Impact of cosmic rays on Population III star formation

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ABSTRACT

We explore the implications of a possible cosmic-ray (CR) background generated during the first supernova explosions that end the brief lives of massive Population III stars. We show that such a CR background could have significantly influenced the cooling and collapse of primordial gas clouds in minihaloes around redshifts of $z \sim 15-20$, provided the CR flux was sufficient to yield an ionization rate greater than about $10^{-19}$ s$^{-1}$ near the centre of the minihalo. The presence of CRs with energies $\lesssim 10^7$ eV would indirectly enhance the molecular cooling in these regions, and we estimate that the resulting lower temperatures in these minihaloes would yield a characteristic stellar mass as low as $\sim 10 \, M_\odot$. CRs have a less-pronounced effect on the cooling and collapse of primordial gas clouds inside more massive dark matter haloes with virial masses $\gtrsim 10^9 \, M_\odot$ at the later stages of cosmological structure formation around $z \sim 10-15$. In these clouds, even without CR flux the molecular abundance is already sufficient to allow cooling to the floor set by the temperature of the cosmic microwave background.

Key words: molecular processes – stars: formation – galaxies: formation – cosmology: theory – early Universe.

1 INTRODUCTION

How did the cosmic dark ages end, and how was the homogeneous early Universe transformed into the highly complex state that we observe today (e.g. Barkana & Loeb 2001; Miralda-Escudé 2003; Ciardi & Ferrara 2005)? One of the key questions is to understand the first stars, the so-called Population III (or Pop III), since they likely played a crucial role in driving early cosmic evolution (e.g. Bromm & Larson 2004; Glover 2005). For instance, the radiation from the first stars contributed to the reionization of the intergalactic medium (IGM) (e.g. Kitayama et al. 2004; Whalen, Abel & Norman 2004; Alvarez, Bromm & Shapiro 2006; Johnson, Greif & Bromm 2007), leading to the end of the cosmic dark ages. After the first stars exploded as supernovae (SNe), they spread heavy elements through the IGM, thereby providing the initial cosmic metal enrichment (e.g. Madau, Ferrara & Rees 2001; Mori, Ferrara & Madau 2002; Bromm, Yoshida & Hernquist 2003; Wada & Venkatesan 2003; Norman, O’Shea & Paschos 2004; Greif et al. 2007). Furthermore, this IGM metallicity was likely crucial in governing the transition from early high-mass star formation to the normal-mass star formation seen today (e.g. Omukai 2000; Bromm et al. 2001; Schneider et al. 2002; Bromm & Loeb 2003; Frebel, Johnson & Bromm 2007; Smith & Sigurdsson 2007).

SN explosions are also thought to be the site of cosmic-ray (CR) production (e.g. Ginzburg & Syrovatskii 1969), and the CR background built up from the first SNe may affect the cooling and collapse of primordial gas clouds. Earlier studies (Shchekinov & Vasiliev 2004; Vasiliev & Shchekinov 2006) have shown how the presence of CRs at high redshifts could have lowered the minimum mass at which primordial gas could cool and collapse into virialized dark matter (DM) haloes (e.g. Haiman, Thoul & Loeb 1996; Tegmark et al. 1997). This earlier study assumes the existence of ultra-heavy X particles that decay into ultra-high-energy CRs which interact with the cosmic microwave background (CMB), leading to the production of ionizing photons. In this paper, however, we instead examine CR effects which occur through direct collisional ionization of neutral hydrogen. The free electrons created from the ionization act as a catalyst for the formation of H$_2$ (McDowell 1961). Because molecular hydrogen emits photons through its rovibrational transitions, H$_2$ is able to cool primordial gas at temperatures below the threshold for atomic hydrogen cooling ($\lesssim 10^4$ K). Furthermore, CR ionization can indirectly lead to higher abundances of HD, which also acts as a cooling agent in low-temperature primordial gas (see Johnson & Bromm 2006). However, the presence of CRs can lead to ionization heating, as well. Whether this direct heating effect is strong enough to counter the additional cooling must be determined and will depend on the high-redshift CR energy density. Similar to our work, Rollinde et al. (2005) and Rollinde et al. (2006) used models of early star formation to estimate the CR energy density in the early Universe, constraining it with the observed Li abundances in metal-poor Galactic halo stars.

