HOW LONG CAN TINY H I CLOUDS SURVIVE?
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ABSTRACT

We estimate the evaporation timescale for spherical H I clouds consisting of the cold neutral medium surrounded by the warm neutral medium. We focus on clouds smaller than 1 pc, which corresponds to tiny H I clouds recently discovered by Braun & Kanekar and Stanimirović & Heiles (2005). This new population of clouds is providing a challenge to our understanding of the evolution of the interstellar medium (ISM). The origin of the clouds is still under debate, but it is presumably the thermal instability (Field 1965; Balbus 1986) in turbulent gas (e.g., Vázquez-Semadeni et al. 1995; Koyama & Inutsuka 2002, 2004; Kritsuk & Norman 2002a, 2002b; Hennebelle & Audit 2005) and/or the fragmentation of cold clouds crushed by interstellar shocks (Nakamura et al. 2006). Apart from the formation process, we can extract a great deal of information about the surrounding ISM via investigation of the evolution and statistics of the clouds. Here we thus study the evaporation process of the tiny H I clouds, which are key quantities characterizing the fate of the clouds.

One of the simplest models is that of clouds that consist of the cold neutral medium (CNM) surrounded by the warm neutral medium (WNM) under pressure equilibrium. The two phases are thermally stable, balancing the heating rate with the cooling rate. This is based on a widely accepted two-phase description of the ISM (Field et al. 1969; Wolfire et al. 2003). The evolution of clouds is therefore described as the motion of the interface, or front, between the CNM and WNM, driven by the thermal conduction. On the basis of this picture, Zel'dovich & Pikel’ner (1969) and Penston & Brown (1970) considered isobaric, steady motion of fronts in plane-parallel geometry. They clarified that the motion is determined by pressure and that there is a saturation pressure for which a static front can exist. Their work has been substantially extended by many authors (e.g., Elphick et al. 1991, 1992; Ferrara & Shchekinov 1993; Hennebelle & Pérault 1999).

Compared to the analysis of the frontal motion in plane-parallel geometry, the evolution of spherical clouds has been much less analyzed. Graham & Langer (1973) extended the work by Zel’dovich & Pikel’ner (1969) and Penston & Brown (1970) to numerically compute isobaric flows in three-dimensional spherical geometry. Cowie & McKee (1977) and McKee & Cowie (1977, hereafter MC77) considered a model of spherical clouds surrounded by the hot ionized medium and derived an analytic formula of the evaporation rate of clouds. MC77 also considered cold clouds surrounded by the WNM, which can be directly compared with the work by Graham & Langer (1973). Nagashima et al. (2005) considered the growth of spherical clouds and showed a new way to obtaining approximate analytic formula of the evaporation rate. In this Letter we show the evaporation rate for a realistic cooling function by using a full numerical simulation and compare it with the approximate analytic formula and with the result of MC77.

Our aim is thus to estimate the evaporation timescale of cold H I clouds as a function of the cloud size and the pressure of the ambient WNM. This Letter is outlined as follows. In § 2 we compute the evaporation rate of cold clouds by using a full numerical simulation. In § 3 we show the evaporation timescale. In § 4 we discuss the dependence of the timescale on the size and pressure. We derive a new description of the evaporation rate as an improvement on MC77. In § 5 we provide conclusions.

1 INTRODUCTION

Recently discovered “tiny H I clouds” have very small sizes, \(10^{-2} \text{ pc}\), and small H I column densities, \(10^{18} \text{ cm}^{-2}\) (Braun & Kanekar 2005; Stanimirović & Heiles 2005). This new population of clouds is providing a challenge to our understanding of the evolution of the interstellar medium (ISM). The origin of the clouds is still under debate, but it is presumably the thermal instability (Field 1965; Balbus 1986) in turbulent gas (e.g., Vázquez-Semadeni et al. 1995; Koyama & Inutsuka 2002, 2004; Kritsuk & Norman 2002a, 2002b; Hennebelle & Audit 2005) and/or the fragmentation of cold clouds crushed by interstellar shocks (Nakamura et al. 2006). Apart from the formation process, we can extract a great deal of information about the surrounding ISM via investigation of the evolution and statistics of the clouds. Here we thus study the evaporation process of the tiny H I clouds, which are key quantities characterizing the fate of the clouds.

