Photoevaporation of Minihalos During Cosmic Reionization: Primordial and Metal-enriched Halos

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Abstract

The density distribution of the intergalactic medium is an uncertain but highly important issue in the study of cosmic reionization. It is expected that there are abundant gas clouds hosted by low-mass “minihalos” in the early universe, which act as photon sinks until being photoevaporated by the emerging ultraviolet background (UVB) radiation. We perform a suite of radiation-hydrodynamic simulations to study the photoevaporation of minihalos. Our simulations follow hydrodynamics, nonequilibrium chemistry, and the associated cooling processes in a self-consistent manner. We conduct a parametric study by considering a wide range of gas metallicities \( (0 < Z_e < Z \lesssim 10^{-3} Z_\odot) \), halo mass \( (10^5 M_\odot \lesssim M \lesssim 10^8 M_\odot) \), UVB intensity \( (0.01 \lesssim J_\gamma \lesssim 1) \), and turn-on redshift of ionizing sources \( (10 \lesssim z_{\text{fin}} \lesssim 20) \). We show that small halos are evaporated in a few tens of millions of years, whereas larger mass halos survive 10 times longer. The gas mass evolution of a minihalo can be characterized by a scaling parameter that is given by a combination of the halo mass, background radiation intensity, and redshift. Efficient radiative cooling in metal-enriched halos induces fast condensation of the gas to form a dense, self-shielded core. The cold, dense core can become gravitationally unstable in halos with high metallicities. Early metal enrichment may allow star formation in minihalos during cosmic reionization.

Unified Astronomy Thesaurus concepts: Hydrodynamical simulations (767); Metallicity (1031); Star formation (1569); Chemical enrichment (225); Reionization (1383)

Supporting material: machine-readable table

1. Introduction

A broad range of observations, including cosmic microwave background radiation anisotropies, have established the so-called standard cosmological model in which the universe mainly consists of two unknown substances called dark energy and dark matter. According to the standard model, small primeval density fluctuations generated in the very early universe grow by gravitational instability, to form nonlinear objects called dark matter halos.

Dark halos with mass of \( \sim 10^5 - 10^6 M_\odot \) are thought to be the birthplace of the first generation of stars (Tegmark et al. 1997; Yoshida et al. 2003). Star-forming gas clouds are formed through condensation of primordial gas by molecular hydrogen cooling. Population III stars are formed in the primordial gas clouds when the age of the universe is a few tens/hundreds of million years old (e.g., Abel et al. 2002; Naoz et al. 2006; Fialkov et al. 2012).

Ultranarrow (UV) and X-ray radiation from the first stars and their remnants ionize and heat the intergalactic medium (IGM), to initiate cosmic reionization. Recent observations including the measurement of the electron scattering optical depth (e.g., Planck Collaboration et al. 2016), the Gunn–Peterson trough in quasar spectra (Gunn & Peterson 1965; Fan 2006; Bañados et al. 2018), and Ly-\( \alpha \) emission from star-forming galaxies (Mason et al. 2018) suggest that the process of reionization completes by \( z \sim 6 \) (e.g., Naidu et al. 2020; Weinberger et al. 2020). It is expected that early stages of reionization can be directly probed by future radio telescopes such as the Square Kilometre Array (SKA; Koopmans et al. 2015) and the Hydrogen Epoch of Reionization Array (HERA; DeBoer et al. 2017) through observation of redshifted 21 cm emission from neutral hydrogen in the IGM (see, e.g., Barkana 2016; Mesinger 2019, for a recent review).

In the early phase of reionization, the density distribution of the IGM, or the so-called gas clumping, is an important factor that critically sets the UV photon budget necessary for reionization. Dense gas clouds hosted by cosmological minihalos can be significant photon sinks, and their existence and abundance affect the process and duration of reionization. Unfortunately, it is nontrivial to derive the abundance of the gas clouds and to estimate the effective gas clumping factor, because small gas clouds are effectively photoevaporated by the emerging UV background radiation. At the same time, stars can be formed in the gas clouds, which then act as UV photon sources.

A number of studies have investigated early star formation under various environments (e.g., Low & Lynden-Bell 1976; Omukai 2000; Schneider et al. 2002, 2003; Glover & Jappsen 2007; Jappsen et al. 2007, 2009a, 2009b; Smith et al. 2009; Chiaki et al. 2018; Chiaki & Wise 2019). Metal enrichment affects the evolution of gas clouds and subsequent star formation process through enhanced radiative cooling by heavy element atoms and dust grains (e.g., Omukai et al. 2005; Hartwig & Yoshida 2019). Although details of low-metallicity star formation have been explored by recent numerical simulations (Chiaki et al. 2018), the effect of metal enrichment on halo photoevaporation has not been systematically studied.
It is important to study the evolution of minihalos with a wide range of metallicities in order to model the physical process of cosmic reionization in a consistent manner.

Reionization begins as a local process in which an individual radiation source generates an HII region around it. Shapiro (1986), Shapiro & Giroux (1987), and Donahue & Shull (1987) used a one-dimensional model under spherical symmetry to study the ionization front (I-front) propagation through the IGM in a cosmological context. Radiative transfer calculations have been performed in a post-processing manner by using the density field realized in cosmological simulations (Abel et al. 1999; Razoumov & Scott 1999; Sokasian et al. 2001; Cen 2002; Hayes & Norman 2003). Fully coupled radiation-hydrodynamic simulations have been used to study reionization in a cosmological volume (Gnedin 2000; Gnedin & Abel 2001; Ricotti et al. 2002; Susa & Umemura 2004a, 2004b; Wise et al. 2014; Xu et al. 2016). Recent simulations of reionization explore the process of cosmic reionization employing large cosmological volumes of ~10^5 comoving Mpc^3 while resolving the first galaxies in halos with mass of ~10^8 M☉ (Semelin et al. 2017; Ocvirk et al. 2020). However, even these state-of-the-art simulations do not fully resolve minihalos nor are they able to follow the process of photoevaporation, and thus it still remains unclear how the small-scale gas clumping affects reionization.

Shapiro et al. (2004) and Iliev et al. (2005) studied the dynamical evolution of minihalos irradiated by UV radiation. They performed 2D radiation-hydrodynamic simulations including the relevant thermochemical processes. They explore a wide range of parameters such as halo mass, redshift, and the strength of UV radiation. It is shown that gas clumping at subkiloparsec scales dominates the absorption of ionizing photons during the early phase of reionization. Park et al. (2016) ran radiation-hydrodynamic simulations starting from realistic cosmological initial conditions to study the feedback of an isotropic UV background on small-scale structures. They showed that recombination per hydrogen atom can be as large as or larger than one, and that the net photon budget for cosmic reionization can be increased. A more recent study by D’Aloisio et al. (2020) examined the impact of small-scale structures on reionization using a suite of radiation-hydrodynamic simulations. An important question is whether or not the minihalos survive under a strong UVB for a long time, over a significant fraction of the age of the universe. Another interesting question is whether or not stars are formed in metal-enriched minihalos. If massive stars are formed, they also contribute to reionization and may thus imprint characteristic features in the 21 cm signals (Cohen et al. 2016).

In the present paper, we perform a large set of high-resolution radiation-hydrodynamic simulations of minihalo photoevaporation. We aim at investigating systematically the effects of metal enrichment on the photoevaporation. We evaluate the characteristic photoevaporation time and study its metallicity dependence. We also develop an analytic model of photoevaporation and compare the model prediction with our simulation result. The rest of the paper is organized as follows. In Section 2, we explain the details of our computational methods. The simulation results are presented in Section 4. We develop an analytical model that describes the physical process of minihalo photoevaporation in Section 3. We discuss the physics of minihalo photoevaporation in Section 5. Finally, we provide a summary and concluding remarks in Section 6.

Throughout the present paper, we assume a flat Lambda cold dark matter cosmology with (Ωm, ΩΛ, h) = (0.27, 0.73, 0.046, 0.7) (Komatsu et al. 2011). All the physical quantities in the following are given in physical units.

2. Numerical Simulations

We perform radiation-hydrodynamic simulations of minihalo photoevaporation by external UV radiation. We run a set of simulations systematically by varying the gas metallicity, dark matter halo mass, intensity of the radiation, and redshift.

Our simulation setup is schematically shown in Figure 1.

The numerical method is essentially the same as in Nakatani & Yoshida (2019), where we studied the dynamical evolution of molecular gas clouds exposed to an external UV radiation field. Briefly, we use a modified version of PLUTO (version 4.1; Mignone et al. 2007), which incorporates ray-tracing radiative transfer of UV photons and nonequilibrium chemistry. The details of the implemented physical processes are found in Nakatani et al. (2018a, 2018b).

The simulations are configured with 2D cylindrical coordinates. The governing equations are

\[ \frac{\partial \rho_b}{\partial t} + \nabla \cdot (\rho_b \mathbf{v}) = 0, \]  
\[ \frac{\partial \rho_b v_R}{\partial t} + \nabla \cdot (\rho_b v_R \mathbf{v}) = - \frac{\partial P}{\partial R} - \rho_b \frac{\partial \Phi}{\partial R}, \]  
\[ \frac{\partial \rho_b v_x}{\partial t} + \nabla \cdot (\rho_b v_x \mathbf{v}) = - \frac{\partial P}{\partial x} - \rho_b \frac{\partial \Phi}{\partial x}, \]
\[ \frac{\partial \rho_b}{\partial t} + \nabla \cdot (\rho_b \mathbf{v}) = -\rho_b \mathbf{v} \cdot \nabla \Phi + \rho_b (\Gamma - \Lambda), \]  
\[ \frac{\partial n_{H_2}}{\partial t} + \nabla \cdot (n_{H_2}) \mathbf{v} = n_{H_2} C_i, \]
where \( R \) and \( x \) are the physical radial distance and vertical height, respectively, and \( t \) is the proper time in the frame of the halo. We denote the gas density, velocity, and pressure as \( \rho_b, \mathbf{v}, \) and \( P. \) In the equation of motion, \( \Phi \) is the external gravitational potential of the host dark halo. In the energy equation, \( E \) and \( H \) are the total specific energy and total specific enthalpy, and \( \Gamma \) and \( \Lambda \) are the specific heating rate and specific cooling rate. The abundance of ith chemical species, \( y_i, \) is defined by the ratio of the species’ number density \( n_i \) to the hydrogen nuclei number density \( n_H. \) The total reaction rate is denoted as \( C_i. \) The gas is composed of eight chemical species: \( \text{H}, \text{H}^+, \text{H}_2, \text{H}_2^+, \text{He}, \text{CO}, \text{C}^+, \text{O}, \) and \( e^+. \) The elemental abundances of carbon and oxygen are normalized by the assumed metallicity \( Z \) as \( 0.926 \times 10^{-4} Z/Z_c \) and \( 3.568 \times 10^{-4} Z/Z_c, \) respectively (Pollack et al. 1994; Omukai 2000).