In this paper, we will investigate the importance of the CR feedback on Pop III star formation by modelling its effect on the cooling of primordial gas in two cases: collapse within minihaloes, and shocks associated with the virialization of more massive DM haloes.

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during the later stages of structure formation. The outline for this paper is as follows. In Section 2, we discuss CR acceleration and propagation in the high-redshift Universe and how these might differ from the present-day case. Section 3 describes the evolution of primordial gas in minihaloes and in virialization shocks when accounting for the effects of CRs. For the minihalo case, we further discuss how the fragmentation scale could change for a sufficiently high CR flux. We present our conclusions in Section 4.

2 COSMIC RAYS IN THE HIGH-Z UNIVERSE

2.1 Population III star formation

Though CR effects will be examined for a range of star formation rates, the typical Pop III star formation rate is taken to be that found in Bromm & Loeb (2006) at $z \approx 15$, which is approximately $\Psi_\star \approx 2 \times 10^{-5} M_\odot$ yr$^{-1}$ Mpc$^{-3}$ in a comoving volume. This rate was derived using the extended Press–Schechter formalism (Lacey & Cole 1993) to model the abundance and merger history of cold dark matter (CDM) haloes. In a neutral medium, before the redshift of reionization, Bromm & Loeb (2006) assume that star formation occurs only in haloes that have become massive enough to enable atomic line cooling with virial temperatures above approximately $10^4$ K. Recent work (see Greif & Bromm 2006) suggests that these more massive haloes were indeed the dominant site for star formation. Greif & Bromm (2006) argue that about 90 per cent of the mass involved in metal-free star formation initially cooled through atomic line transitions.

At higher redshifts such as $z \approx 20$, however, the first stars are thought to have formed inside $\sim 10^6 M_\odot$ minihaloes through molecular cooling, and this mode of star formation is much more significant at this time. For the minihalo case at $z \approx 20$, Yoshida et al. (2003) estimate the rate of star formation through H$_2$ cooling to be $\Psi_\star \sim 10^{-3} M_\odot$ yr$^{-1}$ Mpc$^{-3}$. To account for such differences in these determinations of star formation rates, we examine a range of values spanning multiple orders of magnitude.

The Pop III initial mass function (IMF) currently remains highly uncertain, so for this study we do not perform our calculations using a specific IMF. For simplicity, we instead assume that those Pop III stars whose masses lie in the pair instability SN (PISN) range (140–260 $M_\odot$) have an average mass of 200 $M_\odot$ (e.g. Heger et al. 2003). A Pop III IMF that extends over a large range of masses would imply that only a fraction of these stars were in the PISN range. Thus, here we assume that only slightly less than half of Pop III stars are in this mass range, leading to a somewhat more conservative value for the CR energy density. Our estimate generally corresponds to an IMF peaked around 100 $M_\odot$. Due to their high mass, the first stars had very short lifetimes of about 3 Myr (e.g. Bond, Arnett & Carr 1984), and we therefore assume instantaneous recycling of the stellar material. As described below, this overall picture of Pop III star formation will be used to estimate the average CR energy density in the high-redshift Universe.

2.2 Cosmic-ray production

For this paper, CRs are assumed to have been generated in the PISNe that may have marked the death of Pop III stars. The CRs are accelerated in the SN shock wave through the first-order Fermi process by which high-energy particles gain a small percentage increase in energy each time they diffuse back and forth across a shock wave. This yields a differential energy spectrum in terms of CR number density per energy (e.g. Longair 1994):

$$\frac{dn_{CR}}{d\epsilon} = \frac{n_{norm}}{\epsilon_{min}} \left( \frac{\epsilon}{\epsilon_{min}} \right)^{-x} ,$$

where we will use $x = -2$ for definiteness, a typical value given by the Fermi acceleration theory (e.g. Bell 1978a). Here, $\epsilon$ is the CR kinetic energy, $\epsilon_{min}$ is the minimum kinetic energy, and $n_{CR}$ is the CR number density. Note that our results are somewhat sensitive to the choice of $x$. Choosing values for $x$ that are closer to what is observed in the Milky Way (MW), such as $x = -2.5$ or $-3$, or using a CR spectrum similar to that given in Rollinde et al. (2005) and Rollinde et al. (2006) will increase CR heating and ionization, if the overall CR energy density is held constant. For these steeper power laws, a higher fraction of the total CR energy resides in the lower energy CRs, which are the ones that contribute the most to these CR effects, as discussed in Sections 3.1 and 3.4 and in the discussion for Fig. 3. Such steeper spectral slopes can in fact yield CR heating and ionization rates up to an order of magnitude higher than that for $x = -2$. Using $x = -2$ is therefore a more conservative choice that does not assume any modifications to the standard Fermi acceleration theory.

