Supermassive star formation in a massive cloud with $H_2$ molecules

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ABSTRACT

Recent three-dimensional cosmological simulations of protogalaxy formation have suggested that supermassive stars (SMSs) can form in warm gas clouds in which $H_2$–cooling is overwhelmed by dynamical heating (Wise et al. 2019), but they stopped short of the following growth of a central protostar. Here we examine whether accretion on the protostellar core in this cloud is sufficiently rapid, in the face of the radiation feedback, to produce a SMS. We perform one-dimensional radiation-hydrodynamical simulations of the hot collapsing cloud with non-equilibrium chemical reactions directly adopting the cloud properties from Wise et al. (2019) as an initial condition. We find that the stellar Lyman-Werner (LW) radiation dissociates $H_2$ in the inner regions of the gas flow, increasing gas temperature and thermal pressure, and temporarily stopping the accretion. However, this negative feedback ceases when the self-gravity and inward ram pressure force on larger scales push the gas inward. The central protostar is unable to create an HII region due to the high density, and grows to a mass of $\gtrsim 10^5 M_\odot$. Our results suggests the successful formation of SMSs, and resulting massive ($\sim 10^5 M_\odot$) remnant black holes, in dynamically heated haloes, but need to be confirmed in two- or three-dimensional simulations.

Key words:

1 INTRODUCTION

Observations of distant quasars over the past two decades have shown that supermassive black holes (SMBHs) with masses $\gtrsim 10^9 M_\odot$ exist at redshift $z \gtrsim 6$ (e.g., Fan et al. 2001; Mortlock et al. 2011; Wu et al. 2015; Matsuoka et al. 2016; Bañados et al. 2018; Reed et al. 2019; Onoue et al. 2019; Yang et al. 2020). The existence of these SMBHs means that a billion solar mass or more can rapidly accumulate in a small region within 1 Gyr. The physical mechanism(s) for when and how this occurs remain unknown.

One of the SMBH formation scenarios is that they grow from lower-mass BHs with $\sim 10^{-10^5} M_\odot$ exist at redshift $z \gtrsim 6$ (e.g., Fan et al. 2001; Mortlock et al. 2011; Wu et al. 2015; Matsuoka et al. 2016; Bañados et al. 2018; Reed et al. 2019; Onoue et al. 2019; Yang et al. 2020). The existence of these SMBHs means that a billion solar mass or more can rapidly accumulate in a small region within 1 Gyr. The physical mechanism(s) for when and how this occurs remain unknown.

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Several authors pointed out that the subsequent growth of seed BHs through gas accretion can be suppressed by radiation feedback (Ciotti & Ostriker 2001; Milosavljević et al. 2009; Novak et al. 2011; Park & Ricotti 2011; but see also Inayoshi et al. 2016 for the possible solution of super-Eddington BH accretion). To alleviate the slow growth caused by this suppression, one possibility is the so-called “direct collapse”, in which larger seed BHs with masses $\gtrsim 10^5 M_\odot$ are produced (Loeb & Rasio 1994; Oh & Haiman 2002; Bromm & Loeb 2003; Begelman et al. 2006). In this model, a protostellar core forms in the center of a gas cloud surrounded by a dark matter halo with a virial mass of $\gtrsim 10^7 M_\odot$. These so-called “atomic cooling” haloes are larger than the “minihaloes” ($\sim 10^{5-6} M_\odot$) where the first Pop III

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stars form. The gas in the atomic-cooling halo cools via H\textsubscript{2} lines, and the central protostar can grow via rapid gas accretion into a supermassive star (SMS) with a mass \(\gtrsim 10^5 M_\odot\). The SMS gravitationally collapses to a BH with a similar mass due to general relativistic instability (Umeda et al. 2016; Woods et al. 2017; Haemmerlé et al. 2018).

One of the conditions for SMS formation in atomic-cooling haloes is to avoid H\textsubscript{2}–cooling induced fragmentation. This can be achieved by irradiating the halo by an unusually strong external far-ultraviolet Lyman-Werner (LW) radiation emitted from nearby star-forming galaxies (Omukai 2001; Dijkstra et al. 2008; Regan et al. 2014; Sugimura et al. 2014; Inayoshi et al. 2014; Becerra et al. 2015; Latif et al. 2016; Chon et al. 2016). A sufficiently intense LW radiation fully dissociates H\textsubscript{2} molecules and suppresses H\textsubscript{2} cooling in some rare situations: the halo has a nearby neighbouring halo with highly synchronised star-formation (Visbal et al. 2014; Regan et al. 2017), so that the LW radiation flux exceeds a critical flux (Wolcott-Green & Haiman 2019), but tidal forces do not disperse the halo. The gas in the halo also needs to remain extremely metal-poor. If the H\textsubscript{2} molecules are fully dissociated, the only effective coolant is atomic hydrogen. The gas temperature can not fall below \(\sim 8000\ K\), resulting in elevated sound speeds and higher accretion rates than in cooler gas in minihaloes.

In addition to radiative processes, the dynamical effect of collapsing gas into a massive DM halo as the halo is assembled is expected to play a crucial role on the formation of SMSs, by suppressing H\textsubscript{2}–cooling (Fernandez et al. 2014) and heating the halo gas via compression and shocks (Yoshida et al. 2003), especially in haloes with unusually rapid assembly histories. Recently, Wise et al. (2019, hereafter W19) have shown, using three-dimensional cosmological hydrodynamical simulations, that such dynamical heating helps keep the gas warm, and may lead to the formation of SMSs with masses \(\gtrsim 10^5 M_\odot\) in massive haloes, even if strong external LW radiation is absent, and H\textsubscript{2} molecules are not fully dissociated. Similarly, unusually large baryonic streaming velocities can delay the collapse of gas into less massive haloes and induces violent mergers of gaseous haloes into more massive DM haloes (Hirano et al. 2017; Schauer et al. 2017; Inayoshi et al. 2018), inducing dynamical heating and helping to produce the conditions required for SMS formation.