We consider halos with a wide range of masses of \( 10^{9} M_{\odot} \leq M \leq 10^{10} M_{\odot}. \) Each halo has a Navarro–Frenk–White density profile scaled appropriately (Navarro et al. 1997). We assume that the halo potential is fixed and is given by a function of the spherical radius \( r \equiv \sqrt{R^2 + x^2} \) as
\[ \rho_{\text{DM}} \propto \frac{1}{c_N \xi (1 + c_N e^3)^2}, \]
where \( c_N \) is the concentration parameter, and \( x \) is the spherical radius normalized by the virial radius, i.e., \( \xi \equiv r/r_{\text{vir}}. \) The virial radius of a halo collapsing at redshift \( z \) is
\[ r_{\text{vir}} = 1.51 \left( \frac{\Omega_m h^2}{0.141} \right)^{-1/3} \left( \frac{M}{10^9 M_{\odot}} \right)^{1/3} \right] \times \left( \frac{\Delta_c}{18\pi^2} \right)^{-1/3} \left( 1 + \frac{e^c}{10} \right)^{-1} \text{kpc}, \]
where \( \Delta_c \) is the overdensity relative to the critical density of the universe at the epoch. We adopt \( \Delta_c = 18\pi^2. \) Then the halo potential \( \Phi \) is explicitly given by
\[ \Phi(r) = -V_c^2 \ln (1 + c_N \xi) \frac{c_N}{c_N \xi} \ln (1 + c_N) - \frac{c_N}{1 + c_N}, \]
\[ V_c \equiv \left( \frac{GM}{r_{\text{vir}}} \right), \]
\[ T_{\text{vir}} = \frac{GM \mu m_p}{2r_{\text{vir}} k_B}, \]
where \( \mu \) is mean molecular weight, \( m_p \) is the proton mass, and \( k_B \) is the Boltzmann constant. The initial density profile is
\[ \rho_b(r) = \hat{\rho}_b \exp \left[ -\frac{\Phi}{k_B T_{\text{vir}} / \mu m_p} \right], \]
where \( \hat{\rho}_b \) is the normalization factor
\[ \hat{\rho}_b \equiv \frac{M \Omega_b \Omega_m^{-1}}{\int_0^c \left( \frac{e}{h \nu_T / \mu m_p} \right) \exp \left( -\frac{e}{k_B T_{\text{vir}} / \mu m_p} \right) \nu \, d\nu}, \]
with this normalization, the mass ratio of baryons to dark matter within \( r_{\text{vir}} \) equals the global cosmic baryon fraction \( f_b = \Omega_b / \Omega_m. \) The initial density profile is specified by \( M, c_N, \) and \( \xi. \) Note that \( \hat{\rho}_b \) is not the central density but is a geometry-weighted average density, which is independent of \( M \) but scales as \( \propto (1+z)^3 \) for fixed \( \Delta_c \) and \( c_N. \) Adopting \( c_N = 10, \) \( \hat{\rho}_b \approx 7.2 \times 10^{-28} (1+z)/11 ^2 \text{g cm}^{-3}. \)

The initially fully atomic gas in the halo is exposed to plane-parallel UV radiation as illustrated in Figure 1. We follow photoionization by extreme UV (EUV; 13.6 eV < \( \nu < 100 \text{eV} \) photons and photodissociation by far-UV (FUV) photons in the Lyman–Werner (LW) band (11.2 eV < \( \nu < 13.6 \text{eV} \)). The UV spectrum is given by
\[ J(\nu) = J_{21} \left( \frac{\nu}{\nu_{11}} \right)^{-\alpha} \times 10^{-21} \text{erg} \, \text{s}^{-1} \, \text{cm}^{-2} \, \text{Hz}^{-1} \, \text{sr}^{-1}, \]
where \( \nu_{11} \) is the Lyman limit frequency (i.e., \( h\nu_{11} = 13.6 \text{eV}. \) We set the UV spectral slope \( \alpha = 1 \) and consider \( J_{21} \) in the range of \( 0.01 \leq J_{21} \leq 1 \) (Thoul & Weinberg 1996). We calculate the photodissociation rate with taking into account the self-shielding of hydrogen molecules (Draine & Bertoldi 1996; Lee et al. 1996). The self-shielding factor may need to be corrected if the velocity gradient of the bulk motion is significant within \( H_2 \)-rich regions, because the LW resonance wavelength can be shifted more than a thermal Doppler width. In our study, the bulk motion of the molecular gas is driven by the halo’s gravity, and therefore a correction could be necessary near the central region where the \( H_2 \) formation time and the freefall time are the shortest. However, the Sobolev length (Sobolev 1957) of the infalling gas is estimated to be \( \sim 0.1 r_{\text{vir}} (T/r_{\text{vir}})^{0.5} \) at the center, and it increases with the radial distance. The local Sobolev length is much larger than the length scale over which the density significantly varies, i.e., \( d(\ln \rho_b)/d\nu. \) Typically, the density decreases by about two orders of magnitude at a distance of \( \sim 0.1 r_{\text{vir}}. \) Since \( H_2 \) formation is efficient in the central compact region within which the Doppler shift due to the bulk motion can be ignored, we do not consider the minor correction to the self-shielding factor in the present study.

Heating and cooling rates are calculated self-consistently with the nonequilibrium chemistry model of Nakatani & Yoshida (2019). The major processes are photoionization heating, Lyα cooling, radiative recombination cooling, C II line cooling, O I line cooling, \( H_2 \) line cooling, and CO line cooling. The corresponding heating/cooling rates are found in Nakatani et al. (2018a, 2018b). FUV-induced photoelectric heating is not effective with the FUV intensity and metallicity of interest in this study. We omit it from our thermochemistry model. For the present study, we also implement the Compton cooling by the cosmic microwave background (CMB) photons interacting with free electrons with physical density \( n_e \) as
\[ \Lambda_{\text{Comp}} = 5.65 \times 10^{-36} n_e (1+z)^4 (T - T_{\text{CMB}}) \text{ erg cm}^{-3} \text{s}^{-1}. \]
where \( T_{\text{CMB}} \) is the CMB temperature given by \( T_{\text{CMB}} = 2.73(1+z) \) K.

Our computational domain extends 0 kpc \( \leq R \leq r_{\text{vir}} \) and \( -r_{\text{vir}} \leq x \leq r_{\text{vir}} \). We define computational grids uniformly spaced with the number of grids \( N_R \times N_x = 320 \times 640 \). UV photons are injected from the boundary plane at \( x = -r_{\text{vir}} \). We assume that the cosmological I-front arrives at the plane at a redshift of \( z_{\text{IN}} \). All our runs start at this time denoted as \( t = 0 \). Note that the external halo potential is fixed; we do not consider growth of halo mass. We discuss potential influences of this simplification in Section 5.

We run a number of simulations varying three parameters in the range of \( 0 \leq Z \leq 10^{-3} \), \( 0.01 \leq J_{21} \leq 1 \), \( 10^5 M_0 \leq M \leq 10^9 M_0 \), and \( 10 \leq z_{\text{IN}} \leq 20 \), respectively. A total of 495 \( (=5 \times 3 \times 11 \times 3) \) simulations are performed. We also perform an additional six simulations where three different spectral shapes are adopted. The results of the additional runs are shown in Section 5.5. Hereafter, we dub each run based on the assumed values of the parameters. A simulation with \( (Z, J_{21}, M, z_{\text{IN}}) = (10^{-3}Z_0, 10^5, 10^9 M_0) \) is referred to as “Zab9Mcdz.” For example, \( Z \sim 10 M_5 \) indicates \( Z = (Z_0, 1, 10^5 M_0) \).

### 3. Photoevaporation Model

Before presenting the results from a number of simulations, it would be illustrative to describe the basic physics of photoevaporation. It also helps our understanding of the overall evolution of a UV-irradiated minihalo. To this end, we develop an analytic model and evaluate the photoevaporation rate numerically so that it can be compared directly with our simulation results.

Photoevaporation is driven by gas heating associated with photoionization. Incident UV radiation ionizes the halo gas, forming a sharp boundary between the ionized and neutral (self-shielded) regions. Photoevaporative flows are launched from the outermost layer of the self-shielded region, as schematically shown in Figure 2. The number of UV photons incident on the self-shielded region is equal to that of evaporating gas particles. The gas mass evolution of a halo can be described as

\[
\frac{dM_i}{dr} = -m \int_{\partial V_i} dS_i \cdot \hat{x} J = -m \int_{V_i} dV \nabla \cdot \hat{x} J = -m \int_0^{R_i} 2\pi R J (R, x_i) dR,
\]

where \( M_i \) is the total gas mass in the self-shielded region, \( m \) is the total gas mass in the self-shielded region, \( V_i \) is the volume of the self-shielded region, \( \hat{x} \) is a unit vector in the \( x \)-direction, \( J \) is total photon number flux, \( R_i \) is the maximum radial extent of the self-shielded region (i.e., the radial position of the I-front), and \( x_i(R) \) gives the locus of the I-front on the \( R \)-\( x \) plane.