By equating the total CR energy density $U_{CR}$ with the integral of the differential CR spectrum over all energies, the normalizing density factor $n_{norm}$ is estimated to be

$$n_{norm} = \frac{U_{CR}}{\epsilon_{min} \ln(\epsilon_{max}/\epsilon_{min})} \approx \frac{1}{10} \frac{U_{CR}}{\epsilon_{min}} ,$$

where we get approximately 1/10 for the coefficient choosing $\epsilon_{max} = 10^{15}$ eV and $\epsilon_{min} = 10^6$ eV. The differential energy spectrum is therefore

$$\frac{dn_{CR}}{d\epsilon} = \frac{\epsilon^{x}}{\epsilon_{min} \ln(\epsilon_{max}/\epsilon_{min})} \left( \frac{\epsilon}{\epsilon_{min}} \right)^{-2} ,$$

with

$$U_{CR}(z) \approx \frac{p_{CR}}{\epsilon_{min}} E_{SN, f_{PISN}} \Psi_\star(1 + z)^3.$$  

This can also be written as

$$U_{CR}(z) \approx 2 \times 10^{-15} \text{ erg cm}^{-3} \left( \frac{p_{CR}}{0.1} \right) \left( \frac{E_{SN}}{10^{52} \text{ erg}} \right) \left( \frac{1 + z}{21} \right)^{3/2} \times \left( \frac{f_{PISN}}{2 \times 10^{-2} M_\odot \text{ yr}^{-1} \text{ Mpc}^{-3}} \right),$$

where $p_{CR}$ is the fraction of SN explosion energy, $E_{SN}$, that goes into CR energy, and $f_{PISN}$ is the number of PISNe that occur for every solar mass unit of star-forming material. Here, we take the above star formation rate to be constant over a Hubble time $t_H$, where

$$t_H(z) \approx 2 \times 10^7 \text{ yr} \left( \frac{1 + z}{21} \right)^{-3/2} ,$$

evaluated at the relevant redshift. We assume that each star quickly dies as a PISN with $E_{SN} = 10^{52}$ erg, appropriate for a 200-$M_\odot$ star (Heger & Woosley 2002), 10 per cent of which is transformed into CR energy (e.g. Ruderman 1974). The value of 10 per cent is derived from MW energetics, and here we have simply extrapolated this to PISNe. Very little is known about what value of $p_{CR}$ applies to PISNe specifically, so assuming their shock structure to be similar to that of local SNe appears to be a reasonable first guess. We choose $1/500 M_\odot$ for $f_{PISN}$, so that there is one PISN for every 500 $M_\odot$ of star-forming material. This implies that somewhat less than half of the star-forming mass falls within the PISN range.

Compared to PISNe, the usual core-collapse SNe (CCSNe) thought to accelerate CRs in the MW have an explosion energy...
that is lower by about an order of magnitude. However, the masses of their progenitor stars are also much lower, ranging from \(\sim 10\) to \(40 \, M_\odot\). Thus, if we have an IMF extending to this lower mass range, there will be approximately an order of magnitude more SNe per unit mass of star-forming material. However, because \(E_{\text{SN}}\) for low-mass \((\sim 10 \, M_\odot)\) CCSNe will be lower by a similar amount, these two effects may cancel out and result in a total \(U_{\text{CR}}\) that is comparable to the PISN case. Furthermore, the difference in shock velocities, \(M_{\text{sh}}\), between a PISN and a CCSN should not be more than a factor of a few. Though CCSNe have a lower explosion energy, this energy is used to accelerate roughly 10 times less ejected mass, \(M_\text{ej}\), than in a PISN. Simple energy conservation, \(E_{\text{SN}} \approx \frac{4}{3} M_{\odot} \epsilon_{\text{SN}}\), then indicates that the shock velocities of these different SNe will be similar. This also leads to similar estimates for \(\epsilon_{\text{max}}\) for both cases since the minimum CR energy is expected to depend on shock velocity. For a given star formation rate, the results of our study would therefore change little if the sources of CRs were CCSNe instead of PISNe.