Although dynamical heating in unusually rapidly assembling haloes is one of the promising mechanisms for the formation of SMSs, it remains unclear whether SMSs do indeed form in these haloes. In particular, radiation feedback from the growing protostar can stunt its growth, when the accretion rate is lower than a critical value of \(\approx 0.04 M_\odot\ yr^{-1}\) (Hosokawa et al. 2013; Schleicher et al. 2013; Sakurai et al. 2015). Above this critical rate, the rapid gas accumulation, as well as heat input owing to rapid accretion prevent the stellar surface from contracting via thermal emission on the Kelvin-Helmholtz (KH) timescale. The rapidly accreting protostars evolve to ‘super-giants’ which have inflated radii of \(\sim 100\ AU\). Conversely, since the KH timescale at the surface is \(10^{10–4}\ yr\), if the protostar grows at rates lower than the critical value at the beginning of its evolutionary stage the protostar cannot evolve into the super-giant protostar. The star contracts to a small radius, and develops a correspondingly high effective temperature \(\sim 10^5\ K\), emitting copious amounts of ionizing photons, which cause the radiation feedback. At their last resolved snapshots of the collapsing gas in the three-dimensional simulations by W19, the accretion rates in the innermost regions fall below this critical value (see the bottom right panel of their Figure 4), leaving the fate of the protostar unclear.

In this study, we explore the evolution of gas inflows around the growing protostar, past the epoch simulated in W19, using one-dimensional radiation hydrodynamical simulations. We adopt the initial conditions for the cloud properties directly from W19, and employ non-equilibrium chemical reactions. We include the radiation emitted by the growing protostar (as well as a tentative circumstellar disk). The main goal of this study is to assess whether the stellar radiation suppresses the accretion rate below the critical value, or if accretion remains sufficiently rapid to produce a SMS.

The rest of this paper is organised as follows. In §2, we describe the setup and details of the radiation-hydrodynamical simulations. In §3, we present the evolution of the gas clouds for cases with and without stellar radiation. In §4, we discuss our results, including their implications and some caveats. Finally, in §5, we summarise our main conclusions.

2 METHODS

2.1 Hydrodynamical simulations

In order to explore gas inflows around growing protostars, we use the hydrodynamical simulation code ZEUS (Stone & Norman 1992), including multifrequency radiation transfer, photoionization and heating, and a primordial chemical network (Inayoshi et al. 2016; Sakurai et al. 2016b). Assuming spherical symmetry, we use ZEUS to solve the continuity equation and the equation of motion in one dimension,

\[ \frac{\partial \rho}{\partial t} + \frac{1}{r^2} \frac{\partial}{\partial r} (r^2 \rho v) = 0, \]

\[ \rho \left( \frac{\partial v}{\partial t} + v \frac{\partial v}{\partial r} \right) = -\frac{\partial p}{\partial r} - \frac{\partial \Phi}{\partial r} + f_{\text{rad}}, \]

where \(t\) is time, \(r\) is the radial coordinate, \(\rho\) is the gas density, \(v\) is the velocity (defined to be negative for inflows), \(p = (\gamma - 1) \rho c^2\) is the gas pressure, \(\gamma = 5/3\) is the adiabatic index, \(e\) is the specific internal energy of the gas, \(\Phi\) is the gravitational potential (including contributions from both the growing protostar as a point mass, and from the self-gravity of the gas) and \(f_{\text{rad}}\) is the radiation pressure force. The specific internal energy \(e\) is determined by the energy equation

\[ \rho \left( \frac{\partial e}{\partial t} + v \frac{\partial e}{\partial r} \right) = -\frac{1}{r^2} \frac{\partial}{\partial r} (r^2 v) - \Lambda + \Gamma, \]

where \(\Lambda\) and \(\Gamma\) are the cooling and heating rates. \(\Lambda\) includes line cooling of H, H\textsubscript{2}, H\textsuperscript{+} and He, recombination cooling of H\textsuperscript{+} and He\textsuperscript{+}, free-free emission, collisional ionization cooling of H and He and H\textsubscript{2} dissociation cooling (Glover & Jappsen 2007; Glover & Abel 2008),

\[ \Lambda = \Lambda_\text{H} + \Lambda_\text{H}_2 + \Lambda_{\text{H}_2^+} + \Lambda_{\text{He}} + \Lambda_{\text{He}_+} + \Lambda_{\text{He}^+} + \Lambda_{\text{He}_{col}} + \Lambda_{\text{H}_{col}} + \Lambda_{\text{H}_2\text{dia}}. \]
We omit the He cooling rate by the $2^1S$ metastable excitation state which is proportional to $n_2^2 n_{He^+}$ since in our density regime this cooling rate can be invalid (see equation 14.19-20 in Draine 2011). For the heating rate $\Gamma$, we include H$_2$ formation heating, photoionization heating of H, He, He$^+$ and H$_2^-$ photo-detachment heating, and H$_2$ photodissociation heating (Abel et al. 1997; Omukai 2000).

