Dispersal of Gaseous Circumstellar Discs around High-Mass Stars

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ABSTRACT
We study the dispersal of a gaseous disc surrounding a central high-mass stellar core once this circumstellar disc becomes fully ionized. If the stellar and surrounding EUV and X-ray radiations are so strong as to rapidly heat up and ionize the entire circumstellar disc as further facilitated by disc magnetohydrodynamic (MHD) turbulence, a shock can be driven to travel outward in the fully ionized disc, behind which the disc expands and thins. For an extremely massive and powerful stellar core, the ionized gas pressure overwhelm the centrifugal and gravitational forces in the disc. In this limit, we construct self-similar shock solutions for such an expansion and depletion phase. As a significant amount of circumstellar gas being removed, the relic disc becomes vulnerable to strong stellar winds and fragments into clumps. We speculate that disc disappearance happens rapidly, perhaps on a timescale of $\sim 10^3 - 10^4$ yr once the disc becomes entirely ionized sometime after the onset of thermal nuclear burning in a high-mass stellar core.

Key words: circumstellar matter – hydrodynamics – plasmas – shock waves – stars: early type – stars: winds, outflows

1 INTRODUCTION

There are ample evidence that young stars are often accompanied by circumstellar discs, as part of star formation processes. Discs are most frequently detected for low-mass T Tauri protostars. The probability of detecting a circumstellar disc decreases with an increasing stellar mass. Natta et al. (2000) concluded that a lack of discs around Herbig Be stars may result from their faster evolution. Fuente et al. (2003) revealed the first evidence of dusty discs around Herbig Be stars, yet the disc-to-star mass ratio $M_d/M_*$ is at least an order of magnitude lower than those of Herbig Ae and T Tauri stars, suggesting that discs may evolve much more rapidly around very massive stars. Suppose that high-mass stars form similarly as low- or intermediate-mass stars in the sense of involving collapse (e.g., Lou 1996; Lou & Shen 2004; Bian & Lou 2005; Yu & Lou 2005; Yu, Lou, Bian & Wu 2006; Lou & Gao 2006) and accretion (e.g., McKee & Tan 2002; Blum et al. 2004), then a circumstellar disc should form during the early evolution phase. Observational evidence for accretion and massive discs during high-mass star forming stages were indeed reported (e.g., Sundell et al. 2003; Pestalozzi et al. 2004). While massive stars of late O or early B types may have disc signatures, more massive stars lack clear evidence for discs (e.g., Blum et al. 2004). Unless a massive star forms through other processes without involving a circumstellar disc, the key issue is to understand the mechanism for a fast and efficient dispersal (primarily hydrogen gas) of circumstellar discs around the most massive stars.

Circumstellar discs are important because they are sites where a companion star or planets may form. For a massive star, a considerable mass fraction of its dense circumstellar disc appears to be rapidly removed. The currently estimated lifetime of circumstellar discs for low-mass stars is of the order of several million years (e.g., Strom et al. 1993; Briceño et al. 2001). For intermediate- and high-mass stars, the timescale for the existence of circumstellar discs can be much shorter because of faster evolution and stronger radiative power output from central stars (e.g., Natta et al. 2000). Theoretically, the dispersal or destruction of a circumstellar disc may be modelled through viscous disc accretion onto a central star, photoevaporation by either stellar or external UV radiations, interactions with other stars or consumption by planet formation, or any combinations of these processes (e.g., Hollenbach et al. 2000). The photoevaporation and accretion models have been widely applied to disc dispersals of
low-mass T Tauri stars (e.g., Johnstone et al. 1998; Störzer & Hollenbach 1999; Clarke et al. 2001).