It should be pointed out that the explosion mechanism of the highest-mass CCSN progenitors, \(\sim 40 \, M_\odot\), is still somewhat uncertain and may be associated with very high explosion energies comparable to that of PISNe, though if most of the CCSNe derive from lower mass progenitors they should still have lower explosion energies as argued above. Furthermore, we emphasize that the choice of PISNe versus CCSNe is not crucial for our study as long as there exists some source of sufficient CR production. These sources do not necessarily have to be only PISNe. Constraints on the number of PISNe in the early Universe, such as provided by the problem of overproduction of metals in the IGM by PISNe (e.g. Venkatesan & Truran 2003; Tumlinson, Venkatesan & Shull 2004; Daigne et al. 2006), will therefore not affect the relevance of this study. Finally, it is worth noting that the CR energy densities used in this paper are well within upper limits placed by previous studies. For instance, one of the constraints that can be used to place an upper limit on \(U_{\text{CR}}\) is the \(^6\text{Li}\) plateau observed in metal-poor Galactic halo stars. While a number ratio of \(^6\text{Li}/\text{H} \approx 10^{-11}\) is observed (Asplund et al. 2006), the much smaller ratio \(^6\text{Li}/\text{H} \approx 10^{-14}\) is predicted by big bang nucleosynthesis (BBN). In Rollinde et al. (2005) and Rollinde et al. (2006), the overabundance of \(^6\text{Li}\) compared to predictions from BBN is assumed to have come from CRs generated by the first stars. In Rollinde et al. (2005), this assumption was used to derive an upper limit on the high-z CR energy density. Assuming a model with a single burst of Pop III star formation, they found CR energy densities at the time of these bursts to range from approximately \(10^{10}\) to \(10^{12}\) erg cm\(^{-3}\) for burst times ranging from \(z = 100\) to 10. Such limits are roughly three orders of magnitude larger than the \(U_{\text{CR}}\) values used in our investigation, so that we do not violate any known constraints on CR production in the early Universe.

2.3 High-redshift Greisen–Zatsepin–Kuzmin cut-off

The Greisen–Zatsepin–Kuzmin (GZK) cut-off is an upper limit to the energy of a CR if it is extragalactic in origin and thus has travelled through the sea of CMB photons (Greisen 1966; Zatsepin & Kuzmin 1966). This limit exists because of interactions between CRs and CMB photons which lead to photo-pion and pair production, reducing the energy of the CR for each interaction. The reactions for pair production and photo-pion production, respectively, are (e.g. Berezinsky & Grigorieva 1988):

\[
p^+ + \gamma \rightarrow p^+ + e^+ + e^-, \tag{7}
\]

\[
p^+ + \gamma \rightarrow \pi^0 + p^+. \tag{8}
\]

\[
p^+ + \gamma \rightarrow \pi^0 + p^+. \tag{9}
\]

Photo-pion production is the most important interaction for CR protons (De Marco 2005). However, this reaction can only take place if the energy of a CMB photon is above the energy threshold of \(\epsilon_\gamma = 140\) MeV, but this is possible in the rest frame of CRs with sufficiently large Lorentz factors (e.g. Longair 1994). The reaction will continue until the CR energy falls below the corresponding threshold, which is \(\epsilon_{\text{GZK}} = 5 \times 10^{19}\) eV in today’s Universe. This cut-off has recently been observed by the HiRes experiment (Abbasi et al. 2007).

In the high-redshift Universe, however, the GZK cut-off will be somewhat lower, as can be seen in the following: while in today’s Universe the average energy of a CMB photon is \(\epsilon_{\text{CMB}} = 2.7\, \text{kB} / (1 + z) \approx 6 \times 10^4\) eV, at higher redshifts this energy will be larger by a factor of \((1 + z)\). In the CR rest frame, the CMB photon energy is

\[
\epsilon_{\text{GZK}} \approx 6 \times 10^4\, \text{eV}\times (1 + z), \tag{10}
\]

\[
\gamma \approx 2 \times 10^{11} \times (1 + z). \tag{11}
\]

We can now calculate the CR energy for which the threshold for photo-pion production is reached:

\[
\epsilon_{\text{GZK}}(z) = \gamma m_\text{p} c^2 \approx 3 \times 10^{20}\, \text{eV} \times \frac{1 + z}{(1 + z)} \text{ eV}. \tag{12}
\]

Thus, at redshifts of 10 or 20, the GZK cut-off is around an order of magnitude smaller than in today’s Universe, giving a robust upper limit to the CR energy.