Our non-equilibrium chemistry incorporates the nine species H, H$^+$, He, He$^+$, H$^{2+}$, e$^-$, H$_2$, H$_2^+$ and H$^-$. The chemical reactions are taken mainly from Nos. 1-32 in Table 1 of Glover & Abel (2008). We adopt the case B recombination rates for H$^+$, He$^+$ and H$^{2+}$ because diffusive photons produced by direct recombination to the ground states are immediately absorbed by the surrounding gas. We also consider photoionization, H$^-$ photodetachment and H$_2$ photodissociation, with the rates adopted from references listed in Table 1. The evolution of the number density of each species $i$ is governed by

$$\frac{\partial n_i}{\partial t} = C_i - D_i n_i,$$

where $C_i$ and $D_i$ are creation and destruction terms of species $i$ respectively. The equation is solved using a semi-implicit method updating each species in order (Anninos et al. 1997; Whalen & Norman 2006). The order of the updates is H, H$^+$, He, He$^+$, H$^{2+}$, H, H$_2$, H$_2^+$ and e$^-$. We set a calculation timestep as the smallest among the Courant crossing time (\cite{Draine2011}), the order of the up-}

$$\gamma \rightarrow H_2^+ + e^-$$

$$H^+ + \gamma \rightarrow H + e^-$$

$$H^+_2 + \gamma \rightarrow H + H^+$$

$$H^+_2 + \gamma \rightarrow 2H^+ + e^-$$

$$H_2 + \gamma \rightarrow H_2^+ + e^-$$

The radiation transfer equation reduces to

$$F_{k,\nu} = \left( \frac{r_k - 1}{r_k} \right)^2 F_{k-1,\nu} \exp \left[ -(r_k - r_{k-1}) \sum \sigma_{\nu,i} n_i \right],$$

where the subscript $k$ marks the radial cell and $\sigma_{\nu,i}$ is the absorption cross section for species $i$ (Whalen & Norman 2006). We do not consider production of diffuse EUV and FUV photons owing to recombination and radiative de-excitation in Eq. (8). Instead of considering photoionization by diffuse EUV photons, we here adopt the on-the-spot approximation where the case B recombination rate coefficient is used. The flux of diffuse FUV radiation is negligible compared to those from the central protostar and an external LW background flux (see below). We adopt the cross sections from the references shown in Table 1, except for LW radiation.

For LW radiation, we replace the exponential factor in Eq. (8) by a shielding factor that includes both self-shielding of H$_2$ and shielding of H$_2$ by neutral H,

$$f_{sh,H_2} = \frac{0.965}{(1 + \nu/b_2)^{0.8}} + \frac{0.035}{(1 + \nu)^6} \exp \left( -8.5 \times 10^{-4} (1 + \nu) 0.5 \right)$$

$$f_{sh,H^+} = (1 + x_{HI})^{-1.6} \exp(-0.15 x_{HI})$$

where $f_{sh,H_2} = f_{sh,H_2} f_{sh,H^+} f_{sh,HI}, x \equiv N_{HI}/(5 \times 10^{14} \text{ cm}^{-2}), b_2 \equiv \sqrt{\nu/(\nu_{th}^H)}/(10^6 \text{ cm s}^{-1})$ and $x_{HI} \equiv N_{HI}/(2.85 \times 10^{23} \text{ cm}^{-2})$ (Walcolt-Green et al. 2011). To obtain local estimates of the H$_2$ and HI column densities, we adopt the 'Sobolev' length defined as $L_{ Sob }' \equiv \rho / |d\rho / dr|$ and use the relations $N_{HI} = n_{HI} L_{ Sob }'$ and $N_{HI} = n_{HI} L_{ Sob }'$. We include the LW radiation from the central protostar, as well as an external LW background with specific intensity $J_{LM} = 3 \times 10^{-21} \text{ erg cm}^{-2} \text{s}^{-1} \text{ Hz}^{-1} \text{ sr}^{-1}$ (the same value as in W19). Since the background radiation irradiates the cloud from the outside in, whereas the stellar radiation irradiates the cloud from the center, we define the self-shielding factor outside-in for the background radiation, and inside out for the stellar radiation. Specifically, the self-shielding factor for the background radiation is computed as $f_{sh,k-1} = \min(f_{sh,H_2,k-1} f_{sh,H^+,k-1} f_{sh,HI,k-1}, f_{sh,k})$ with $f_{sh,k,max} = f_{sh,H_2} f_{sh,H^+} f_{sh,HI}$. The reaction rates $k_i$ and photo-heating rates $\Gamma_i$ for photoionization, H$^-$ photo-detachment and H$_2$ photodissociation (Table 1) are computed using a photon-conserving scheme (Whalen & Norman 2006) as

$$k_i = \int_{\nu_{th,i}}^{\nu_{th,i}'} \frac{4\pi}{h\nu} \sigma_{\nu,i} E_{heat,i} d\nu,$$

$$\Gamma_i = n_i \int_{\nu_{th,i}}^{\nu_{th,i}'} \frac{4\pi}{h\nu} \sigma_{\nu,i} E_{heat,i} d\nu,$$

where $\nu_{th,i}$ is the threshold frequency for species $i$, $\tilde{\nu}_i$ is the mean intensity (over solid angles) and $E_{heat,i} = h(\nu - \nu_{th,i})$. The photon conservation method here means that all the excess energy is thermalized and deposited into the gas. For the two-step H$_2$ dis-
sociation by LW photons (often referred to as the "Solomon process"; Field et al. 1966), we adopt the reaction rate

\[ k_{\text{LW}} = 1.1 \times 10^8 \frac{F_{\nu, \text{LW}}}{\text{erg s}^{-1} \text{cm}^{-2} \text{Hz}^{-1}} \text{s}^{-1} \]  

(14)

and the heating rate

\[ \Gamma_{\text{LW}} = 6.4 \times 10^{-13} n_{\text{H}_2} k_{\text{LW}} \text{erg s}^{-1} \text{cm}^{-3}. \]  

(15)

(Abel et al. 1997). The radiation pressure force is

\[ f_{\text{rad}} = \frac{n_e}{c} \int \sigma_{\text{rs}} F_\nu d\nu + \frac{\Gamma}{c}, \]

(16)

where \( \sigma_{\text{rs}} \) is the Thomson cross section and \( \Gamma \) is the total bound-free photoheating rate.