We define the following dimensionless quantities: \( \tilde{\xi} = (\xi, \xi_R) \equiv (x, r_{\text{vir}}, R_{\text{vir}}, t/t_{\text{vir}}) \), \( \tilde{J} = J/\langle J \rangle \), \( M = M_i/\langle m \rangle \), and \( J = J/\langle J \rangle \), and rewrite Equation (15) in a dimensionless form

\[
\frac{d\tilde{M}_i}{d\tilde{r}} = -\tilde{J}_0 \tilde{r}_i \frac{m}{\tilde{f}_0} \int_0^{\tilde{\xi}_i} 2\pi \tilde{R} \tilde{J}(\tilde{\xi}_R, \tilde{\xi}_x, \tilde{\xi}_R) d\tilde{\xi}_R.
\]

The ionizing photon number flux at the I-front, \( J(R, x_i) \), is equal to \( J_0 \) minus the total recombinations along the ray up to the I-front. Note that \( J(R, x) \) depends on the recombination coefficient and the density, and also on the velocity profile of the photoevaporative flows.

In our analytic model, we approximate the I-front to be a hemispheric surface with radius \( R_i \) facing toward the incident radiation in the region \( x < 0 \). In the other region \( x \geq 0 \), the I-front lies at the surface of a cylinder with radius \( R_p \). Our model further assumes that photoevaporative flows are spherically expanding in \( x < 0 \) and that the ionized gas is isothermal with \( T = 10^5 \) K. The wind velocity is assumed to be the sound speed of the ionized gas, \( c_i \). The continuity equation gives the ionized gas density \( n_i \) as

\[
n_i = \frac{J(R, x_i)}{c_i} \frac{R_i^2}{x^2 + R_i^2} \frac{[x]}{\sqrt{x^2 + R_i^2}}.
\]
for \( x^2 + R^2 \geq R_\text{e}^2 \) and \( R \leq R_\text{e} \). The ionizing photon number flux at the I-front \( J(R, x_i) \) is equal to \( J_0 \) minus the sum of recombination along the line of sight

\[
J_0 - J(R, x_i) = \int_{-\infty}^{x_i} dx \, n_i^2 \alpha_B \approx \frac{\alpha_B}{c_i^2} \left( \frac{R_i}{R} \right)^4 \frac{2}{8} \left( J(J(R, x_i)) \right)^2 \times \left( \theta_{R,i} - \frac{1}{4} \sin 4\theta_{R,i} \right),
\]

where \( \alpha_B \) is the case-B recombination coefficient and

\[
\theta_{R,i} = \arccos \sqrt{1 - \left( \frac{R}{R_i} \right)^2} = \arccos \sqrt{1 - \left( \frac{\xi_R}{\xi_i} \right)^2}.
\]

For the integration in Equation (18), we have ignored the variation of \( J(R, x_i) \) over the surface of the I-front. Solving the quadratic equation with respect to \( J(R, x_i) \) yields

\[
J(R, x_i) \approx \frac{2 J_0}{1 + \sqrt{1 + \frac{\alpha_B R_i}{c_i} \Psi(\theta_{R,i})}},
\]

where

\[
\Psi(\theta_{R,i}) = \frac{1}{2 \sin^3 \theta_{R,i}} \left( \theta_{R,i} - \frac{1}{4} \sin 4\theta_{R,i} \right).
\]

The dimensionless geometrical factor \( \Psi(\theta_{R,i}) \) is a monotonically decreasing function for \( 0 \leq \theta_{R,i} \leq \pi/2 \) and varies in the range of \( \pi/4 \leq \Psi \leq \pi/3 \).

We note that diffuse EUV photons can ionize the neutral gas in the shade region \((x_i \geq 0)\); the effect we do not include in our simulations. If the recombination time at the I-front in \( x_i \geq 0 \) is longer than the sound crossing time of photoevaporative flows, i.e., \((n_i \alpha_B)^{-1} \gtrsim c_i/R_i\), the direct radiation from \( x < 0 \) would erode the surface layer in the shade region faster than the diffusion rate. photoevaporation there. The mass loss driven by the diffuse EUV radiation is negligible in this case.

If \((n_i \alpha_B)^{-1} \lesssim c_i/R_i\) at the I-front, photoevaporative flows can be excited by the diffuse field in \( x \geq 0 \). Under the assumption of axisymmetric photoevaporative flows (i.e., \( n_i = J c_i^{-1}(R_i/R) \)), dimensional analysis gives \( J \sim c_i^2/\alpha_B R_\text{e} \) for the ionizing photon number flux at the I-front, which is much smaller than that in \( x < 0 \) if \( \alpha_B R_i/c_i^2 \gg 1 \), and becomes comparable if \( \alpha_B R_i/c_i \sim 1 \). These estimates let us neglect the contribution of mass loss from \( x \geq 0 \) in our analytic model that focuses on the lifetime of minihalos.

With these results, Equation (16) reduces to

\[
\frac{dM}{dt} = -\eta J_0^{3/2} \int_{-\infty}^{x_i} \frac{4 \pi \xi R d\xi R}{1 + \sqrt{1 + q \xi R \Psi(\theta_{R,i})}} ,
\]

where

\[
\eta \equiv \frac{J_0 v_\text{e,m}}{\mathcal{J}_0 M V_\text{c}} \approx 12 J_2 \left( \frac{1 - z}{11} \right)^{1/3} \left( \frac{1 + z}{11} \right)^{-7/2},
\]

\[
q \equiv \frac{\alpha_B R_i}{c_i^2} J_0 \approx 1.5 \times 10^2 J_2 \left( \frac{1 - z}{11} \right)^{1/3} \left( \frac{1 + z}{11} \right)^{-1}.
\]

Note that the dimensionless parameter \( \eta \) effectively measures the ratio of Hubble time to the ionization timescale. The other parameter \( q \) quantifies the magnitude of UV absorption in the photoevaporative flows; absorption is negligible if \( q \ll 1 \), while it is significant if \( q \gg 1 \). Although the differential equation is not solved analytically, we can derive the asymptotic behavior of the gas mass in a few limiting cases. For \( q \gg 1 \), the right-hand-side coefficient is approximately proportional to \( \approx \eta q^{-1/2} \propto J_2^{1/2} M^{-1/2}(1 + z)^{-3} \). It follows that there is a similarity in the gas mass evolution among halos having similar \( \eta q^{-1/2} \) and initial \( \xi_i \). Hence we define the similarity parameter as

\[
\chi \equiv \eta q^{-1/2}.
\]

The initial \( \xi_i \) is determined by the radius at which the I-front turns to the D-critical type from R-type (e.g., Spitzer 1978; Bertoldi 1989). The radius is obtained by numerically solving the integral equation

\[
J_0 \int_{-\infty}^{-\xi} \rho_b^2(\xi_i, \xi) d\xi_i = \frac{J_0}{\mathcal{J}_0 M V_\text{c}} \left( 1 - \frac{2c_i n_0}{J_0} \rho_b(\xi_i) \right)
\]

\[
\rho_0 \equiv \frac{\mathcal{J}_0 M V_\text{c}}{m_\text{e} n_0} \quad \mathcal{M} \equiv \frac{J_0}{\mathcal{J}_0 M V_\text{c}} 
\]

\[
\theta_\text{e} \equiv \frac{2 c_i n_0}{J_0} 
\]

with the initial density profile of Equation (11). Typically, the initial \( \chi \) is larger for lower \( J_0 \), for larger \( M \), and for higher \( n_0 \) (i.e., higher \( z_{\text{IN}} \)). We list \( q, \eta, \chi, \theta_\text{e}, \) and \( \theta_\text{e} \) for each of our runs in Table 1.

In the above model, we have assumed a constant flow velocity, but in practice the flow is accelerated within the ionized boundary layer after launched with a small, negligible velocity. Also, we do not consider gravitational force by the host halo, which can decelerate the photoevaporative flows. In Section 4.6, we will introduce a few corrections to these simplifications and examine carefully the similarity of gas mass evolution by comparing with the simulation results.

4. Simulation Results

We first describe the dynamical evolution of a minihalo in our fiducial case, and examine the effect of metal and dust cooling in Section 4.1. Then, we focus on the photoevaporation rates and study the dependence on metallicity and on halo mass in Section 4.2. The dependence of the photoevaporation rates on radiation intensity and the turn-on redshift is studied in Sections 4.3 and 4.4. We summarize the halo mass evolution in Section 4.6. Then we provide an analytical fit to the derived mass evolution as a function of time in Section 4.7. For convenience, we term halos with \( T_\text{vir} > 10^5 \) K atomic cooling (massive) halos in the following sections. The corresponding mass range is \( M \gtrsim 10^{7.5} M_\odot \) \( (10^8 M_\odot) \) at \( z = 10 \). Lower-mass halos \( (T_\text{vir} < 10^5 \) K \) are referred to as low-mass halos; those with \( M \gtrsim 10^{6.5} \) \( 10^7 M_\odot \) \( (10^7 \) \( 10^8 M_\odot) \) at \( z_{\text{IN}} = 10 \) \( (15–20) \) are specifically called molecular cooling halos.