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2 EVAPORATION RATES

Below we assume for simplicity that clouds are spherically symmetric, so that we compute only radial structure of clouds. Thus, the basic fluid equations are written as

\[
\frac{\partial \rho}{\partial t} + \frac{1}{r} \frac{\partial}{\partial r} r \rho v = 0, 
\]

(1)

\[
\frac{\partial v}{\partial t} + v \frac{\partial v}{\partial r} = - \frac{1}{\rho} \frac{\partial p}{\partial r},
\]

(2)

\[
\frac{\partial \rho e}{\partial t} + \frac{1}{r^2} \frac{\partial}{\partial r} r^2 \rho e v + \frac{p}{r^2} \frac{\partial r^2 v}{\partial r} = -\rho \mathcal{C} + \frac{1}{r^3} \frac{\partial}{\partial r} r^3 \kappa (T) \frac{\partial T}{\partial r},
\]

(3)

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and the equation of state, $p = \rho k_B T$, where $\kappa(T) = 2.5 \times 10^3 \sqrt{T}$ ergs cm$^{-1}$ K$^{-1}$ s$^{-1}$ is the conductivity for neutral gas (Parker 1953), $k_B$ is the Boltzmann constant, $\mu$ is the mean molecular mass so that $\rho = \mu n$, and $L$ is the heat-loss function defined as $\rho L \equiv n^2 \Delta - n \Gamma$, where $\Delta$ and $\Gamma$ are the cooling and heating rates, respectively. We adopt a simple analytic fitting function given by Koyama & Inutsuka (2002),

$$\frac{\Lambda}{\Gamma} = 10^7 \exp \left( -\frac{118,400}{T + 1000} \right) + 1.4 \times 10^{-2} \sqrt{T} \exp \left( -\frac{92}{T} \right) \text{ cm}^3,$$

and $\Gamma = 2 \times 10^{-26}$ ergs s$^{-1}$, which take into account many processes (Koyama & Inutsuka 2000). Throughout this Letter we ignore the effect of magnetic fields. This might alter the processes (Koyama & Inutsuka 2000). Throughout this Letter we ignore the effect of magnetic fields. This might alter the processes (Koyama & Inutsuka 2000).

Before performing full numerical simulations, we show the results for a quasi–steady state (QSS), under which we transform the time derivative to the spatial one, $\partial \phi \partial t = -\mathbf{R} \partial \phi \partial r$, where $R_c$ is the cloud size and $\mathbf{R}$ is the velocity of the cloud size changing. Then the above equations become a set of ordinary differential equations. In Figure 1 several snapshots of evaporation rates, $M(r) = 4 \pi r^2 \rho(r) v(r)$, are shown by thin dashed lines for the case of the saturation pressure, $p = p_{sat} \approx 2823 k_B$. In contrast to the case of plane-parallel geometry, the cloud evaporates by the heat conduction and the size decreases as time passes. The motion of the front shown in this figure is purely driven by the curvature effect (Nagashima et al. 2005; see also § 4).

Next we show the results of full-numerical simulations. The code is based on the second-order Godunov method (van Leer 1977) in the Lagrange coordinate. We impose a constant pressure at the outer boundary. The initial structure of the cloud is computed by the QSS model with a size of about 0.1 pc. Correctly giving the initial condition is very important for this computation. If the deviation of the initial condition from the correct structure is large, spherical sound waves emerge and we could not extract the information on the evaporation. Several snapshots at which the size of the cloud becomes the same as that shown for the QSS are indicated by the thin solid lines in Figure 1. The results given by the full simulation agree well with those given by assuming the QSS, probably because the Mach number of the flow is much below unity. Therefore, we can consider that the QSS is a good assumption for the description of the evaporation of clouds.

The thick solid curve in Figure 1 denotes the evaporation rate as a function of the cloud size, $M(R_c)$, where the edge of the cloud is defined as a radius at which the density becomes $(n_{CSM} + n_{WNM})/2$. We have confirmed that the rate is proportional to the size, $M(R_c) \propto R_c$. In fact, this is the same scaling against the size as that given by MC77. Their evaporation rate is

$$M_{\text{MC77}} = \frac{16 \pi \mu k(T_c) R_c}{5 k_B} = 1.3 \times 10^{15} T_{15}^{1/2} \left( \frac{R_c}{\text{pc}} \right) \text{ g s}^{-1},$$

where $T_c$ is the temperature of the WNM and $T_c = 6.4 \times 10^4$ K for the adopted heat-loss function. This rate is plotted by the thick dashed curve. When we write our evaporation rate in a similar way,

$$M = 3.1 \times 10^{14} T_{15}^{1/2} \left( \frac{R_c}{\text{pc}} \right) \text{ g s}^{-1};$$

therefore, this is about a factor of 4 smaller than that of MC77, in spite of the same dependence on the size. The form of the evaporation rate is discussed in § 4 in a different way from MC77. Note that the above rate is valid only for the case of the saturation pressure.