We here advance a plausible physical scenario for the dispersal of circumstellar discs around most massive stars. By forming a massive protostellar core, the accretion gradually ceases and a massive (perhaps \( \gtrsim 0.1 - 0.3 \, M_\odot \)) centrifugally supported H I disc\(^1\) forms around the core. Near the main-sequence when the stellar core ignites thermal nuclear reactions to produce a profuse amount of extreme-ultraviolet (EUV) photons, an ionization front (IF) travels rapidly outward to ionize the entire H I disc to a high temperature \( \gtrsim 10^4 \, \text{K} \) (e.g., Osterbrock 1989). Meanwhile, X-ray emissions associated with protostellar and disc magnetic activities also ionize gas to \( \sim \times 10^5 \) EUV photons, an ionization front (IF) travels rapidly outward to ionize the entire H I disc to a high temperature \( \gtrsim 10^4 \, \text{K} \) (e.g., Balbus & Hawley 1998; Balbus 2003) in the weak magnetic field regime sustain a level of MHD turbulence to incessantly mix ionized and neutral components and to effectively enhance ionization. To estimate, we assume a diluted neutral H I disc\(^2\) with a surface mass density \( \Sigma = \Sigma_0 (r/R)^{-1/2} \), for \( M_\star \sim 0.02 \, M_\odot \), we derive \( \Sigma_0 = 3 M_\star / (4 \pi R_\star^2) \sim 0.042 \, \text{g cm}^{-2} \) such that the H I column density normal to the disc is \( N_H (r) \sim 2.5 \times 10^{22} (r/R)^{-1/2} \, \text{cm}^{-2} \). The Lyman flux incident on the disc will penetrate a column \( N_L (r) \lesssim 1.8 \times 10^{22} \Phi_{\text{Ly}} (r/R)^{-1/2} \, \text{cm}^{-2} \) (see Appendix A of Hollenbach et al. 1994). Thus an H I disc may be fully ionized\(^3\) for \( \Phi_{\text{Ly}} \sim 2 \), which is readily satisfied for early O stars (e.g., Panagia 1973). For more massive stars with \( \Phi_{\text{Ly}} \sim 10 - 100 \), a denser circumstellar disc of mass \( M_\star \) even higher than \( \sim 0.2 \, M_\odot \) might be ionized completely as facilitated by X-ray fluxes and disc MHD turbulence. If submerged in an external X-ray and EUV radiation field from nearby massive stars within a local star cluster (which is the usual case), a complete ionization of a circumstellar disc becomes more likely. A hot H II disc inevitably expands as a result of the overwhelming pressure gradient. Typically, the gas density is higher in the inner part than in the outer part of a disc, so does the pressure gradient. Hence, the inner disc expands faster than the outer disc does and consequently a shock develops in the disc to travel radially outward. The disc portion behind the shock will continue this expansion to reduce the disc surface mass density.

The main thrust of this Letter is to construct a gas dynamic model for a such disc expansion phase after a circumstellar disc has been ionized and to estimate disc disappearance timescales using available data. While our results are preferably applicable to high-mass stars ( \( \gtrsim 10 \, M_\odot \)) they are also relevant for discs around low- or intermediate-mass stars in binaries where the disc of the secondary may be fully ionized by the X-ray and EUV field of the more massive primary.

### 2 DISC MODEL FORMULATION

We start with a specific case of a neutral H I disc carrying a surface mass density \( \Sigma \propto r^{-1} \), and then extend the results to different power-law indices of \( \Sigma \) profiles (Shen & Lou 2004a; Shen, Liu \\& Lou 2005; Lou \\& Bai 2006; Wu \\& Lou 2006) in the limit when the thermal pressure dominates other forces in the shock evolution of an H II disc. The stellar gravity has weaker effects over distances larger than \( \tilde{r} \sim GM_\star / (2a^2) \) where \( a \) is the sound speed of the H I gas. For \( M_\star \sim 10 \, M_\odot \) and \( a \sim 10 \, \text{km s}^{-1} \), we have \( \tilde{r} \sim 40 \, \text{AU} \) and may neglect the stellar gravity on the main portion of the disc with a typical radius \( R \sim 1000 \, \text{AU} \).

We presume a thin neutral H I circumstellar disc around a protostar in a rotational equilibrium initially with a singular isothermal disc (SID) profile, namely

\[
\Sigma(r) = \frac{a_0^2 (1 + D^2)}{2 \pi G r} , \quad M(r) = \frac{a_0^2 (1 + D^2) r^2}{G} , \quad j(r) = a_0 D r ,
\]

where in cylindrical coordinates \( (r, \theta, z) \), \( \Sigma(r), M(r) \) and \( j(r) \) are the surface mass density, the enclosed mass inside radius \( r \) and the \( \hat{z} \)-component specific angular momentum, respectively; \( G \) is the gravitational constant and \( D \) is the rotational Mach number defined by \( D \equiv V_0 / a_0 \) with \( V_0 \) and \( a_0 \) being the rotation speed\(^4\) and the isothermal sound speed of the H I gas.