4.1. Photoevaporation and Metallicity Dependence

Figure 3 shows the density and temperature distributions for a halo with \( M = 10^{5.5} M_\odot \), irradiated by UV radiation with \( J_2 = 1 \) at \( z_{\text{IN}} = 10 \). We compare the results with two different
The physical parameters of Z = 0 Z⊙ and Z = 10−4 Z⊙ (i.e., Z ∝ J1M5.5z10 and Z4J1M5.5z10). At the beginning of the evolution, an R-type I-front moves in from the diffuse outer part to the denser interior, and soon forms a D-type I-front. This transition is commonly seen regardless of metallicity for given J21, Zvir, and M. The D-type front converges toward the symmetry axis and squeezes the gas regardless of metallicity for given I-front moves in from the diffuse outer part to the denser interior, producing abundant H2 molecules. The central low temperature is roughly comparable for H2, C II, and OI species in metal-rich halos, a sufficient amount of H2 molecules form via the grain-catalyzed reaction and induce gravitational collapse. The H2 abundance increases to nH2 ≈ 3 × 10−17 cm3 s−1 is the reaction rate, by the time the shock sweeps the central gas. The approximate shock crossing time of ~ rvir/ci is longer than the cooling time at the center (see also discussion in Section 4.2). Therefore, the gas density in the metal-rich halos increases through the radiative cooling before the shock compression. The D-type front causes only a minor effect on the formation of the dense core, which is driven primarily by the cooling-induced contraction.

| Label | Z (Z⊙) | J21 | M (M⊙) | zIN | t0 (Myr) | q | η | χ | χ′ | Lvir | αBR | C4 | i4 | p |
|-------|--------|-----|--------|-----|----------|---|---|---|---|-----|------|---|---|---|
| Z ∝ J0M3z10 | 0 | 1 | 103 | 10 | 76.9 | 14.9 | 119 | 30.7 | 29.5 | 1.36 | 0.444 | 0.00659(3.7%) | 0.0228(0.12%) | 3.7(0.38%) |
| Z5H0M3z10 | 10−6 | 1 | 103 | 10 | 76.9 | 14.9 | 119 | 30.7 | 29.5 | 1.36 | 0.444 | 0.00659(3.7%) | 0.0228(0.12%) | 3.7(0.38%) |
| Z5H0M3z10 | 10−5 | 1 | 103 | 10 | 76.9 | 14.9 | 119 | 30.7 | 29.5 | 1.36 | 0.444 | 0.00659(3.7%) | 0.0228(0.12%) | 3.7(0.38%) |
| Z5H0M3z10 | 10−4 | 1 | 103 | 10 | 76.9 | 14.9 | 119 | 30.7 | 29.5 | 1.36 | 0.444 | 0.00659(3.7%) | 0.0228(0.12%) | 3.7(0.38%) |
| Z5H0M3z10 | 10−3 | 1 | 103 | 10 | 76.9 | 14.9 | 119 | 30.7 | 29.5 | 1.36 | 0.444 | 0.00659(3.7%) | 0.0228(0.12%) | 3.7(0.38%) |
| Z5H0M3z10 | 0 | 0.1 | 103 | 10 | 76.9 | 1.49 | 119 | 9.7 | 9.34 | 0.136 | 4.44 | 1.94 × 10−18(−) | 0.0407(0.37%) | 2.840(0.96%) |
| Z6H1M3z10 | 10−6 | 0.1 | 103 | 10 | 76.9 | 1.49 | 119 | 9.7 | 9.34 | 0.136 | 4.44 | 1.94 × 10−18(−) | 0.0407(0.37%) | 2.840(0.96%) |
| Z6H1M3z10 | 10−5 | 0.1 | 103 | 10 | 76.9 | 1.49 | 119 | 9.7 | 9.34 | 0.136 | 4.44 | 1.94 × 10−18(−) | 0.0407(0.37%) | 2.840(0.96%) |
| Z6H1M3z10 | 10−4 | 0.1 | 103 | 10 | 76.9 | 1.49 | 119 | 9.7 | 9.34 | 0.136 | 4.44 | 2.05 × 10−18(−) | 0.0407(0.37%) | 2.840(0.96%) |
| Z6H1M3z10 | 10−3 | 0.1 | 103 | 10 | 76.9 | 1.49 | 119 | 9.7 | 9.34 | 0.136 | 4.44 | 1.75 × 10−18(−) | 0.0407(0.37%) | 2.835(0.95%) |

Note: The physical parameters t0, q, η, χ, Lvir, and αBR are defined in Section 3. The definition of χ′ is given in Section 4.6, and those of the fitting parameters C, i, and p are given in Section 4.7. Parentheses in the columns of the fitting parameters C, i, and p show the fitting errors. Hyphens in these columns mean that no fitting parameters are found with reasonable errors. Note that i4 and p are undervisible when C4 = 1.

(This table is available in its entirety in machine-readable form.)
We note that primordial gas clouds formed under strong LW radiation are suggested to be possible birthplaces of massive black holes (Omukai 2001; Shang et al. 2010; Agarwal et al. 2012; Regan et al. 2014; Hartwig et al. 2015; Regan et al. 2016a; Schauer et al. 2017). In metal-enriched halos, the gas can still cool and condense by metal and dust cooling even under strong UV radiation.

In contrast to low-mass, H2-cooling halos, Lyα cooling dominates in halos with $T_{\text{vir}} \gtrsim 20,000$ K. Since the efficiency of Lyα cooling is independent of metallicity, the gas condensates quickly in the massive halos with $M \gtrsim 10^{7.5} M_\odot$. If the gas is enriched to $Z \gtrsim 10^{-3} Z_\odot$, dust-gas collisional heat transfer causes efficient gas cooling for $n_H \gtrsim 10^7$ cm$^{-3}$ (Omukai et al. 2005, 2008; Chiaki et al. 2016).

Our findings are largely consistent with those of Wise et al. (2014) regarding the correspondence between halo mass and dominant cooling processes. Halos with $M \gtrsim 10^7 M_\odot$ are expected to be metal enriched by Population III supernovae triggered in the progenitor halo or in nearby halos, and thus metal cooling is dominant or comparable to atomic/molecular hydrogen cooling. Following Wise et al. (2014), we call such halos “metal-cooling halos,” which have masses between the atomic cooling limit ($\sim 10^{7.5} - 10^8 M_\odot$) and the upper mass limit of molecular cooling halos ($\sim 10^{6.5} - 10^7 M_\odot$). We find similarly efficient metal cooling for $M \gtrsim 10^7 M_\odot$, but the most important effect of metal enrichment in our cases is to lower the molecular cooling limit by allowing formation of H2 through grain-catalyzed reactions, especially at $Z \gtrsim 10^{-4}Z_\odot$.

The gas in massive halos ($T_{\text{vir}} > 10^4$ K) is gravitationally bound even under strong UV radiation. We find rather small mass loss in the runs with $M = 10^{7.5} M_\odot$. Approximately 10% of the initial gas mass is lost via photoevaporation, but the diffuse, ionized gas moves toward the center. This process slightly recovers the total gas mass within $r_{\text{vir}}$. For halos with $T_{\text{vir}} > 10^4$ K, outgoing flows are not excited, and all of the baryons concentrate to the halo center regardless of the UV strength. The total mass slightly increases from the initial state by accretion of the diffuse gas in the outer part.
4.2. Mass Loss

The photoheated gas flows outward from the surface of the self-shielded region, while the central part continues contracting. The rate of the gas mass loss can be calculated as

$$M_{\text{ph}} = \int_S dS \cdot (\rho_0 v^2), \quad (22)$$

where $S$ is the surface area of the launching base. Note that the right-hand side of this equation is equal to that of Equation (15). The mass-loss rate is essentially determined by the radial extent of the self-shielded region, because the initial velocity of the photoevaporative winds is typically the sound speed of the ionized gas ($\sim 10 \text{ km s}^{-1}$), and the base density is determined by the EUV flux. The mass flux at the base is not strongly dependent of the gas metallicity (Bertoldi 1989; Nakatani & Yoshida 2019).

We measure the total gas mass within a halo as

$$M_b = \int_{r_r \leq r_m} 2\pi r \rho_b \, dr \, dR. \quad (23)$$

The evolution of $M_b$ can be characterized by two phases separated by the time when the diffuse outer part is stripped off and a “naked” dense core is left. In the later phase, the core is directly exposed to UV radiation but the net mass loss is small owing to small geometrical size. The photoevaporation rate decreases rapidly during the transition phase. A similar process is known also in the study of molecular cloud photoevaporation (e.g., Bertoldi 1989; Nakatani & Yoshida 2019). It is difficult to follow the photoevaporation process in detail after the transition phase, because the small core is resolved only with several computational cells in our simulations. We thus calculate $M_b$ only up to the transitional phase.

We empirically determine the transition time by the following conditions:

$$\frac{1}{M_b} \int_{V_{\text{sub}}} 2\pi r \rho_b \, dr \, dR > 0.8 \quad (24)$$

and

$$\rho_{b,\text{max}} \gg \rho_{b,0} \quad (25)$$

Here, $\rho_{b,\text{max}}$ is the maximum density in the computational domain, $V_{\text{sub}}$ is the region where $r \leq r_{\text{vir}}$, $\rho_b > 10^{-3} \rho_{b,\text{max}}$, and $\rho_{b,0} \equiv \rho_b(d = 0, r = 0)$. We adopt the condition $\rho_b > 10^{-3} \rho_{b,\text{max}}$ for $V_{\text{sub}}$ since this dense region initially accounts for about a half of the initial baryonic mass. The left-hand side of Equation (24) (hereafter condensation parameter) is initially $\approx 0.53$ and is independent of $M$ or $z$. It characterizes to what extent the densest region of halos accounts for the total baryonic mass $M_b$. The condensation parameter decreases while the central density increases faster than the halo evaporates. It approaches unity as the diffuse gas continues evaporation leaving the dense core at the center. For the halos that evaporates without leaving the cores, the condensation parameter decreases at the early evolutionary phase when the D-type I-front sweeps the gas; it remains nearly constant afterwards until $\rho_{b,\text{max}}$ drops below $\rho_{b,0}$ (see Equation (25)). For the halos that leaves the dense cores, the parameter quickly decreases to $\lesssim 0.3$ either by the shock or the cooling-triggered condensation. Then, it increases toward unity as the diffuse gas evaporates to reduce $M_b$.