3. EVAPORATION TIMESCALE

Hereafter we define the evaporation timescale as $t_{\text{evap}} \equiv M/M$, where the cloud mass $M$ is estimated by integrating the mass density within the cloud radius, $R_c$; therefore $M_c \approx 4 \pi p_{\text{CSM}} R_c^3$. Figure 2 shows the evaporation timescale against the cloud size $R_c$. The solid straight line indicates $t_{\text{evap}}$ for the case of the saturation pressure, $p = p_{sat}$, and for the QSS. The crosses on the solid line denote the result from the full numerical simulation. Similar to Figure 1, both are in excellent agreement. It is easily confirmed that $t_{\text{evap}} \propto R_c^2 \propto M^{2/3}$ for $p = p_{sat}$ reflecting the previous result $M \propto R_c$ and $M \propto R_c$. Here it should be worth noting that tiny H i clouds with the size of $\sim 10^{-2}$ pc must disappear in $\sim 1$ Myr irrespective of the ambient pressure.

The dashed and dot-dashed lines indicate the evaporation timescale for the cases of $p = 2000$ and 4000 K cm$^{-3}$, respectively. Crosses on the lines also denote the results from the full numerical simulations. Again we can see the good agreement. Note that clouds always evaporate for $p \leq p_{sat}$, but large clouds can grow for $p \geq p_{sat}$, which is shown by the thin dot-dashed line at $R \approx 0.1$ pc. This means that there exists a critical radius for growth, $R_{\text{crit}}$. Hereafter we also use $t_{\text{evap}}$ as the growth timescale. For these pressures, the size dependence of $t_{\text{evap}}$ is partly different from that for the saturation pressure. While it is very similar to that for the saturation pressure, $t_{\text{evap}} \propto R_c^2$ at $R_c \ll R_{\text{crit}}$, it becomes $t_{\text{evap}} \propto R_c$ at $R_c \gg R_{\text{crit}}$. This
feature is different from the result of MC77. In the next section, we explain the reason for the difference.

4. DISCUSSION

Below we derive an approximate evaporation rate in the same way as that of Nagashima et al. (2005). Transforming the energy equation (3) by assuming the isobaric QSS evolution and steep changes of temperature and density only at the front, we obtain the fluid velocity \( u_R \equiv v - \dot{R} \), with respect to the front rest frame,

\[
 u_R = u_R(p) + \frac{\gamma - 1}{\gamma} \frac{\mu}{k_b R} \frac{\kappa_R}{\rho_R},
\]

where the subscript \( R \) stands for values at the front, \( v = R \), and \( u_R(p) \) is the fluid velocity for the plane-parallel geometry that depends only on pressure.

To obtain a formula for the size evolution, we need a further assumption that a convergence factor from \( u \) to \( \dot{R} \) is independent of geometry, \( f \equiv -\dot{R}/u \). This is exactly the case for the thin-front limit. For three-dimensional spherical geometry, from the mass-flux conservation, we obtain \(-4\pi R^2 \rho_{\text{CNM}} \dot{R}_c = 4\pi R^2 \rho_k u_R \), where \( v = 0 \) inside the cloud. Then \( f_{\text{sphere}} = (R^3 \rho_k)/(r^2 \rho_{\text{CNM}}) \). For the plane-parallel geometry, \(-\rho_{\text{CNM}} \dot{c} = \rho_k u_R \), where \( c \) is the speed of the front with respect to the cloud (CNM) rest frame, which corresponds to \( \dot{R} \). Then \( f_{\text{plane}} = \rho_k/\rho_{\text{CNM}} \). If the front is sufficiently thin, and therefore the density \( \rho \) is \( \rho_{\text{CNM}} \) even just inside the front, we obtain \( f_{\text{plane}} \rightarrow f_{\text{sphere}} \). In reality, however, the front has a nonzero thickness. So the same \( f \) is an approximation.

Finally, we thus obtain the size evolution

\[
 \dot{R}_c = c(p) - f \frac{\gamma - 1}{\gamma} \frac{\mu}{k_b R} \frac{\kappa_R}{\rho_R},
\]

where we explicitly write the argument of \( c, p \), to stress that it depends only on pressure. For \( p > p_{\text{sat}} \), \( c(p) > 0 \), and vice versa. This equation tells us the existence of a critical radius at which \( \dot{R}_c = 0 \) when \( c(p) > 0 \),

\[
 R_{\text{crit}} = f \frac{\gamma - 1}{\gamma} \frac{\mu}{k_b R} \frac{\kappa_R}{\rho_R}.
\]