At a certain initial time \( t = 0 \), thermal nuclear reactions start in the stellar core and the resulting radiation rapidly ionizes the circumstellar disc in a short time to reach an isothermal state with a faster sound speed \( a \). To keep the initial SID profile unchanged would require

\[
\Sigma(r, 0^+) = \frac{\epsilon^2 a^2 (1 + D^2)}{2 \pi G r} , \quad j(r, 0^+) = \epsilon a D r ,
\]

where \( \epsilon \equiv a_0 / a \leq 1 \) is the square root of the temperature ratio. Now the pressure and centrifugal forces overwhelm the gravitational force and a circumstellar disc expands inevitably. As the pressure gradient is steeper in the inner disc, a shock will emerge to travel outward. This is a close analog

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\(^1\) Discs around massive stars may be even more massive initially (see Sandell et al. 2003).

\(^2\) Disc radii \( R \) of high-mass stars are expected to be larger than those of T Tauri or Herbig Ae/Be stars, which extend from tens to hundreds of AU's (e.g., McCaughrean 1997; Meyer \\& Beckwith 2000). We take \( R \sim 10^3 \, \text{AU} \) (e.g., Pestalozzi et al. 2004).

\(^3\) Around massive stars, dust grains are effectively evacuated to larger radii (\( \sim 0.1 \, \text{pc} \)) due to sublimation, stellar winds and radiation pressure (e.g., Chini et al. 1987; Churchwell 2002). When dust attenuation is included, the required relic disc mass will be smaller. In the approximation of Franco et al. (1990), the incident Lyman flux \( \Phi_{\text{Ly}} \) is reduced by a factor \( e^{-\tau_2} \) where \( \tau_2 \) is the optical depth at the boundary of H II regions. For \( \tau_2 \gtrsim 1 \), the disc mass that can be fully ionized is reduced by a factor less than 0.6.

\(^4\) We take on a flat rotation curve instead of the usual Keplerian rotation curve. As can be seen later, the centrifugal force can be omitted in the limit when the pressure dominates and thus the choice of different rotation curves does not really matter.
of the so-called ‘champagne flows’ in H II clouds surrounding OB stars (Tenorio-Tagle 1979; Franco et al. 1990). Since the latter may evolve towards a self-similar phase (Shu et al. 2002), it is suggestive of a similar description for a disc expansion with a shock. Here, we focus on axisymmetric similarity shocks relevant to the disc dispersal scenario.

Under the axisymmetry, we use the following nonlinear equations to describe disc dynamics in cylindrical geometry.

\[
\frac{\partial \Sigma}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} (r \Sigma u) = 0 ,
\]

\[
\frac{\partial u}{\partial t} + u \frac{\partial u}{\partial r} - \frac{j^2}{\Sigma} = - \frac{1}{2} \frac{\partial \Pi}{\partial r} - \frac{\partial \Phi}{\partial r} ,
\]

\[
\frac{\partial j}{\partial t} + \frac{\partial j}{\partial r} = 0 ,
\]

\[
\frac{\partial M}{\partial t} + u \frac{\partial M}{\partial r} = 0 ,
\]

where \( u \) is the radial speed, \( \Phi \) is the gravitational potential, \( \Pi \) is the two-dimensional isothermal self-gravitating non-linear gas pressure, \( M(r, t) \equiv \int_0^\infty \Sigma(r', t) 2 \pi r' dr' \) is the enclosed mass within radius \( r \). The Poisson integral is

\[
\Phi(r, t) = \frac{-G \Sigma(r', t)}{(r^2 + r^2 - 2r \cos \theta)^{1/2}} .
\]

We introduce the following similarity transformation

\[
x \equiv \frac{r}{\alpha(t)} , \quad \Sigma(r, t) \equiv \frac{a \alpha(x)}{2 \pi G t} , \quad u(r, t) \equiv \alpha v(x) , \quad j(r, t) \equiv \beta a^2 m(x) , \quad M(r, t) \equiv \frac{a^4 t}{G} m(x) , \quad j(r, t) \equiv \beta a^2 m(x) ,
\]

in the two-dimensional isothermal self-gravitating non-linear ideal gas equations (9 - 10), where \( \alpha(x), v(x) \) and \( m(x) \) are the reduced surface mass density, radial flow speed and enclosed mass, all being dimensionless functions of \( x \). Parameter \( \beta \equiv D/[(1 + D^2)] \) is determined from the equilibrium. The last one in transformation (7) holds as the ratio \( j(r)/M(r) \) remains constant since \( t \to 0 \) by the angular momentum conservation (Li & Shu 1997; Saigo & Hanawa 1998).