Figure 4 shows the evolution of $M_b$ for halos with $M = 10^{2.5-8} M_\odot$ with various metallicities and other parameters.

For given values of $M$, $J_21$, $z_{\text{IN}}$, $M_b$ evolves along the same track on the $t$-$M_b$ plane for various metallicities. We have discussed in Section 4.1 that the effect of metal enrichment appears clearly in the concentrating core, but the strength of EUV-driven photo-evaporation is independent of metallicity. Thus the outer diffuse gas has not cooled in the early evolutionary phase even with nonzero $Z$ cases, and the size of the self-shielded region is similar regardless of $Z$ (compare the rightmost panels in Figure 3). Therefore, the evolution of $M_b$ does not differ significantly until the diffuse envelope gas is lost.

In the low-mass halos ($T_{\text{vir}} \lesssim 10^4 \text{ K}$), $H_2$ cooling is essential to form a dense core at any metallicity. The hydrogen molecules are produced via the grain-catalyzed reaction. Since the $H_2$ self-shielding significantly reduces the flux of LW photons reaching the central region, the photodissociation time and the $H_2$ formation time are longer than the dynamical time of the gas. Then the production rate of $H_2$ is given by $\approx \gamma_{H_2} n_H (Z/Z_\odot)$. The $H_2$ abundance increases to

$$\gamma_{H_2} \sim \gamma_{H_2} (Z/Z_\odot) \frac{r_{\text{vir}}}{c_s}$$

$$\approx 5.5 \times 10^{-3} \left( \frac{n_H}{200 \text{ cm}^{-3}} \right) \left( \frac{Z}{10^{-3} Z_\odot} \right) \left( \frac{p}{10} \right)$$

$$\times \left( \frac{M}{10^6 M_\odot} \right)^{1/3} \left( \frac{1+z}{10} \right)^{-1}$$

by the time the shock reaches the central region. The $H_2$ cooling rate is approximated as $\Lambda_{H_2} \approx 6 \times 10^{-25} \rho_b (T/200 \text{ K})^{3/2} (n_H/200 \text{ cm}^{-3})$ erg s$^{-1}$ for $200 \text{ K} \lesssim T \lesssim 10^4 \text{ K}$, and the corresponding cooling time is

$$t_{\text{cool}} \approx 1.1 \left( \frac{Z}{10^{-3} Z_\odot} \right)^{-1} \left( \frac{1+z}{10} \right)^{-1.3}$$

$$\times \left( \frac{M}{10^6 M_\odot} \right)^{-5.6/3} \left( \frac{n_H}{200 \text{ cm}^{-3}} \right)^2 \text{ Myr.} \quad (26)$$

The nominal $n_H$ is the initial central density of the halo at $z = 10$. Defining an approximate evaporation time as

$$t_{\text{eva}} \sim \Omega_b \Omega_m^{-1} M/(\pi r_{\text{vir}}^2 \rho_0)$$

$$\approx 2.0 J_{21} \left( \frac{M}{10^6 M_\odot} \right)^{1/3} \left( \frac{1+z}{10} \right)^2 \text{ Myr},$$

the condition, $t_{\text{cool}} < t_{\text{eva}}$, reduces to

$$M > 0.8 \times 10^6 J_{21}^{1/3} \left( \frac{Z}{10^{-3} Z_\odot} \right)^{-1/2.2}$$

$$\times \left( \frac{n_H}{200 \text{ cm}^{-3}} \right)^{-1/1.1} \left( \frac{1+z}{10} \right)^{-3/2} M_\odot.$$

Halos satisfying this condition condense faster than losing the gas by photoevaporation. In Figure 4, such halos have the endpoints indicated by the round markers. Because primordial halos contain no metals and dust grains, $H_2$ cooling is inefficient, and thus the endpoints indicated by round markers do not appear for low-mass primordial halos (see yellow lines).

We show the time evolution of $M_b$ for halos with various $M$, including the high mass case ($T_{\text{vir}} > 10^4 \text{ K}$; $M \geq 10^{7.5} M_\odot$) in Figure 4 (Panel b). A large amount of gas is lost from low-mass
Figure 4. Time evolution of the total gas mass relative to the initial gas mass for selected runs. The panels show the dependence of the gas mass evolution on the four simulation parameters \((Z, M, J_{21}, z_{IN})\): (a) metallicity dependence of the gas mass evolution for \(M = 10^{5.5}, 10^6\) \(M_\odot\) with \(J_{21} = 10\) and \(z_{IN} = 10\). The line colors and styles differentiate metallicity and the halo mass, respectively. The round marker points indicate the time at which all of the atmospheric gas has been lost, leaving an concentrated core (see Equation (24) and Equation (25)). Note that the curve for \(Z10^{3,5}M_\odot/10\) overlaps those of \(Z = 10^{3,5}M_\odot/10\) and \(Z=10^{3}M_\odot/10\). (b) Halo mass dependence of the gas mass evolution for \(Z = 0, 10^{-3}Z_\odot\) with \(J_{21} = 1\) and \(z_{IN} = 10\). The line colors correspond to halo masses as annotated. The lines for \(Z = 10^{3}M_\odot/10\) and \(Z=10^{3}M_\odot/10\) overlap. (c) \(J_{21}\) dependence of the gas mass evolution for \(Z = 0, 10^{-2}Z_\odot\) with \(M = 10^6\) \(M_\odot\) and \(z_{IN} = 10\). Solid, dashed, and dotted lines indicate \(J_{21} = 1, 0.1\), and \(0.01\), respectively. (d) \(z_{IN}\) dependence of the gas mass evolution for \(Z = 0, 10^{-3}Z_\odot\) with \(M = 10^{5.5}\) \(M_\odot\) and \(J_{21} = 1\).

halos \((T_{vir} < 10^4\) K\) but the mass-loss rates are significantly smaller for massive halos with \(10^5\) \(M_\odot\). Interestingly, the gas mass increases slightly. The diffuse gas at \(r > r_{vir}\) is accreted while the central part keeps cooling. Photoevaporation is hardly observed in these runs, because the initial temperature, \(T_{vir}\), is higher than the typical temperature of a photoheated gas in the first place. Halo’s gravity is so strong that it retains the photoheated gas. Another important feature is that the gas mass evolution of the massive halos is nearly independent of metallicity. Radiative cooling by hydrogen is dominant in these halos (Section 4.1).

4.3. Radiation Intensity

In photoionized regions, the characteristic ionization time, \(t_{ion} \sim 0.01J_{21}^{-1}\) Myr, is orders of magnitude shorter than the typical crossing time of photoevaporative flows, \(t_{vir} \simeq r_{vir}/10\) km s\(^{-1}\) \(\simeq 10\) \(M_\odot\)^\(1/2\)((1 + z)/10)^\(-1\) Myr with \(M_{halo} \equiv M/10^3\) \(M_\odot\). For weak UV radiation, the I-front is located at the outer part of the halo where the density is low. We find that the boundary where \(n_{HII} = 0.5\) is located at a radius where \(N_{HII} \sim 10^9J_{21}^{1/2}\) cm\(^{-2}\), and the base density is approximately estimated to be \(n_{HII} \sim 10^{-1}, 10^{-0.5}J_{21}\) cm\(^{-3}\), which is consistent with the result of Shapiro et al. (2004).

The UV intensity \(J_{21}\) does not strongly change the qualitative dynamical photoevaporation process. However, the mass-loss rate sensitively depends on \(J_{21}\), because the base is located at a larger \(r\) for lower \(J_{21}\), where the gas density is lower. One may naively expect that the geometrical size can make the photoevaporation rate large, but the small density at the large radius mitigates increase of mass loss (see Equation (22)). The base density is approximately proportional to the UV flux, while the base radius increases by only a small factor because the density rapidly decreases with increasing radial distance. Hence, the mass-loss rate, \(\dot{M}_{halo}\), decreases for smaller \(J_{21}\) as can be seen in Figure 4(c). The mass-loss rate decreases with time, \(d\dot{M}_{halo}/dt < 0\), because the geometrical cross section of halo decreases. We also note that the
characteristic mass-loss time for massive halos is longer than the Hubble time at the epochs considered.

4.4. Turn-on Redshift

Halos forming at different redshifts have different properties. Most notably high-redshift halos are more compact and denser (see Equation (11)). We study cases with different \( z_{\text{IN}} \), the timing of radiation turn-on, when the cosmological I-front reaches the halo. One can consider that different \( z_{\text{IN}} \) effectively corresponds to different reionization histories, or to an inhomogeneous reionization model in which the effective \( z_{\text{IN}} \) differs from place to place.

The process of photoevaporation is essentially the same as those described in Section 4.2 and Section 4.3. The gas density at the photoevaporative flow “base” is primarily set by \( J_{21} \), and is not explicitly dependent of \( z_{\text{IN}} \). However, the relative distance of the base to the halo center, \( \xi(\equiv r/r_{\text{vir}}) \), is larger for halos at higher redshift owing to higher average density, while the size of the neutral, self-shielded region is smaller. These two effects nearly cancel out and yield photoevaporation rates nearly independent of \( z_{\text{IN}} \). We find that \( |\dot{M}_{b}| \) increases only by 20%–30% with \( \Delta z_{\text{IN}} = -5 \).

In the characteristic minihalo case with \( M = 10^{5.5} M_\odot \), \( J_{21} = 1, Z = 10^{-3} Z_\odot \), and \( z_{\text{IN}} = 10, \) about 80% of the initial gas mass is lost, and the mass-loss fraction is only \( \sim 60\% \) with \( z_{\text{IN}} = 20 \). Similar trend of the final core mass is seen in other runs with different \( M \) and \( J_{21} \). This is consistent with the results of Iliev et al. (2005) who show weak dependence of mass-loss timescale on turn-on redshift.