As central radiation sources, we consider contributions from both the growing protostar and from a hypothetical circumstellar disk. We include the disk component, because the gas contracting in the cores of the haloes simulated in W19 have non-negligible angular momenta (see their Extended Data Figure 4), suggesting that a circumstellar disk may form. Both sources are unresolved and located at the origin \( r = 0 \).

The stellar radiation flux at the innermost cell is

\[ F_{\nu, \text{star}} = \pi \left( \frac{R_\star}{r_{\text{min}}} \right)^2 B_\nu(T_\text{eff}), \]

(17)

where \( R_\star \) is the stellar radius, \( T_\text{eff} \) is the effective temperature, \( r_{\text{min}} \) is the radius of the innermost cell, and \( B_\nu \) is the Planck function. The stellar radius and effective temperature are calculated from a stellar evolution model described in §4.3.1 below.

\[ r_{\text{min}} = 3.14 \times 10^{15} \text{Hz} \]

\[ \times \left( \frac{M_\star}{1 M_\odot} \right)^{1/4} \left( \frac{\dot{M}}{10^{-2} M_\odot \text{yr}^{-1}} \right)^{1/4} \left( \frac{R_\star}{1 R_\odot} \right)^{-3/4}. \]

Note that the cutoff frequency \( \nu_{\text{cut}} \equiv \nu_\star \), which corresponds to the frequency of the maximum flux of the optically thick disk, always remains below the maximum frequency \( \nu_{\text{max}} \approx 2.85 \times 10^{16} \text{Hz} \) in our simulations. Since \( h\nu_\star \) is always below \( \sim 5 \text{eV} \), less than the EUV and FUV energies, the disk radiation has only a relatively minor effect on the dynamics of the flow. The total flux is \( F_{\nu, \text{disk}} = F_{\nu, \text{star}} + F_{\nu, \text{disk}} \).

We adopt a spherically symmetric, logarithmically spaced grid in the simulations. The \( k^{th} \) cell of the grid is located at \( r_k = r_{\text{min}} + (r_{\text{max}} - r_{\text{min}})(\epsilon^{k-1} - 1)/(\epsilon^N - 1) \), where \( N \) is the number of cells, \( r_{\text{max}} \) is the radius of the outermost cell and \( \epsilon = (\Delta r_{k+1}/\Delta r_k) \) is the ratio between the radii of consecutive cells. The adopted grid parameters are summarised in Table 2 for convenience. The innermost cell radius \( r_{\text{min}} \) (i.e. the inner boundary of the simulation) is chosen so that it is comparable to the star’s gravitational radius \( R_G = 2GM_\star/c_\infty^2 \) at the initial time, which is \( \sim 8.2 \times 10^{15} \text{cm} \) assuming \( T_\infty = 300 \text{K} \) and \( M_\star = 2 M_\odot \).

### Table 2. Grid parameters for our hydrodynamical simulations.

| \( N \) | \( 600 \) |
| --- | --- |
| \( \nu_{\text{min}} (\text{Hz}) \) | \( 10^{16} \) |
| \( \nu_{\text{max}} (\text{Hz}) \) | \( 2.85 \times 10^{16} \) |
| \( N_\nu \) | 50 |
| \( \epsilon \) | 1.008 |
| \( r_{\text{min}} (\text{cm}) \) | \( 10^{16} \) |
| \( r_{\text{max}} (\text{cm}) \) | \( 10^{20} \) |

We use outflow boundary conditions at both the inner and the outer boundary. The radius of the innermost cell is always larger than the stellar radius of a highly accreting protostar with a bloated envelope with \( r \leq 2 \times 10^{15} \text{cm} \) for \( M_\star \lesssim 10^5 M_\odot \) (Hosokawa et al. 2013).

We adopt a frequency grid so which allows us to follow the relevant radiative processes (Table 1). The frequency range is \( 10^{13} \text{Hz} < \nu < 2.85 \times 10^{16} \text{Hz} \) or \( 0.04 \text{eV} < h\nu < 118 \text{eV} \). The number of frequency bins is \( N_\nu = 50 \). We designed the grid layout to decrease the number of frequency bins making computation time shorter: we choose a fine frequency mesh at energies moderately larger than threshold energy of each of the reactions in Table 1 and space the bins more sparsely at other photon energies.

The initial conditions of our simulations are taken from the spherically-averaged gas cloud profiles of the LWH model in W19 (see dashed curves in their Fig. 4). These include the gas density, velocity, temperature, \( H_2 \) fraction and electron fraction. In their simulations they also show results for the cloud “MMH” which is their most massive halo. The results for LWH and MMH are qualitatively similar in our simulations so in the following we will focus on the results of the LWH cloud for simplicity. The initial profiles we adopted are shown by the black curves in Fig. 2 below. The initial \( H^+ \) fraction is set to that of electron because of the charge neutrality (note that helium is neutral at the initial condition). The ratio of the number density of hydrogen nuclei to that of helium nuclei is 0.0833. Helium is assumed to be initially all neutral.

Although in our simulations we assume a spherically symmetric gas distribution, in the 3-D simulation of W19 the gas distribution is not spherically symmetric. We discuss the possible impact of the spherical assumption on our results in §4.3.1 below.