The GZK cut-off applies only to those CRs that travel large distances through the CMB. A CR can undergo a maximum of approximately 10 interactions with CMB photons before falling below the GZK limit (e.g. Longair 1994). Given an interaction cross-section of \(\sigma = 2.5 \times 10^{-28}\) cm\(^2\) for photo-pion production and a CMB density of \(n_\gamma = 5 \times 10^5\) cm\(^{-3}\) \((1 + z)^3\), the interaction mean-free-path is

\[
\lambda = (\sigma n_\gamma)^{-1} \approx \frac{10^{23}\text{ cm}}{(1 + z)^3}. \tag{14}
\]

The maximum distance from which ultrahigh energy CRs (UHECRs) with \(\epsilon_{\text{CR}} \geq 5 \times 10^{19}\, \text{eV}/(1 + z)\) could have originated is therefore

\[
d_\text{max} \approx 10 \lambda \approx 10^{26}\, \text{cm} \times \frac{1}{(1 + z)^3} \approx 3\, \text{kpc} \left(\frac{1 + z}{21}\right)^{-3}. \tag{15}
\]

Thus, at \(z = 20\) a UHECR impinging on a primordial gas cloud must have been accelerated in a source no more than a proper distance of \(~3\) kpc away. The origin of UHECRs remains unknown even in today’s Universe, though possible sources range from pulsar winds to active galactic nuclei (AGN) and gamma-ray bursts (GRBs) (e.g. de Gouveia Dal Pino & Lazarian 2001; Waxman 2001; Torres et al. 2002). Though structures at \(z = 15\) and 20 are not yet massive enough to form AGN, pulsars and GRBs may be plausible UHECR sources at these redshifts (e.g. Bromm & Loeb 2006).

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\[\text{p}^+ + \gamma \rightarrow \pi^+ + e^+ + e^-, \tag{7}\]

\[\text{and}\]

\[\text{p}^+ + \gamma \rightarrow \pi^0 + p^+. \tag{8}\]

\[\text{p}^+ + \gamma \rightarrow \pi^0 + p^+. \tag{9}\]
2.4 Magnetic fields

Magnetic fields are an important component of CR studies due to their effects on both CR acceleration and propagation through the Universe. For instance, the maximum energy a CR can reach with the Fermi acceleration process has a linear dependence on the ambient magnetic field strength if the growth rate of the CR’s energy is limited by its Larmor radius (Lagage & Cesarsky 1983). Under this assumption, strengths of $\sim 10^{-10}$ G are necessary to accelerate a CR to $10^9$ eV (Zweibel 2003).

The strength, generation, and dispersal of magnetic fields at high redshifts, however, are still highly uncertain. One of the most widely held views is that magnetic seed fields were created soon after the big bang and were later amplified in higher density structures through dynamo mechanisms. Seed fields are thought to form through various processes including galaxy-scale outflows in the early Universe and the Biermann battery mechanism in regions such as shock waves and ionization fronts (e.g. Kronberg, Lesch & Hopp 1999; Gnedin, Ferrara & Zweibel 2000). For one illustrative example, Ichiki et al. (2006) propose a mechanism by which seed fields are created before recombination through second-order cosmological perturbations. They calculate the seed field at the redshift of recombination, $z_{\text{rec}} \approx 10^3$, to be

$$B_0(z_{\text{rec}}) \approx 10^{-14} \left( \frac{\lambda}{10 \, \text{kpc}} \right)^{-2} \, \text{G},$$

where $\lambda$ is the comoving size of the structure in question for scales less than about 10 Mpc. If we take into account magnetic flux freezing as structures become dense and the Universe expands, we get a seed field at virialization of

$$B_0(z_{\text{vir}}) \approx 10^{-20} \left( \frac{\lambda}{10 \, \text{kpc}} \right)^{-2} \left( \frac{\rho_0}{\rho_s} \right)_{z_{\text{vir}}}^{2/3} \left( 1 + z_{\text{vir}} \right)^2,$$

