacity drops by one order of magnitude at this temperature. This is responsible for the "opacity knee" in the pre-heating phase, where the gas becomes optically thin.