4.5. Gas Mass Evolution

We have shown that the gas mass evolution depends most sensitively on \( M \) and \( J_{21} \). Physically, these are the most relevant quantities to the gravitational force and the mass flux, respectively (see Equation (15)). The results of our numerical simulations can be characterized by two quantities: the half-mass time, \( t_{1/2} \), at which the gas fraction decreases to 0.5, and the remaining mass fraction, \( f_{b,\text{rem}} \), which is the mass fraction of the “remnant” condensed core. Figure 5 shows the distribution of \( t_{1/2} \) and \( f_{b,\text{rem}} \). This summarizes the overall dependence of gas mass evolution on the simulation parameters \( (Z, J_{21}, M, \text{and } z_{\text{IN}}) \). Massive halos at lower \( z_{\text{IN}} \) reflect lower average densities. The mass (symbol size) is larger toward the upper right corner in each panel of Figure 5, indicating that the remaining mass fraction \( f_{b,\text{rem}} \) increases for higher halo mass. Also halos with higher metallicity have higher \( f_{b,\text{rem}} \) owing to efficient cooling (Section 4.2). For higher \( J_{21} \), both \( t_{1/2} \) and \( f_{b,\text{rem}} \) are smaller because of the faster mass loss for stronger UV radiation (Section 4.3).

4.6. Similarity in Mass Loss

We have shown that the mass-loss rate has weak metallicity dependence at least until the bulk of the diffuse halo gas is stripped off. In this section, we first derive nontrivial similarity of the gas mass evolution for metal-free halos. We then apply this model to other low-metallicity cases.

In Section 3, we have developed an analytic model with a key parameter \( \chi \) that characterizes the gas mass evolution of a photoevaporating halo. There, the effect of the host halo’s gravity has not been incorporated (Equations (15), (16), and (20)). We expect that deceleration by gravity becomes important for halos whose virial temperatures are comparable to the typical temperature of the photoionized gas, \( \sim 10^4 \text{K} \). In such cases, assuming \( c_s = 10 \text{ km s}^{-1} \) as the photoevaporative flow velocity overestimates the photoevaporation rate (see Equation (20) and the description above it). To account for the
deceleration, we adopt a “reduced” parameter defined as
\[
\chi' \equiv \frac{c - \frac{V}{c}}{c_i}
\] (27)
in the following discussion. The derived values are listed in Table 1.

The top panel of Figure 6 shows the gas mass evolution of metal-free halos with various parameter sets in the dimensionless form. Halos with similar \( \chi \) or \( \chi' \) evolve on essentially the same track in the \( \tilde{M} - \tilde{t} \) plane. We explain the similarity further by using a specific example as follows. The simulation parameters of \( Z \sim J0M6.5z10 \) (cyan-dotted line; \( \chi = 0.54, \chi' = 0.24 \)) are close to those of \( Z \sim J0M6.5z20 \) (yellow dashed line; \( \chi = 0.078, \chi' = 0.018 \)), but the mass evolution significantly deviates from each other on the dimensionless plane. On the other hand, it is closer to those of \( Z \sim J1M4.5z20 \) (cyan-green solid line; \( \chi = 0.25, \chi' = 0.21 \)) and \( Z \sim J2M5z10 \) (cyan dashed line; \( \chi = 0.31, \chi' = 0.25 \)). A more straightforward case is \( Z \sim J2M5z10 \) (light blue dashed line; \( \chi = 0.54, \chi' = 0.48 \)), as expected from the close parameter values. Interestingly, the \( Z \sim J2M4.5z10 \) (blue dashed line; \( \chi = 0.54, \chi' = 0.48 \)) run is also very close to \( Z \sim J1M5.5z10 \) (cyan dashed-dotted line; \( \chi = 0.54, \chi' = 0.4 \)).

We find that gravitational deceleration of the photoevaporating gas is an important factor. The evolution in Run \( Z \sim J0M6.5z10 \) (cyan dotted line) is close to other cases with green lines, but its \( \chi \) value \((\approx 0.54)\) is actually closer to those of runs indicated by the blue lines. Also, \( \chi \approx 0.31 \) for Run \( Z \sim J0M7z10 \) (yellow-green solid line) is close to those of runs indicated by the green lines, but the actual evolution apparently deviates. Gravitational deceleration of the photoevaporating gas reduces the mass-loss rate for these relatively massive minihalos with virial temperature of 2400–5000 K. Clearly, it is important to incorporate the correction of \( \chi \) owing to deceleration. We conclude that \( \chi' \) is the essential parameter to characterize the gas mass evolution of photoevaporating halos.

The bottom panel of Figure 6 shows a correlation between \( \chi' \) and the dimensionless half-mass time \( \tilde{t}_{1/2} \equiv t_{1/2}/t_0 \) for metal-free halos. There is a tight correlation given by \( \tilde{t}_{1/2} \approx 0.2 \chi^{r-0.75} \). The correlation confirms the importance of \( \chi' \) in characterizing the gas mass evolution of photoevaporating halos. Note that the same correlation holds for low-metallicity halos. Thus, the fit can be applied for halos with any metallicity that lose more than a half of the initial mass.

4.7. Fitting Function

Based on the similarity studied so far, we derive a fit of \( \tilde{M} \) that can be readily used in seminumerical models (e.g., Fialkov et al. 2012, 2013, 2015; Sobacchi & Mesinger 2013a; Cohen et al. 2016). From the result shown in Figure 4, we propose the function
\[
\tilde{M}_{\text{fit}} = f(\tilde{t}) = \frac{1 - C_1}{(\tilde{t}/\tilde{t}_0)^p + 1} + C_i, \tag{28}
\]
where \( p, C_1, \tilde{t}_0 \) are fitting parameters, which control the steepness of mass decrease, the remaining mass fraction, and the dimensionless time at which \( \tilde{t}_0 = (1 + C_i)/2 \), respectively. We restrict the parameter ranges to \( 0 \leq C_i \leq 1, 0 \leq \tilde{t}_0, \) and \( 0 \leq p \), in order to avoid unphysical fitting results.

We list the best fit values in Table 1. The excellent accuracy of the fit can be seen in Figure 7 in comparison with the simulation results. The three-parameter function fits the simulation results well until the mass decreases to \( \tilde{M} \approx 0.01 \). For halos with \( T_{\text{vir}} > 10^4 \) K, the gas mass increases owing to the halo’s strong gravity. We do not consider the slight increase.

![Figure 6](image-url)
of gas mass in deriving the fit of Equation (28), and simply set $M_{\text{fit}} = 1$.

Since we have not followed the evolution of the dense core after the outer diffuse gas is photoevaporated, the resulting $M_{\text{fit}}$ does not strongly depend on metallicity. The photoevaporation becomes inefficient in the late phase when the concentrated core is directly exposed to the UV radiation, because of its small geometrical cross section (see Equation (15)). In order to follow the late phase evolution more accurately, we will need to run simulations with a much higher spatial resolution so that the small core can be fully resolved. However, we expect that the mass evolution would not differ significantly from those obtained in this section, because the remaining mass fraction is already small in the late phase.

5. Discussion

5.1. Star Formation in Metal-enriched Halos

As the surrounding gas cools and falls on to the center, the central gas density further increases but the temperature remains low, and thus the dense core can be gravitationally unstable to induce star formation. In metal-enriched halos, cooling by metal atoms/ions and by dust grains enable the gas to condense even under the influence of strong UV radiation. Hence, metal enrichment can effectively lower the mass threshold of star-forming halos ($\sim 10^5$–$10^6 M_\odot$; e.g., Tegmark et al. 1997; Machacek et al. 2001; Bromm et al. 2002; Yoshida et al. 2003). In this section, we study the relation between metallicity and the minimum star-forming halo mass, $M_{\text{min}}$.

We assume that star formation occurs when an enclosed gas mass,

$$M_{\text{enc}}(r) = \int_{r_{\text{in}}}^r \rho_v \, dV,$$

exceeds the Bonnor–Ebert mass (Ebert 1955; Bonnor 1956),

$$M_{\text{BE}}(r) \approx 1.18 \frac{c_s^4}{P_c G^2},$$

at a spherical radius of $r$. Here, $c_s$ is the average sound speed of gas, and $P_c$ is the confining pressure. We calculate $c_s$ and $P_c$ by integrating the pressure within an enclosed volume $V$ and over the corresponding enclosing surface $\partial V$, respectively,

$$c_s(r) = \left( \frac{M_{\text{enc}}}{V} \int_{r_{\text{in}}}^r P \, dV \right)^{1/2},$$

$$P_c(r) = \frac{1}{4\pi r^2} \int_{\partial V} P \, dS.$$

Since we are interested in star formation within dark halos, we set the enclosing radius to the scale radius (core radius), $r_{\text{c}} = r_{\text{vir}}/\epsilon_{N}$, in Equations (29)–(32). We regard a halo as star-forming if it has $M_{\text{enc}}(r_{\text{c}})/M_{\text{BE}}(r_{\text{c}})$ larger than unity at a certain point during the evolution. Figure 8 shows star-forming halos and non-star-forming halos defined by this condition. We also provide a fit to the resulting $M_{\text{min}}$ as a function of $\epsilon_{N}$, $J_{21}$ and $Z$ in the Appendix.

In molecular cooling halos, the gas cools to satisfy the unstable condition, $M_{\text{enc}}/M_{\text{BE}} > 1$, at any metallicity, including the primordial case. The halos retain more than 10% of the initial gas. The remaining gas mass is larger for higher halo mass and for lower $J_{21}$. It is nearly unity for atomic cooling halos ($T_{\text{vir}} > 10^4$ K).

We find strong impact of metal enrichment in halos whose mass is lower than the atomic cooling limit. The gas cools via H$_2$ cooling faster than the bulk of the gas is photoevaporated. This effect is clearly seen in higher-metallicity halos (Figure 8).