where $(\rho_0/\rho_s)_{z_{\text{vir}}} \approx 200$ is the overdensity of the structure in question compared to the average density of the Universe at $z_{\text{vir}}$, the redshift at which the structure first virializes. For a minihalo with a comoving size of around 10 kpc and virialization redshift $z_{\text{vir}} = 20$, this gives $B_0 \approx 10^{-18} \, \text{G}$. When dynamo effects are taken into account, the magnetic field is further amplified and can grow exponentially on a time-scale determined by differential rotation and turbulence within the structure (e.g. Field 1995; Widrow 2002). Magnetic field amplification within the accelerating region of a SN remnant itself may also increase its strength by up to two orders of magnitude (e.g. van der Lann 1962; Bell & Lucek 2001). This magnetic field growth can occur through processes such as flux freezing in the compressed regions of the SN remnant and non-linear amplification through growth and advection of Alfvén waves generated by the pressure of CRs themselves. Within structures that have already experienced star formation, magnetic fields can further be built up through field ejection in stellar winds, SN blastwaves, and protostellar jets (e.g. Machida et al. 2006).

For these fields to be spread into the general IGM, however, there must be a sufficient degree of turbulent mixing and diffusion in intergalactic regions. While such processes are effective within structures, it is less obvious that they are also effective in the IGM at early times (Johnson & Bromm 2006; see also Mackey, Bromm & Hernquist 2003). The $z \approx 20$ collapsing minihalo case, typical halo masses are $\sim 10^5 \, M_\odot$. The minihalos form during hierarchically merging halos at velocities too low to cause any shock or ionization in the minihalo, so the electron-catalyzed molecule formation and cooling is insufficient to allow stars less than around $100 \, M_\odot$ to form (e.g. Bromm, Coppi & Larson 1999; Bromm, Coppi & Larson 2002; Abel, Bryan & Norman 2002). CR ionization and heating, however, can potentially alter the thermal and chemical evolution of primordial gas.

In our model, the minihalo has an initial ionization fraction of $x_e = 10^{-4}$ and undergoes free-fall collapse. Its density therefore evolves according to $dn/dt = n/t_{ff}$, with the free-fall time being

$$t_{ff} = \left( \frac{3\pi}{32G\rho} \right)^{1/2},$$

where $\rho = \mu m_n n \approx m_n n$ and $\mu = 1.2$ is the mean molecular weight for neutral primordial gas. The initial density was taken to be the density of baryons in DM haloes at the point of virialization (e.g. Clarke & Bromm 2003):

$$n_0 \approx 0.3 \, \text{cm}^{-3} \left( \frac{1 + z}{20} \right)^3.$$

The initial temperature of the gas was taken to be 200 K, and the initial abundances were the primordial ones (e.g. Bromm, Coppi & Larson 2002).

As our second case, we consider strong virialization shocks that arise in later stages of structure formation during the assembly of the first dwarf galaxies. The corresponding DM haloes virialize at $z \sim 10-15$, and have masses ranging from $\sim 10^6$ to $\sim 10^{10} M_\odot$. The conditions during the assembly of the first dwarf galaxies show some key differences from the former case. The virial velocities of these DM haloes are much greater, implying merger velocities that are now high enough to create a shock that can partially ionize the primordial gas (Johnson & Bromm 2006 and references therein). The post-shock evolution is taken to be roughly isobaric (e.g. Shapiro & Kang 1987; Yamada & Nishi 1998). We use an initial post-shock...
temperature of
\[ T_{\text{ps}} = \frac{m_p u_{\text{m}}^2}{3 k_B}. \]  
(20)

For an initial density of \( n_{\text{ps}} \), again found from equation (19), the temperature and density will follow the relation \( T_{\text{ps}} n_{\text{ps}} \propto T n \).