#### 2.2 Stellar evolution

The growth of the central protostar during each time-step \( \Delta t \) is calculated simply from \( \Delta M_\star = \dot{M} \Delta t \), using the mass flux at the innermost cell as the accretion rate \( \dot{M} \) onto the protostar. The protostellar evolution is computed by fitting stellar evolution data. Depending on the accretion rate, the evolution of a rapidly growing protostar is divided into two phases: if the accretion rate is lower than \( \dot{M}_{\text{crit}}(\equiv 0.04 M_\odot \text{yr}^{-1}) \), the star is in a compact zero-age main sequence (ZAMS) phase, and otherwise it is in a bloating phase (Hosokawa et al. 2013). Even if the accretion rate drops below \( \dot{M}_{\text{crit}} \), the star may be still in the bloating phase for several thousand years (Sakurai et al. 2015). However, we do not model this sustained bloating phase, and instead assume the star is in the ZAMS phase whenever
\[ M < M_{\text{crit}} \]. As seen in §3, this treatment makes our conclusion conservative, because the EUV luminosity from the compact ZAMS model with the same mass is substantially higher than that in the bloating phase with a lower effective temperature of \( T_{\text{eff}} \sim 3000 \text{K} \). Specifically, for \( M < M_{\text{crit}} \) we use the data for ZAMS stellar evolution from Marigo et al. (2001) and Bromm et al. (2001), as summarised in Table 3. For \( M > M_{\text{crit}} \), we use the model data of a supergiant protostar growing at a constant mass accretion rate of \( \dot{M} = 10^{-5} M_{\odot} \text{yr}^{-1} \). The data of the stellar radii and effective temperature are generated by using a stellar evolution code STELLAR that was originally developed in Yorke \\& Bodenheimer (2008) and used in Sakurai et al. (2015) (see Table 4). We note that the evolution of a highly accreting protostar hardly depends on the detailed time-evolution of the mass infall rate as long as \( \dot{M} > 0.04 M_{\odot} \text{yr}^{-1} \) is satisfied.

We set the initial stellar mass to \( M_{*} = 2 M_{\odot} \), which is chosen so that the dynamical timescale at the inner-most cell \( (r_{\text{min}} = R_{B} \propto M_{*}) \) is not too short to follow gas dynamics over a wide range of spacial scales. The choice of a smaller initial mass would not make the result qualitatively different because radiative feedback does not affect the mass accretion rate until the star grows to \( \sim 4 M_{\odot} \), as seen in §3. For stellar radii and temperatures in-between the masses or outside the mass range in Tables 3 and 4, we interpolate/extrapolate linearly in the logarithmic quantities.

### 3 RESULTS

#### 3.1 Simulations with and without radiation

In the left panel of Fig. 1, we show the evolution of the accretion rates with time. The solid and dashed curves indicate simulations without radiation and with radiation respectively. The horizontal dotted line shows the critical accretion rate \( \dot{M} = 0.04 M_{\odot} \text{yr}^{-1} \) (see §1 and §2.2). In the right panel, we show the time evolution of the growing protostellar masses.

In the no-radiation case, the accretion rate remains above \( \sim 0.005 M_{\odot} \text{yr}^{-1} \) and the stellar mass grows monotonically without interruptions. In contrast, when stellar radiation is included, the radiation stops the accretion onto the protostar temporarily for \( \sim 10^{4} \text{yr} \). The growth of the stellar mass is halted at \( \sim 4 M_{\odot} \). At \( t \gtrsim 10^{8} \text{yr} \) accretion resumes, the rate eventually increases to \( \gtrsim 10^{-2} M_{\odot} \text{yr}^{-1} \), and the protostar rapidly increases its mass by \( \sim 2 \) orders of magnitude in \( \sim \) a few \( \times 10^{7} \text{yr} \). During this rapid accretion episode, the accretion rate reaches the critical rate \( \dot{M}_{\text{crit}} \) at \( t \sim 3 \times 10^{4} \text{yr} \), but decreases after the peak because the radiative luminosity due to mass accretion onto the compact ZAMS marginally exceeds the Eddington limit.

In Fig. 2, for the simulation with the radiation from the central source included, we show snapshots of the radial profiles of gas density (top left), temperature (top right), velocity (middle left), H$_{2}$ fraction (middle right), accretion rates (bottom left) and electron fraction (bottom right) when the stellar masses are \( M_{*} = 2, 4, 10, 100, 1000 \) and \( 10^{4} M_{\odot} \). The initial density profile (black line in the top left panel) has a slope \( \rho \propto r^{-1.5} \) at \( r \gtrsim 10^{19} \text{cm} \), which is shallower than the power-law \( r^{-2} \) for isothermal collapse. The shallower slope is due to sheet-like structures seen in the 3-D simulation of W19. We also show the profiles of the cooling/heating rates at the first and last snapshot in Fig. 3, where the different colors indicate the H$_{2}$ line cooling rate (black), compressional heating rate (green), atomic hydrogen line cooling rate (blue) and photoheating rate (orange).