In the so-called monopole approximation,5 the similarity ordinary differential equations (ODEs) become

\[
\frac{[x - v]^2 - 1}{x - v} \frac{d e}{d x} = \frac{\alpha(x - v)[1 - \beta^2 \alpha(x - v)] - 1}{x} ,
\]

\[
[1 - \frac{[x - v]^2}{x - v} \alpha] \frac{d \alpha}{d x} = \frac{\alpha[1 - \beta^2 \alpha(x - v)] - (x - v)}{x} ,
\]

which, together with \( m = x(x - v) \alpha \), are to be solved for specified ‘asymptotic conditions’.

The leading asymptotic solutions at \( x \to +\infty \) for either \( t \to 0^+ \) or \( r \to +\infty \) are

\[
v \to V + \frac{1 - A + \beta^2 A^2}{x} \alpha \to A , \quad m \to Ax ,
\]

where \( V \) and \( A \) are two constants and higher-order terms

5 A multipole treatment can be found elsewhere (Li & Shu 1997). This monopole approximation does not affect our results much because when \( \epsilon \ll 1 \), the self-gravity becomes negligible (Shu et al. 2002; Shen & Lou 2004b).

The three rows without values for \( D \) and \( \epsilon \) are unphysical cases.

| Table 1. Parameters for numerical solutions |
|---------------------------------------------|
| \( B \) | \( x_s \) | \( A \) | \( \beta \) | \( D \) | \( \epsilon \) |
|---------------------------------------------|
| 0.1 1.92 0.136 0.0 0.369 |
| 1.93 0.138 0.5 0.189 0.365 |
| 1.96 0.144 2.0 1.166 0.247 |
| 0.01 1.99 0.0155 0.0 0.124 |
| 2.00 0.0158 2.0 0.260 0.122 |
| 2.00 0.0158 8.0 – – |
| 10\(^{-3}\) 2.00 1.58 × 10\(^{-3}\) 0 0 0.0397 |
| 2.00 1.58 × 10\(^{-3}\) 8.0 0.335 0.0377 |
| 2.00 1.58 × 10\(^{-3}\) 32 – – |
| 10\(^{-4}\) 2.00 1.58 × 10\(^{-4}\) 0 0 0.0126 |
| 2.00 1.58 × 10\(^{-4}\) 32 0.439 0.0115 |
| 2.00 1.58 × 10\(^{-4}\) 128 – – |

The result of solving the ODEs (5, 6) numerically, we specify \( \beta \) leading to a relation for \( D \) and \( \epsilon \) by \( \beta \equiv D/[(1 + D^2)] \). A shock solution is constructed as follows: (i) Integrate ODEs (5, 6) from the origin using solution (11) for various \( B \); (ii) Impose shock condition (12) at each integration step to cross the sonic critical line \( x = 1 \) and continue for larger \( x \); (iii) Pick out the one that matches condition (10) as \( x \to +\infty \) with \( V = 0 \). By matching a proper \( A \) value, this last step leads to another relation of \( D \) and \( \epsilon \) by \( A = \epsilon^2(1 + D^2) \). Both \( D \) and \( \epsilon \) are then determined accordingly.

Remarkably, it is found that in the limit \( B \to 0 \) for a disc heated to high temperatures, there exists an invariant shock solution not affected by \( \beta \). The \( A \) value is uniquely determined by \( B \) with \( A \to 0 \) as \( B \to 0 \), while the ratio \( A/B \) approaches a constant \( \sim 1.58 \) (Table 1). Physical solutions require \( A \beta^2 < 1 \). This bears a strong resemblance to the spherical similarity solutions of ‘champagne flows’ (Shu et al. 2002). This invariant form describes a situation when the thermal pressure overwhelms in the disc expansion with negligible gravity and centrifugal forces. The limit of \( x_s \) \( = 2.00 \) gives a shock speed twice the ionized sound speed \( a \). The downstream (post-shock) gas becomes more uniformly distributed and a linear expansion emerges from small \( x \).