### 3.2 Thermal and chemical evolution

In calculating the evolution of the primordial clouds, we solve the comprehensive chemical reaction network for all the species included in Johnson & Bromm (2006), and consider cooling due to H, H\(_2\), and HD. The temperature of the CMB sets the lower limit to which the gas can cool radiatively (e.g. Larson 1998). Assuming that CRs with the above energy spectrum impinge on the primordial gas cloud, we add the respective heating and ionization rates. Once a low-energy CR enters the high-density region of the cloud, each time it ionizes an H atom, an electron with average energy \( E = 35 \) eV is released (Spitzer & Tomasko 1968). Including the ionization energy of 13.6 eV implies that a CR proton loses approximately 50 eV of kinetic energy upon each scattering. This places a limit on the number of scatterings a CR can undergo in a cloud as well as a limit on the distance into the cloud that it can reach before it loses all its energy to ionization. This distance can be described by a penetration depth

\[ D_p(\epsilon) \approx \frac{\beta c}{-(\text{de/dr})_{\text{ion}}}, \]

where (Schlickeiser 2002)

\[ -(\text{de/dr})_{\text{ion}} = 1.82 \times 10^{-7} \text{ eV s}^{-1} n_{H^0} f(\epsilon), \]

(22)

\[ f(\epsilon) = (1 + 0.0185 \ln \beta) \left( \frac{2 \beta^2}{\beta_0^2 + 2 \beta^2} \right)^{-1/2}, \]

(23)

and

\[ \beta = \left[ 1 + \left( \frac{\epsilon}{m_p c^2} + 1 \right)^{-2} \right]^{1/2}. \]

(24)

Here, \( m_p \) is the mass of a proton, \( \beta = v/c \), and \( -(\text{de/dr})_{\text{ion}} \) is the rate of CR energy loss due to ionization. The cut-off value of \( \beta_0 \approx 0.01 \) is appropriate for CRs travelling through a medium of atomic hydrogen, since \( \beta_0 c = 0.01c \) is the approximate orbital velocity of electrons in the ground state of atomic hydrogen. When the velocity of CRs falls below \( \beta_0 c \), the interaction between the CRs and electrons will sharply decrease, as will the ionization rate (Schlickeiser 2002), but for our study all CRs are assumed to be above this critical velocity. Thus, for a given distance \( D \) into a cloud, a CR has an effective optical depth of \( D/D_p \). Fig. 1 shows the dependence of \( D_p \) on CR energy for a neutral hydrogen density of \( n_{H^0} = 1 \) cm\(^{-3} \), which is typical for densities in the minihalo case. Densities will of course greatly increase towards the end of the freefall evolution, but the size of the collapsing gas cloud, and thus the distance CRs must travel, will decrease. Although the overall attenuation would be slightly greater if the time-dependent density evolution were accounted for instead of using a constant attenuation value, when comparing these two cases the difference is not large enough to yield a significant variation in the minihalo’s temperature evolution. For simplicity, only the typical density was therefore used in calculating the attenuation.

As can be seen in Fig. 1, the lowest-energy CRs do not get attenuated until they travel a distance of about a few hundred pc, so in gas clouds of this size or smaller the CR flux will not be significantly attenuated. This also shows that the low-energy CRs are the ones that will have the greatest ionization and heating contribution to the cloud, as they more readily release their energy into the gas. In contrast, higher energy CRs will quickly travel through a minihalo without transferring much of their energy into the gas. They instead lose energy more slowly over much longer distances. Accounting for the attenuation yields CR ionization and heating rates of

\[ \Gamma_{\text{CR}}(D) = \frac{E_{\text{heat}}}{50 \text{ eV}} \int_{n_{\text{min}}}^{n_{\text{max}}} \left( \frac{\text{de}}{\text{dr}} \right)_{\text{ion}} \frac{n_{H^0}}{e^{-D/D_p} \text{de}}, \]

(25)

and

\[ \xi_{\text{CR}}(D) = \frac{\Gamma_{\text{CR}}}{n_{H^0} E_{\text{heat}}}, \]

(26)

respectively. These rates can also be written as

\[ \Gamma_{\text{CR}}(D) = 5 \times 10^{-26} \text{ erg cm}^{-3} \text{s}^{-1} \left( \frac{U_{\text{CR}}}{2 \times 10^{-15} \text{ erg cm}^{-3}} \right) \]

\[ \times \left( \frac{E_{\text{heat}}}{6 \text{ eV}} \right) \left( \frac{n_{H^0}}{1 \text{ cm}^{-3}} \right) \left( \frac{\epsilon_{\text{min}}}{10^6 \text{ eV}} \right)^{-1} I(\epsilon), \]