The evaporation rate is obtained from the above equation,

\[
 M_c = -4\pi R^2 \rho_{\text{CNM}} c(p) \propto R^3;
\]

\[
 M_c = 4\pi R^2 \rho_{\text{CNM}} f \frac{\gamma - 1}{\gamma} \frac{\mu}{k_b R} \frac{\kappa_R}{\rho_R} \propto R_c.
\]

From the different dependences on \( R_c \), we can see that the pressure term dominates for large clouds, \( R_c \gg R_{\text{crit}} \), and that the curvature term does for small clouds, \( R \ll R_{\text{crit}} \). Using \( f = \rho_k/\rho_{\text{CNM}} \) and \( \gamma = 5/3 \), we obtain

\[
 M_c = \frac{16\pi\mu k_R R^2}{5 k_b}.
\]

Thus, the difference from the result of MC77 for the curvature term is

\[
 \frac{M_c}{M_{M77}} = \frac{\kappa_R}{\kappa(T)} = \sqrt{\frac{T_g}{T_{g,0}}},
\]

Thus, the difference between our result and that of MC77, which should be caused by the fact that they neglected the effect of radiative cooling that compensates for the plane-parallel term of the thermal conductivity, can be partly explained in addition to the thin-front approximation to derive the curvature term. We should also note that MC77 took into account only the curvature term.

The size dependence of the evaporation timescale is thus dependent on the size. For small clouds, \( R \ll R_{\text{crit}} \), the timescale is \( \tau_{\text{evap}} \sim M_c/M \propto R^3 \). For large clouds, \( R_c \gg R_{\text{crit}} \), it is \( \tau_{\text{evap}} \sim M_c/M_c \propto R_c \). Thus, the size dependence of the evaporation timescale shown in Figure 2 is well explained by the above approximate evaporation rate we derived.

Our choice of the thermal conduction coefficient is based on the formula for neutral particles. If the ionization degree is sufficiently high, a reduced value should be more appropriate. Here we simply note that \( M_c \) and \( R_{\text{crit}} \) scale linearly to the value of \( \kappa \).

A further important point is the relationship between the critical radius, \( R_{\text{crit}} \), and the Jeans length for gravitational instability, \( R_J \equiv \sqrt{4\pi G \rho_{\text{sat}}}/c_s^2 \), where \( c_s \) is the sound velocity. Figure 3 shows \( R_{\text{crit}} \) and \( R_J \) as a function of the ambient pressure.

We find that there can exist clouds growing only by condensation for the case of \( p > p_{\text{sat}} \). The timescale until the cloud size exceeds \( R_J \) is, however, quite long, as shown in Figure 2.

5. CONCLUSION

We have investigated the evaporation rate and timescale of tiny H I clouds by using numerical simulations. We confirmed that the results are almost recovered by assuming the evolution under the quasi-steady state because the fluid velocity is much slower than the sound velocity. We have found that clouds with a size of about 0.01 pc evaporate in approximately 1 Myr almost
Fig. 3.—Critical radius $R_{\text{crit}}$ plotted by the solid line against pressure of the WNM. Jeans length $R_{J}$ is also plotted by the dashed line. [See the electronic edition of the Journal for a color version of this figure.]

independent of pressure of the ambient WNM. This suggests that there might be some mechanisms to continuously form tiny H i clouds, or, that the ambient pressure around the clouds is much higher than the standard ISM pressure, if their existence is ubiquitous. The evaporation timescale of clouds larger than 0.1 pc, however, depends strongly on the pressure. Clouds larger than a critical radius can even grow if the pressure is higher than the saturation pressure.

In order to physically understand the simulation results, we derived an analytic formula for evaporation by assuming the isobaric evolution and a structure of the interface independent of geometry. The obtained evaporation rate has two separate terms: pressure and curvature terms. The former is independent of geometry, that is, it emerges even in analysis for plane-parallel geometry, and it becomes zero for the saturation pressure. The latter is proportional to the size of clouds, that is, the curvature. We have found that the obtained evaporation rate is a natural extension of an evaporation rate obtained by MC77, which corresponds to our curvature term, but the rate of MC77 rate is a little higher than ours. For lower pressure than the saturation pressure, the signs of the two terms are the same. Therefore, clouds always evaporate. For higher pressure, on the other hand, the signs are different. Therefore, a critical size exists and the size of clouds determines whether they grow or evaporate. We have confirmed that the simulation results show the same dependence on the size as the analytic formula.

We would like to stress that our analysis can be adapted irrespective of the origin of the tiny H i clouds. In order to obtain information on the formation process, statistical properties such as mass function of clouds should be required. This will be done in future work.

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