Following the above procedure by ignoring gravity and centrifugal forces, one can extend these results to a power-law surface mass density \( \Sigma(r) \propto r^{-\alpha} \) where \( 0 < \alpha \leq 2 \) (Shen et al. 2005). Similar to the \( n = 1 \) case, the downstream behind the shock approaches a uniform density yet with the enclosed mass \( \propto t^{-\eta} \) and the flow speed becomes linear in the radii, viz., \( v \to \alpha x/2 \) as \( x \to 0 \).
Figure 1. A shock solution for $v(x)$, $a(x)$ and $m(x)$ versus $x$. The dash-dotted line in the top panel is the sonic critical line $x - v = 1$. The reduced radial speed $v(x)$ and density $a(x)$ encounter discontinuities at shock location $x_c = 2.00$, while the reduced enclosed mass $m(x)$ is continuous by the mass conservation (13) and (14). Downstream (post-shock) and upstream (pre-shock) portions are shown in heavy and light solid curves; the dotted curves denote the leading-order terms of asymptotic solutions (15).

3 CIRCUMSTELLAR DISCS

We use our self-similar invariant shock solutions to estimate a circumstellar disc dispersal timescale. To elucidate the concept, we introduce the following fiducial solution using the SID model. For a centrifugally supported neutral equilibrium SID, we may choose $D = 2$ for a supersonic rotation against self-gravity\(^6\). The Lyman ionization may lead to a $a_0/a$ ratio of $\epsilon = 0.02$ (e.g., $a_0 \sim 0.2 \text{ km s}^{-1}$ and $a \sim 10 \text{ km s}^{-1}$). We then have $A = 2 \times 10^{-3}$ which gives a $B = 1.26 \times 10^{-3}$ and a $\beta = 20$. Yet this solution remains essentially the same as that of $\beta = 0$ and $B = 1.26 \times 10^{-3}$.

That is, there exists an invariant shock solution (Fig. 1) not affected by $\beta$ in the limit of $B \to 0$ with a shock at $x_c = 2.00$.

Since a similarity solution does not carry characteristic timescales or lengths, one needs a starting point to apply the solution in proper spatial and temporal domains. For example, at a particular radius $R \sim 1000 \text{ AU}$, i.e., the ionized disc radius after some time of expansion, the corresponding similarity variable is $x_c = R/(x_c a t_c)$ where $t_c$ is an input timescale that a self-similar evolution has lapsed and marks the starting epoch for the subsequent evolution. Here, $t = t_c$ is a calibration point. We focus on $R$ and presume that for $t \geq t_c$, all materials leaving $R$ would be regarded as lost from a circumstellar disc. The key issue of concern is the evolution of the enclosed disc mass $M(R, t)$ for $t \geq t_c$ where $t_c$ is a starting time when the mass and radius of a disc are specified. The information may be extracted from the bottom panel of Fig. 1, i.e., the downstream portion of $m(x)$ in heavy solid curve. For a small $x < 0.5$, the reduced mass is $m(x) \cong B x^2/2$ and the evolution of $M(R, t)$ may be estimated by

$$M(R, t) \cong a B R^2 / (2 G t).$$

(13)

As a self-similar evolution starts from $t = 0$, it takes a timescale of $t_c = R/(2 a) \sim 250 \text{ yr}$ for a shock to reach $R$, where $t_c$ is the calibration time corresponding to a similarity coordinate $x_c = 2.00$. At this epoch, the disc mass behind a shock is $M(R, t_c) = a^3 t_c m(2) / G \sim 0.2 M_\odot$ where $m(2) = 3.4 \times 10^{-3}$ by the heavy solid curve of $m(x)$ in Fig. 1. From now on, the shock continues to travel into distant regions and eventually merges into ambient HII clouds; meanwhile, the subsequent self-similar evolution of the disc is well described by the downstream (post-shock) portion. We may then use interpolated values of $m(x)$ with $x \leq 2$ and $M(R, t) = a^3 t m(x) / G$ to estimate $M(R, t)$ for $t \geq t_c$. For a sufficiently small $x$ or large $t$, we may directly use equation (13) that shows an enclosed disc mass proportional to $t^{-1}$. At later times of $5 t_c$ ($x = 0.4$), $10 t_c$ ($x = 0.2$), $20 t_c$ ($x = 0.1$), $100 t_c$ ($x = 0.02$), corresponding to $\sim 1250 \text{ yr}$, $\sim 2500 \text{ yr}$, $\sim 5000 \text{ yr}$, $\sim 25000 \text{ yr}$, the enclosed disc mass becomes $\sim 0.03 M_\odot$, $\sim 0.015 M_\odot$, $\sim 0.0075 M_\odot$, $\sim 0.0015 M_\odot$, respectively.