Interestingly, the minimum collapse mass is lowered even with very small metallicities ($Z \lesssim 10^{-5} Z_\odot$), and becomes as small as $M_{\text{min}} \sim 10^4 M_\odot$ with $Z \gtrsim 10^{-4} Z_\odot$.

Wise et al. (2014) showed that star formation is active in the metal-cooling halos, and the ionizing photons of the formed stars can provide up to 30% of ionizing photons responsible for reionization. Metal-cooling halos are a heavy analog of molecular cooling halos and have masses slightly below the atomic cooling limit ($\sim 10^8 M_\odot$). We have found a light analog of molecular cooling halos in which the gas cools by H$_2$.
molecules that are formed by grain-catalyzed reactions. Recent numerical simulations suggest formation of massive or even very massive stars in metal-enriched halos (Chiaki et al. 2016; Fukushima et al. 2018). These stars can allow the beginning of reionization to occur earlier than without the metal effects. It does not likely change the redshift where reionization completes (Norman et al. 2018).

5.2. Implications of the 21 cm Line Observations

The hyperfine spin-flip signal of atomic hydrogen (the so-called 21 cm line) is a promising probe of the neutral IGM in the Epoch of Reionization. Because the strength of the 21 cm signals depend crucially on astrophysical processes at the early epochs, observations of the emission/absorption against the CMB will provide invaluable information on star formation and the physical state of the IGM.

Seminumerical models have been used to predict the large-scale fluctuations of the 21 cm signal (e.g., Mesinger et al. 2011; Visbal et al. 2012; Fialkov & Barkana 2014). The results of such models are often utilized to derive upper limits on the high-redshift astrophysics (e.g., Monsalve et al. 2018, 2019; Ghara et al. 2020; Mondal et al. 2020; Greig et al. 2020), and make forecasts for ongoing measurements of the 21 cm power spectrum with experiments such as HERA (DeBoer et al. 2017), the Low-Frequency Array (e.g., Mertens et al. 2020), Murchison Widefield Array (e.g., Trott et al. 2020), Owens Valley Radio Observatory Long Wavelength Array (Eastwood et al. 2019), and the future SKA (Koopmans et al. 2015), and of the 21 cm sky-averaged (global) signal using the Large-aperture Experiment to Detect the Dark Age (Price et al. 2018), Shaped Antenna measurement of the background Radio Spectrum (Singh et al. 2018), Experiment to Detect the Global EoR Signature (Bowman et al. 2018), Probing Radio Intensity at high-Z from Marion (Philip et al. 2019), Mapper of the IGM Spin Temperature, and Radio Experiment for the Analysis of Cosmic Hydrogen.

The speed and convenience of the seminumerical methods come along with their poor spatial resolution, which is compensated by extensive use of sub-grid models. For example, Sobacchi & Mesinger (2013a) and Cohen et al. (2016) studied the effect of the UV background radiation in terms of the gas cooling threshold $M_{\text{cool}}$. Halos more massive than $M_{\text{cool}}$ are regarded as star-forming halos. The simple prescription in previous studies, however, does not take into account the effect of metal enrichment. As we have shown in Section 5.1, star formation can occur in metal-enriched minihalos even during reionization. Cohen et al. (2016) showed that star formation in such metal-enriched minihalos affects the 21 cm signal from high redshifts. In particular, signatures of baryon acoustic oscillation (BAO) imprinted on the 21 cm signal is amplified and can possibly be detected over a wide range of redshift. Cohen et al. (2016) set a threshold mass for cooling (and thus star formation) similar to that of molecular cooling halos. Interestingly, this somewhat conservative assumption is well justified by our results, where metal enrichment lowers $M_{\text{min}}$ from the molecular cooling limit by an order of magnitude for $Z \gtrsim 10^{-4} Z_\odot$ even under the effects of UVB. Our results support the predicted enhancement and survival of BAO signature imprinted on 21 cm signals due to effects of metal enrichment.

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**Figure 8.** Star formation vs. photoevaporation for all our runs. The left, middle, and right panels correspond to $z_{\text{IN}} = 10, 15,$ and 20, respectively. The horizontal and vertical axes are halo mass and $J_{\text{21}}$ in all the panels. Each of the three panels is divided into $11 \times 3$ rectangular blocks indicating combinations of $M$ and $J_{\text{21}}$. Each block is further divided into five square sections to show the dimension of metallicity; the corresponding metallicity is $Z = 0, 10^{-6}, 10^{-5}, 10^{-4},$ and $10^{-3} Z_\odot$ from top to bottom. We represent star-forming halos by filling the sections with colored squares as $10^{-3} Z_\odot$ (navy), $10^{-4} Z_\odot$ (purple), $10^{-5} Z_\odot$ (pink), $10^{-6} Z_\odot$ (orange), and $0 Z_\odot$ (yellow). The vertical blue dashed line shows the molecular cooling limit, which we derive using an expression in Fialkov et al. (2013) (see Machacek et al. 2001; Fialkov et al. 2012), for a reference; note that we have not taken into account the correction to the molecular cooling limit due to relative velocity between baryons and cold dark matter. Incorporating the correction does not significantly change the limit mass for the redshifts of interest here.

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7 http://www.physics.mcgill.ca/mist/
8 https://www.kicc.cam.ac.uk/projects/reach
5.3. Internal Stellar Feedback Effect

Massive stars formed in metal-enriched halos affect the host halo by UV radiation and stellar winds, and by supernova explosions. Then the halo gas would not only photoevaporate by external UV radiation but also can be dispersed by these internal processes.

The stellar feedback is effective if (i) the enclosed gas mass $M_{\text{enc}}$ (Equation (29)) exceeds the Bonnor–Ebert mass $M_{\text{BE}}$ (Equation (30)); (ii) stellar feedback energy deposited by massive stars, $E_{\text{dep}}$, is larger than the gravitational binding energy of the neutral gas in the self-shielded regions

$$ E_{\text{dep}} \gtrsim \frac{G(M + M_*)}{r_n} M_*, $$

where $r_n$ is the size of the neutral gas clump. We first consider only the supernova explosion energy for simplicity. The deposited energy by other feedback processes can be easily accounted for by increasing $E_{\text{dep}}$ by a suitable factor.

The cold gas that satisfies condition (i) is assumed to form stars with a star formation efficiency of $c_{\text{SN}}$ and with an initial mass function, $\Psi(M_*)$. Let $\varepsilon_{\text{SN}}$ be the average supernova explosion energy. The deposited feedback energy is estimated as

$$ E_{\text{dep}} = c_{\text{SN}} M_{\text{cold}} \varepsilon_{\text{SN}} \left( \int M_\Psi \psi M_\Psi \right)^{-1} \int M_\Psi \psi M_\Psi, $$

where $M_{\text{cold}}$ is the mass of enclosed cold gas, and $M_{\text{th}}$ is a threshold stellar mass above which stars cause supernova explosion. We adopt $c_{\text{SN}} = 0.1$, $\varepsilon_{\text{SN}} = 10^{51}$ erg, $M_{\text{th}} = 8 M_\odot$, and use the initial mass function of Chabrier (2003).

When conditions (i) and (ii) are met and stars are formed over a freefall time, the self-shielded region is expected to evolve differently from what has been shown in Section 4. If the feedback effects are strong enough to disrupt the entire halo, the halo evolution described in Section 4 is valid only up to one freefall time. Figure 9 shows the same plots as Figure 4(a), but the lines extend to the time when the stellar feedback is assumed to be effective. Massive, metal-rich halos are likely to be self-destructed before the gas is lost via photoevaporation. Note that the fraction of the cold gas that is converted to stars is small.

In metal-free halos, the stellar feedback can also be important, especially for massive halos. We expect that the stellar feedback is unimportant in low-mass halos with $T_{\text{vir}} \lesssim 100$ K ($M \lesssim 10^{13} M_\odot$ at $z = 10$) for any metallicity. In these halos, most of the gas is quickly lost by photoevaporation before star formation. In conclusion, the deposited energy due to the stellar feedback can disperse the gas from the host halo, similarly to the well-known feedback effect in dwarf galaxies devoid of H I content (Grcevich & Putman 2009; Spekkens et al. 2014).

5.4. X-Ray Effects

X-rays are attenuated by a larger column ($\sim 10^{21} \text{cm}^{-2}$) compared to EUV (Wilms et al. 2000). They can reach halos and pre-ionize/pre-heat the gas before the ionization front hits the halos. A larger attenuation column also implies larger penetration depths. Higher-density photoevaporative flows would be driven if the gas temperature increases sufficiently to allow it escape from gravitational binding. Accordingly, mass-loss rates would be significantly larger than those of EUV-driven photoevaporation.

X-ray possibly delays concentration of the self-shielded regions and thereby star formation, if it efficiently heats the gas. On the other hand, X-ray ionization can promote H$_2$ formation via the electron-catalyzed reactions: H + e$^-$ $\rightarrow$ H$^-$ + $\gamma$ and H$^-$ + H $\rightarrow$ H$_2$ + e$^-$ (Haiman et al. 1996; Bromm et al. 1999; Glover & Brand 2003; Hummel et al. 2015; Inayoshi & Omukai 2011; Inayoshi & Tanaka 2015; Glover 2016; Regan et al. 2016b). If X-rays are strong, very strong LW intensities are required to photodissociate H$_2$ in the entire halos. We expect X-rays to have significant effects on evolution of irradiated halos and star formation activities. X-ray chemistry is already implemented in our code (Nakatani et al. 2018b) and we plan to investigate their influence on halo photoevaporation in the future.