The gas inflow is stopped at \( M_{*} \sim 4 M_{\odot} \) from \( t \sim 600 \text{yr} \) to \( \sim 7000 \text{yr} \); the infall velocity and the accretion rate become zero at \( r \lesssim 2 \times 10^{16} \text{cm} \) (purple curves in Fig. 2). At \( t \gtrsim 10^{8} \text{yr} \), the gas inflow resumes and the stellar mass increases to \( M_{*} \gtrsim 10 M_{\odot} \). In the accretion phase, the slope of the density profile in the inner regions \( r \lesssim 10^{17} \text{cm} \) gradually evolves to \( \rho \propto r^{-1} \). This change in slope occurs because the stellar gravitational radius \( R_{B} \) increases to \( \gtrsim 10^{18} \text{cm} \), making gas free-fall in the star’s point-like gravitational potential (left middle panel). The temperature reaches \( \sim 8000 \text{K} \) for \( M_{*} \gtrsim 10^{3} M_{\odot} \) in the inner region, where atomic-hydrogen cooling becomes effective (see blue curve in the bottom panel of Fig. 3). Once the core becomes hot enough to collisionally dissociate H$_{2}$ \( (T \gtrsim 3000 \text{K}) \) and the H$_{2}$ column density drops within the core, LW radiation produced from the central star propagates out and effectively dissociates H$_{2}$ in the outer region \( (r \lesssim 2 \times 10^{19} \text{cm}) \). Within the stellar influence radius \( (r < R_{B} \sim 4 \times 10^{18} \text{cm}) \), the temperature increases inward due to compressional heating (bottom panel of Fig. 3) from the equilibrium temperature of H$_{2}$ cooling \( \sim 200 \text{K} \).
As the stellar mass reaches \( \sim 10^4 \, M_\odot \), where the accretion rate is still below the critical value, a fully-ionized (\( T > 10^6 \, \text{K}; \ r \lesssim 2 \times 10^{16} \, \text{cm} \)) and partially ionized region (\( T \approx 8000 \, \text{K}; \ 2 \times 10^{16} \, \text{cm} \lesssim r \lesssim 2 \times 10^{17} \, \text{cm} \)) form. While in the partially ionized region, the electron fraction is determined by the balance between collisional ionization of neutral hydrogen (by electrons) and radiative recombination of hydrogen, the inner-most hot region is created because of photo-ionizing radiation from the \( \sim 10^4 \, M_\odot \) star with \( T_{\text{eff}} \sim 10^6 \, \text{K} \). Despite the strong stellar EUV radiation, the gas is not fully ionized to form a large H\textsc{ii} region. This is because the gas density is so high (\( n \approx 10^5 \, \text{cm}^{-3} \) at \( r \approx r_{\text{min}} \)) that the hydrogen recombination rate is faster than the ionization rate, and the H\textsc{ii} region is unable to propagate away from the stellar surface. The H\textsc{ii} fraction still remains as low as \( 10^{-11} - 10^{-7} \) because of H\textsc{ii} collisional dissociation at \( r < 2 \times 10^{17} \, \text{cm} \) and LW photodissociation at \( r > 2 \times 10^{17} \, \text{cm} \).

We examine the reason why the gas inflow is temporarily stopped and then resumes. We show the cooling/heating rates at \( t \sim 10^3 \, \text{yr} \) and \( M_* \sim 4 \, M_\odot \) when the gas inflow stops in the top panel of Fig. 3. The H\textsc{ii} cooling rate is suppressed for \( r \lesssim 10^{17} \, \text{cm} \) where the H\textsc{ii} molecules are dissociated by LW radiation (see the middle right panel of Fig. 2). Since H\textsc{ii} cooling is inefficient, and photoheating and compressional heating are effective, the temperature increases adiabatically inward following the virial temperature \( T_{\text{vir}}(\propto r^{-1}) \). Since the sonic point moves inward with this increase of the temperature, the outward gas pressure gradient force overcomes the inward gravitational force on the gas, and decelerates the infalling gas. This is seen directly in Fig. 4, where the top panel shows the radial profiles of the pressure and gravitational forces at 610 yr, at the time when the accretion first stops, and reveals that the outward pressure gradient force in the inner region becomes dominant.

The accretion rate finally resumes at \( t \gtrsim 10^4 \, \text{yr} \), because the self-gravity of the gas builds up as the outer shells fall in and accumulate. The over-pressurised region moves steadily outward from the inner core, until it reaches \( r \sim 10^{17} \, \text{cm} \) at \( t \sim 10^4 \, \text{yr} \). At the same time, the gas accumulating due to infall from larger radii increases both the inward gravitational and ram pressure forces. For example at \( r = 10^{17} \, \text{cm} \) the gravity increases from \( 1.7 \times 10^{-6} \, \text{cm} \, \text{s}^{-2} \) at \( t \sim 600 \, \text{yr} \), just after the accretion stops, to \( 2.4 \times 10^{-6} \, \text{cm} \, \text{s}^{-2} \) at \( t \sim 7000 \, \text{yr} \), just before the accretion recovers (upper vs. lower panel in Fig. 4). At \( t \approx 1.4 \times 10^4 \, \text{yr} \) (not shown in Fig. 4) the outward pressure force becomes subdominant at all radii, allowing accretion to resume.

For comparison, in Fig. 5 we show radial profiles of the density, temperature, velocity, H\textsc{ii} fraction, accretion rates and electron fraction for the no-radiation case, when the stellar masses are \( M_* = 2, 11, 100, 1000 \) and \( 10^4 \, M_\odot \). Likewise, in Fig. 6, we show profiles of cooling/heating rates when the stellar masses are \( M_* = 100 \, M_\odot \) and \( 10^4 \, M_\odot \). As time elapses and the stellar mass grows to \( M_* \gtrsim 10^4 \, M_\odot \), the slope of the density profile in the inner regions \( r \lesssim 10^{16} \, \text{cm} \) gradually evolves from the isothermal one \( \rho \propto r^{-2} \) to the free-fall one \( \rho \propto r^{-1.5} \) as seen in the case with radiation field. Because of the lack of stellar radiation feedback, the inflow velocity is accelerated to the free-fall value monotonically at all radii, and the accretion rate \( (\dot{M} \propto r^2 \rho/|v|) \) tracks its evolution without suppression. The gas temperature gradually increases inward but is saturated once it reaches \( \sim 2 \times 10^4 \, \text{K} \), because the compressional heating rate is balanced with the H\textsc{ii}-line cooling rate in the case without stellar radiation (see Fig. 6). In the central core (\( n \gtrsim 10^6 \, \text{cm}^{-3} \)), the H\textsc{ii} fraction rises rapidly through the three-body reaction (\( 3\text{H} \rightarrow \text{H}_2 + \text{H} \)) as seen in pristine star forming clouds (Yoshida et al. 2006; Turk et al. 2011).
SMS formation in a massive cloud with $H_2$ molecules