Diagnostically, it is not easy to set a criterion for the ‘disc disappearance’ by decreasing enclosed disc mass $M(R, t)$. We propose tentatively that a ‘disc disappearance’ corresponds to a disc mass dropped between one tenth and one half of the mass at $t = t_c$. This is partly due to the fact that a disc depletion leads to fainter Hα emissions from ionized gas and infrared/submillimeter emissions from dusts, and partly due to the possibility that a rarefied disc becomes unstable by strong stellar wind shears and breaks into clumps. Thus, a timescale of disc disappearance is estimated as $t_{dis} \sim 10 t_c - 100 t_c$ by expression (13). For $t_c \sim 250 \text{ yr}$, the ‘disappearance timescale’ is $t_{dis} \sim 2.5 \times 10^3 - 2.5 \times 10^4$ yr is very much shorter than the typical timescale of several Myr for low-mass T Tauri stars.

4 DISCUSSIONS

While the self-similar process described here is highly idealized, it may catch several gross features of a disc dispersal process not fully explored so far. First, we assume a fast traveling IF together with environmental X-ray and EUV radiation field that heat up the entire disc to an isothermal state in a short time. Secondly, rotation and gravity are ignored in the limit of an overwhelming thermal pressure. Our model of invariant shock is more applicable to discs around very massive stars or to circumstellar discs exposed to intense X-ray and EUV radiation from nearby external sources (e.g., evolved massive stars or supernova explosions; Chevalier 2000). Thirdly, we ignore the disc thickness. In reality, such a disc should expand both radially and vertically besides rotation, with photoevaporation being concurrent. A thorough analysis of these aspects seems worthwhile.

We now summarize highlights of our model. First, a circumstellar disc as a whole is expanding rather than fixed or shrinking shortly after the ignition of nuclear reactions in the stellar core. This expansion induced disc mass-loss rate is modelled as self-similar and differs from classic photoevaporation models. For very massive stars submerged in an intense X-ray and EUV radiation field, circumstellar discs become too short-lived to be observed after entering the

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\(^6\) $D$ should be $\geq 1$, but may not be too large. Our basic results derived here are not very sensitive to $D$ variations.
main-sequence. Secondly, the fast moving shock serves to disperse disc materials and the rarefied disc may become vulnerable to fragmentations by stellar wind shear. These processes accelerate the disc destruction. Thirdly, this self-similar expansion tends to establish a more uniform density distribution even though the initial density profile is non-uniform. Fourthly, during this self-similar expansion phase, a considerable mass fraction of the disc is expelled or displaced to distant places; these processes might explain the presence of appreciable amounts of ionized gas at intermediate radii (i.e. $10^3 \sim 10^4$ AU) in ultracompact H II regions (e.g., Wood & Churchwell 1989; Shu et al. 2002). Moreover, as time goes on, these dispersed disc materials eventually cool down and may become a reservoir for producing Edgeworth-Kuiper Belt-like objects around a central star.

We now comment on several aspects of further model development. First, even for an isothermal model, the temperatures across a shock can be allowed to be different (Shen & Lou 2004b; Bian & Lou 2005; Yu et al. 2006). Secondly, the isothermal condition can be replaced by the more general polytropic approximation (Wang & Lou 2006; Lou & Gao 2006). Thirdly, it is of considerable interest to incorporate the effect of a magnetic field (Shen et al. 2005; Yu & Lou 2005; Lou & Zou 2004, 2006; Lou & Wu 2005; Wu & Lou 2006; Lou & Bai 2006).

Finally, we speculate that grossly spherical ‘champagne flows’ with shocks traveling in H II clouds (e.g., Franco et al. 1990; Shu et al. 2002; Shen & Lou 2004b; Bian & Lou 2005) on much larger spatial scales encompassing a circumstellar disc can be driven by an intense central stellar radiation, stellar winds and outflows from the disc sustained by photo-evaporation. Triggered by shocks and various flow or thermal instabilities, fragmentations would occur in clouds; by cooling and coagulation with time, these clumps may evolve into comet-like objects. This might be the origin of the Oort cloud around our solar system and should be common in other protostellar systems as well.

Not only for young stellar objects, circumstellar discs are also observed around main-sequence stars or even more evolved stars (e.g., Zuckerman 2001). These debris dusty discs are exposed to intense radiation from central objects during the late evolutionary stage. Disc expansions and shocks can be initiated in these discs and evolve in a self-similar manner. Even with considerable idealization, the analysis is technically challenging as rotation, central gravity and disc self-gravity need to be taken into account.

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