5.5. Spectral Shape

We have adopted a fixed UV spectral index of $\alpha = 1$ (Equation (13)) in our fiducial models. It is worth examining different types of radiation sources, i.e., different spectral shapes. The relative rates of H$_2$ photodissociation inside a halo and of ionization in the outer part vary depending on the exact source spectrum. In order to study the effect on the photoevaporation, we select two specific models, $Z \propto 0.5z10$ and Z4J0M5.5z10. The results with the fiducial spectrum are shown in Figure 3. We perform additional simulations with three other UV spectral energy distributions: (i) $\alpha = 1$ but the numbers of the LW photons is increased by about a factor of 4 so as to be equal to the number of ionizing photons, (ii) $\alpha = 3.55$, which yields almost the same number of photons in the LW and EUV bands, and (iii) $\alpha = 5$, with which the number of ionizing photons is 60% of the LW photons. Hereafter, we refer to the three spectral models as “eq-$\alpha1$,” “eq-$\alpha3.55$,” and “$\alpha5.0$.”
Figure 10. Snapshots of ZJ0M5.5z10 (top panels) and Z \( \propto \) J0M5.5z10 (bottom panels) with the spectral models (i), (ii), and (iii), from left to right, specified in Section 5.5. The snapshots are taken at \( t = 43 \) Myr and are shown as in Figure 3. The gray dotted-dashed lines in the temperature map indicate the ionization degree of 0.1, 0.5, and 0.9, and the black dashed lines of the density map are the H\(_2\) abundance of 0.5 \( \times \) 10\(^{-5}\) and 0.5 \( \times \) 10\(^{-4}\). Note that the blue lines appear only in the central regions of the upper panels.

Figure 10 shows the snapshots at \( t = 43 \) Myr and compares the density and temperature structures for the three spectral models. The H\(_2\) abundance exceeds \( \sim 10^{-4}\) in the central regions of the metal-enriched halo (see blue dashed lines in the density maps) for all the spectral models. Molecular cooling is efficient there, and the gas collapses to form a dense core. The LW flux in the runs with eq-\( \alpha \)-3.55 and \( \alpha \)-5.0 is higher than in the fiducial model only by about 50%. The resulting H\(_2\) mass differs only by a small factor, and thus the core evolves in a very similar manner to the fiducial case. In the run with eq-\( \alpha \)-1, the LW flux is larger by a factor of 4, and the H\(_2\) mass increases slowly. Typically the H\(_2\) mass is an order of magnitude smaller than that of the fiducial model through the evolution.

The ionizing radiation drives photoevaporation for all the spectral models. Since the average mean-free-path of the photons decreases for a steeper UV spectral energy distribution (SED), the eq-\( \alpha \)-3.55 and \( \alpha \)-5.0 models yield a sharper ionization front than the other models. Also a steeper SED yields a lower ionizing photon number flux for a given \( J_{21} \). The ionizing photon number flux of the eq-\( \alpha \)-3.55 and \( \alpha \)-5.0 models is \( \approx 20\%\) and 30\% of that of the fiducial model, respectively. The smaller EUV flux results in a geometrically more extended I-front (Figure 10) and a smaller mass-loss rate. In the runs with eq-\( \alpha \)-3.55 and \( \alpha \)-5.0, the half-mass time \( t_{1/2} \) is longer by a factor of 2 and 2.5 than that of the fiducial case, respectively. We note that \( t_{1/2} \) of the eq-\( \alpha \)-3.55 is almost the same as that of Z \( \propto \) J1M5.5z10. Consequently, for a given ionizing photon number flux, a shallower UV SED causes more effective heating in the deep, dense interior of the halo, and thus the gas dispersal is enhanced.

In summary, at a fixed \( J_{21} \), the effects of steepening the UV spectral index are to prolong the halo lifetimes and to reduce the H\(_2\) abundance inside metal-enriched halos. The UV spectral index is an important parameter that determines the dynamical evolution of minihalos via photoevaporation. However, even with the very different LW fluxes explored here, the central portion condenses to form a dense core in metal-enriched halos.

5.6. Limitations of the Model

We have studied the gas mass evolution for a wide range of model parameters. Exploring the large volume of the parameter space is essential in order to understand the evolution of a population of photoevaporating halos during reionization. Since we have adopted a few simplifications to save computational times, it would be worth examining the limitations of our results.

Halo’s gravitational potential.—We fix the dark matter halo potential throughout our simulations. In practice, halos grow in mass by mergers and accretion, which deepens the halo’s gravitational potential (e.g., Sobacchi & Mesinger 2013b). This effect can be important for a halo that has a longer mass-loss time than its growth time, namely, when the halo merger timescale, \( t_{\text{merg}} \sim \left( GM/(4\pi r_{\text{vir}}^3/3)\right)^{-0.5} \), is longer than the mass-halving time \( t_{1/2} \) for \( \chi^2 \gtrsim 0.05 \) (Section 4.6). The parameter range satisfying this condition is represented by the shaded regions in Figure 11. A halo can grow faster than being photoevaporated for \( M \gtrsim 10^{6.5} M_\odot \) with \( J_{21} = 0.01-1 \) at \( z = 10-20 \). When such a halo grows to have \( T_{\text{vir}} \gtrsim 10^4 \) K, photoevaporation is significantly suppressed so that it can sustain the bulk of the baryonic mass despite photoevaporation. On the other hand, lower-mass halos likely lose the bulk of the
A small amount of metals is engulfed by the intergalactic I-front. Note that even a very small metallic enrichment may have proceeded in some halos already before being hit by the minihalo. In reality, chemistry and thermal evolution of metalrich minihalos is switched on at the same time as the I-front arrival. The UV intensity is also a factor that affects the thermochemistry of the photoevaporating halos that evolve in a similar manner to those photoevaporating in their parental molecular cooling halos.

The concentrating cores of the molecular clouds are formed in the central portions of the halos. In our model, the internal chemothermal evolution is switched on at the same time as the I-front hits the minihalo. In reality, chemistry and thermal evolution must have proceeded in some halos already before being engulfed by the intergalactic I-front. Note that even a very small amount of metals is sufficient to form H$_2$ molecules to trigger collapse of the central portion. Ahn & Shapiro (2007) adopted a 1D model to investigate the influence of the “delay” for primordial halos. They showed that a delayed shock can still disrupt collapsing cores in halos with $M \sim 10^8 M_\odot$. It appears that the exact timing of I-front arrival is another important parameter that affects the star formation and mass loss via photoevaporation. It would be interesting to quantify the disruption/photoevaporation efficiency as a function of delay time in future studies.

Metallicity—Finally, metallicity is also fixed in individual runs of our simulations. Time-dependent metallicity might affect the thermochemistry of the photoevaporating halos that survive for a long time compared to the timescale on which metallicity increases by orders of magnitude. However, we expect that the metallicity-dependent trend would not significantly differ from what we have reported in this study. We also note that metallicity effect would be apparent in the formation of the concentrating cores but not in the evaporation time.

6. Conclusions and Summary

Photoheating and metal enrichment of halos during the epoch of reionization can affect star formation efficiency, and thus, change the course of reionization. The effects of metal enrichment, however, have never been explored systematically in prior works. We have run a suite of 495 hydrodynamic simulations of photoevaporating minihalos ($T_{\text{vir}} \lesssim 10^4$ K) irradiated by a planer UV, covering a wide range of metallicity, halo mass, UV intensity, and turn-on redshift of UV sources.

Our main findings are summarized as follows:

1. In low-mass minihalos with $T_{\text{vir}} < 10^4$ K, the gas cools mainly via C II and H$_2$ line emission to a temperature of $\lesssim 100$ K if $Z > 10^{-4} Z_\odot$. H$_2$ molecules are produced through grain-catalyzed reactions in the self-shielded, neutral region. The cooled gas concentrates toward the potential center to form a dense core.
2. The evolution of the gas mass is qualitatively the same at any metallicity. The photoevaporation rate decreases after the bulk of the diffuse gas is lost. The dispersal of the diffuse gas completes at an earlier time for halos with higher metallicity.
3. In halos with $T_{\text{vir}} > 20,000$ K, the gas cools by hydrogen Lyman$\alpha$ cooling, and thus the overall evolution of photoevaporating halos does not depend on metallicity.
4. The photoevaporation rate depends only weakly on turn-on redshift, and it is slightly smaller for higher $z_{\text{turn}}$.
5. There is a simple scaling relation for the gas mass evolution of photoevaporating minihalos. The time evolution is characterized by a parameter ($\chi'$) scaling as $\propto J_1^{1/2} M^{-1/2} (1 + z)^{-3}$. It indicates that the obtained evolution applies to any halos with the same $\chi'$. We give a fit to the gas mass evolution as a function of time.
6. The concentrating cores of the molecular/atomic cooling halos are likely a suitable environment for star formation. The efficient cooling of metal-rich halos fastens the concentration, and it results in lowering the molecular cooling limit by a factor ($< 1$ dex) for small metallicities ($Z \lesssim 10^{-5} Z_\odot$) and by an order of magnitude for metal-rich cases ($Z \gtrsim 10^{-4} Z_\odot$). Stellar feedback of the formed stars may be significant enough to disperse the baryons of the parental molecular cooling halos.

Our study suggests the existence of small mass, metal-enriched halos in which stars are formed even under the influence of emerging UV background radiation.

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Appendix

Fit to the Minimum Mass

The maximum value of $M_{enc}/M_{BH}$ increases with increasing mass of the star-forming halo. The mass ratio can be well approximated as $log(M_{enc}/M_{BH})_{\text{max}} = a(logM)^2 + b logM + c$, where $a$, $b$, and $c$ are fitting parameters. We derive $M_{\text{min}}$ by finding a root of $a(logM)^2 + b logM + c = 0$ for each set of $(Z, J_{21}, z_{\text{in}})$. By further fitting the derived $M_{\text{min}}$ as a function of $(Z, J_{21}, z_{\text{in}})$, we obtain

$$log M_{\text{min}} = c_1 z_{\text{in}} + c_2 (z_{\text{in}} + 1) + 7.4,$$

$$c_1 z_{\text{in}} = -0.080(z_{\text{in}} + 1) + 7.4,$$

and $Z = Z/Z_{\odot}$.

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