Figure 2. Radial profiles of density (top left), temperature (top right), velocity (middle left), $H_2$ fraction (middle right), accretion rates (bottom left) and electron fraction (bottom right) for the simulation with the radiation from the central source included. We show snapshots of the profiles when the stellar masses are $2$, $4$, $10$, $100$, $1000$ and $10^4 M_\odot$. The dashed line in the density panel indicates the slope $\propto r^{-1.5}$ expected for steady accretion via free-fall. The dotted line in the velocity panel indicates the free-fall velocity at $M_\star = 10^4 M_\odot$, $|v_{\text{ff}}| = \sqrt{2G(M_\star + M_{\text{enc}})/r}$. Note that for $M_\star \gtrsim 10^3 M_\odot$ the velocity of the flow is highly supersonic near the center.

4 DISCUSSION

4.1 Impact of the parameters

The main result of our simulations is that formation of SMSs can take place via rapid mass accretion because an Hii region is unable to propagate to large radii and hinder the inflow. We expect that this result may be changed by differences in the simulation setup and the initial conditions. Specifically, if the cloud had a lower initial density, the accretion rate would be smaller and the SMS formation could be inhibited if the Hii region could expand. Also, a stronger LW radiation may help gas heat more, and hinder the gas infall.

To explore how these effects would impact our conclusions (i.e. whether an SMS finally forms), we performed two variants of our fiducial simulation. First, in Fig. 7, we show the evolution of the stellar mass for a simulation in which radiation is included, but the $H_2$ self-shielding against LW radiation is ignored (dashed curve). In this case, the recovery of the accretion rate is delayed from $\sim 10^4$ yr to $\gtrsim 1$ Myr. The accretion rate then increases to $\gtrsim 1 M_\odot \text{yr}^{-1}$ and the stellar mass rapidly increases from $\sim 4 M_\odot$ to $\gtrsim 10^5 M_\odot$ within $\sim 1$ Myr. In this no-shielding model, the radiation feedback process is similar to the fiducial model: the LW radiation dissociates $H_2$ molecules, $H_2$ cooling becomes in-
efficient, the gas temperature increases and the outward gas pressure gradient force overwhelms the inward gravitational force. The longer pause in the stellar growth than in the fiducial model is due to more efficient H$_2$ dissociation by the stronger (unshielded) LW radiation.

We next show the evolution of the stellar mass for a simulation in which the initial density profile was assumed to be 10 times lower than in the fiducial case (dotted curve in Fig. 7). In this case, gas accretion is suppressed for $t \lesssim 1.5 \text{Myr}$, because the inner region of the cloud is initially gravitationally stable due to the lower density. After $t \gtrsim 1.5 \text{Myr}$, the accretion rate increases and reaches $\gtrsim 0.1 M_\odot \text{yr}^{-1}$, because gas from large scales falls inward and gravitational instability develops. The stellar mass increases from $\sim 2 M_\odot$ to $\gtrsim 10^4 M_\odot$ until $t \sim 3 \text{Myr}$.

We conclude that the SMS formation is viable if the density is larger than 0.1 times the density in the fiducial profile taken from W19, and that the limiting factor is the self-gravity of the gas in the core, rather than the radiative feedback.

4.2 Comparison to a PopIII formation case

The study presented in this paper is analogous to previous work, which assessed, via radiation-hydrodynamical simulations, the final masses of PopIII stars forming in primordial gas that has cooled and contracted inside lower-mass minihaloes (Hirano et al. 2015, 2017). The two major differences are that (i) here we consider proto-stellar cores and their surrounding initial density profiles extracted from simulations of more massive atomic-cooling haloes, and (ii) we assume spherical symmetry, and perform 1D simulations, rather than the 2D treatment in Hirano et al. (2015, 2017).
Figure 5. The same as Fig. 2 but for the simulation without radiation. We show the profiles when the stellar mass is 2, 11, 100, 1000 and $10^4 \, M_\odot$.

To check how our 1D simulation might give a different result compared to a multidimensional simulation, we perform a 1D simulation as above, but with the initial conditions adopted for a PopIII star forming cloud from Hirano et al. (2015) - specifically their cloud ID=4 and with $J_{21} = 0$ (see their Table 1). This 1D simulation can be directly compared to the 2D simulation performed in their paper. We find that in our case, the cloud forms a star with a final mass of $\sim 2000 \, M_\odot$, whereas in the 2D RHD simulation by Hirano et al. (2015) the final mass is $\sim 50 \, M_\odot$. The difference can be attributed to the fact that in the 1D case, the cloud is spherical, and radiation feedback is more easily suppressed. In the 2D simulations, radiation can escape along the lower-density polar regions, and subsequently ionize and heat the gas farther away, and more easily suppress the accretion rates at these larger distances (Tan & McKee 2001; McKee & Tan 2008). This suggests that our 1D treatment may overestimate the final stellar mass.

4.3 Caveats

4.3.1 Spherical assumption

Although we assume a spherically symmetric gas distribution in our simulations, the gas distribution in the halo LWH of W19 has non-negligible angular momentum and an asymmetric morphology. From the Extended Data Figure 4 in W19, the circular velocity of the cloud is comparable to the Keplerian velocity at radii where the enclosed mass is $M_{\text{enc}} \gtrsim 2000 \, M_\odot$, following a self-similar solution of a grav-
The same as Fig. 3 but for the simulation without radiation. We show the snapshots when the stellar mass is $M_*=100 \, M_\odot$ (top panel) and $10^4 \, M_\odot$ (bottom panel). The total cooling rate is comparable to the H$_2$ cooling rate.

Evolution of the stellar mass. Solid curve: the fiducial simulation with radiation included. Dashed: the same simulation but with LW self-shielding ignored. Dotted: the fiducial simulation with the initial density profile reduced by a factor of ten.

In the context of present-day massive star formation, in addition to stellar radiation, collimated outflows and magnetic fields are also known to suppress stellar growth (Cunningham et al. 2011; Kuiper et al. 2015, 2016; Rosen & Krumholz 2020). Outflows suppress the stellar growth rate by sweeping up interstellar material in polar directions of the star and ejecting the material from the star-forming system, as well as by decreasing the density in the polar directions and making stellar radiation feedback more effective. Strong magnetic fields also decelerate the growth rate since magnetic pressure slows down gravitational collapse of the cloud. Magnetic fields also enhance angular momentum transport by magnetic braking and inhibit the formation of a gravitationally unstable accretion disc which can cause fragmentation. These effects, however, may also suppress the accretion rate, rather than help the stellar growth (Section 3.2 of Rosen & Krumholz 2020). The effects of outflows and magnetic fields in the SMS formation case need to be investigated to clarify if they could be obstacles for the SMS formation.

In order to more robustly judge whether the protostar emerging in the core of a dynamically heated, atomic-cooling halo, can grow to a SMS, multi-dimensional simulations will need to be performed, incorporating the asymmetric distribution and nonzero angular momentum of the nearby gas.
4.3.2 Resolution of the simulations

We set the innermost cell radius to $r_{\text{min}} = 10^{16}$ cm, which is comparable to the initial protostar’s gravitational radius $R_{\text{BH}} \sim 8.2 \times 10^{15}$ cm, assuming $T_{\infty} = 300$ K and $m_\star = 2 M_\odot$ (§ 2.1). We compare this radius to the size of the HII region estimated from an equilibrium St"omgren sphere,

$$R_S = \left( \frac{3 Q_{\text{ion}}}{4 \pi n_e^2 \alpha} \right)^{1/3},$$

where $Q_{\text{ion}}$ is the ionising photon emissivity and $\alpha = 2.6 \times 10^{-13} (T_{\text{ion}}/10^4 \text{K})^{-0.85} \text{cm}^{-2} \text{s}^{-1}$ is the case-B recombination rate of hydrogen. We find $R_S \sim 1.4 \times 10^{-16}$ cm for $T_{\text{ion}} = 10^4$ K, $m_\star = 5 M_\odot$ and $Q_{\text{ion}} \sim 10^{45} \text{s}^{-1}$ (Schaerer 2002), which is much smaller than either the gravitational influence radius or the resolution of our simulation. If we estimate $R_S$ assuming a density profile $\rho \propto r^{-1.5}$ instead of a constant density, $R_S$ is even smaller, and becomes comparable to the initial stellar radius. If we resolved a region as small as $r < R_S$ in the simulations, an HII region may begin to drive an outflow and expand in the early phase. However, the HII region size (computed assuming the $\rho \propto r^{-1.5}$ profile) is about four orders of magnitude smaller than the growing protostar’s Bondi radius. We conclude that even if the HII region is less compact due to an early outflow, it is unlikely to decelerate the inflow of neutral gas in the region $R_S \lesssim r \lesssim R_B$ and radiation feedback would not be effective to hinder gas accretion (Inayoshi et al. 2016; Sakurai et al. 2016b). Even as the stellar masses grow during the evolution in our simulation, $R_S$ remains below $R_B$ by at least two orders of magnitude. Although numerical limitations preclude us from using a smaller minimum cell radius $r_{\text{min}}$ and resolving the initial ultra-compact HII-region, we expect that our main result, i.e. that the radiation feedback is ineffective, is not compromised by this limitation.

5 SUMMARY

W19 argued, based on three-dimensional cosmological simulations, that SMSs may form in large atomic-cooling haloes in which H$_2$ molecules are not fully dissociated by external FUV radiation. In this work, we followed the evolution of a protostar identified in one of the haloes (specifically, the halo “LWH” in W19), beyond the point where their simulation stopped. We performed 1-D radiation hydrodynamical simulations to explore if radiation feedback suppresses the growth of this protostar. We solved the non-equilibrium chemical reactions of nine primordial species, and included the radiation of the central source derived from stellar evolution models, as well as radiation from a circumstellar disk.

We found that a SMS with a mass of $\gtrsim 10^5 M_\odot$ forms, even though stellar radiation feedback temporarily halts the accretion for $\sim 10^5$ yr. This feedback is caused by LW radiation from the protostar. The LW radiation dissociates H$_2$ in the inner region, increasing the gas temperature and the gas pressure gradient force which opposes gravity. The feedback stops after $\sim 10^5$ yr, because the gas self-gravity and inward ram pressure force of the gas building up on larger scales overcome the outward pressure gradient force. Although the stellar UV radiation is strong, no HII region forms during the evolution because of the high density and efficient hydrogen recombination. We conclude that the protostar can grow to $M_\star \gtrsim 10^5 M_\odot$, as long as the central density is at least $\sim$10% of the value found in W19. The main caveat to this conclusion is our assumption of spherical symmetry; radiation may have a stronger effect on an asymmetric collapse. Multi-dimensional simulations will be required to include these effects and to assess the robustness of our results.

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