(27)

and

\[ \xi_{\text{CR}}(D) = 5 \times 10^{-18} \text{s}^{-1} \left( \frac{U_{\text{CR}}}{2 \times 10^{-15} \text{ erg cm}^{-3}} \right) \]

\[ \times \left( \frac{\epsilon_{\text{min}}}{10^6 \text{ eV}} \right)^{-1} I(\epsilon), \]

(28)

respectively, where

\[ I(\epsilon) = \int_{n_{\text{min}}}^{n_{\text{max}}} f(\epsilon) \left( \frac{\epsilon}{\epsilon_{\text{min}}} \right)^x e^{-D/D_p} \text{de} \]

(29)

The factor \( E_{\text{heat}} \) is equal to 6 eV for the minihalo case because, though CRs lose about 50 eV of energy after each ionization, only about 6 eV of that energy goes towards heating in a neutral medium (see Spitzer & Scott 1969; Shull & van Steenberg 1985). The value of \( E_{\text{heat}} \) increases for media with larger ionization fractions due to an increase in Coulomb interactions between the newly freed electrons and the medium, and for the ionization fractions typical of the virialization shock case, one has \( E_{\text{heat}} \approx 26 \) eV.
The minimum kinetic energy, \( \epsilon_{\text{min}} \), with which the CRs impinge on the gas clouds can be roughly estimated by assuming that during Fermi acceleration a particle will acquire the velocity of the shock wave itself after crossing it a single time (see Bell 1978b). This gives a minimum energy \( \epsilon_{\text{min}} \approx \frac{1}{2} m_H u_{\text{sh}} \approx 10^6 \) eV, where \( u_{\text{sh}} \approx 10^9 \) km s\(^{-1}\) is the PISN shock velocity in its initial blast-wave stage. The effects of changing \( \epsilon_{\text{min}} \) will be discussed later. Due to the power-law distribution of CR energies, the value of the maximum CR kinetic energy, \( \epsilon_{\text{max}} \), has less importance since the number densities and ionization rates of highly relativistic CRs are much smaller than those of non-relativistic and marginally relativistic ones. We therefore choose a default value of \( \epsilon_{\text{max}} = 10^{15} \) eV, a typical maximum value determined from the Fermi acceleration theory (e.g. Blandford & Eichler 1987).

Fig. 2 shows the gas cloud evolution for the minihalo collapse case for models with and without CR effects at \( z = 20 \). The heating rate is evaluated at \( D = 100 \) pc, and the characteristic distance to the cloud centre, as this is where star formation is expected to take place. Various combinations of \( \Psi_\star = 2 \times 10^{-3} \) or \( 2 \times 10^{-2} \) M\(_\odot\) yr\(^{-1}\) Mpc\(^{-3}\) and \( \epsilon_{\text{min}} = 10^6, 10^5 \) or \( 10^4 \) eV are considered. The resulting CR ionization leads to an increase in the electron abundance in the cloud. These electrons act as the only catalyst for H\(_2\) formation since at this early time in the Universe there are no dust grains on which H\(_2\) could form. The increased electron abundance thus allows for more H\(_2\) to form. H\(_2\), in turn, can also be used in the main reaction that creates HD (see, e.g. Johnson & Bromm 2006). With the increase in molecular abundance due to CR ionization, the cloud is able to cool to temperatures much closer to the CMB floor if \( \Psi_\star = 2 \times 10^{-3} \) M\(_\odot\) yr\(^{-1}\) Mpc\(^{-3}\) and \( \epsilon_{\text{min}} = 10^6 \) or \( 10^5 \) eV. This indirect cooling effect is stronger than the direct CR heating effect even for star formation rates up to 100–1000 times higher than those shown here. Furthermore, for \( \Psi_\star \) as low as \( 2 \times 10^{-4} \) M\(_\odot\) yr\(^{-1}\) Mpc\(^{-3}\), the CR effects are negligible for all \( \epsilon_{\text{min}} \) values explored in this study.

To further illustrate how the level of CR-induced cooling can vary with CR parameter values, Fig. 3 shows the minimum temperature \( T_{\text{min}} \) reached by the minihalo gas versus \( U_{\text{CR}} \) (top panel) and spectral slope \( x \) (bottom panel). The dotted lines denote the CMB temperature, and the dashed line in the upper panel is the upper bound on \( U_{\text{CR}} \) based on constraints from the \(^6\)Li measurements. Here \( \epsilon_{\text{min}} \) is kept constant at \( 10^6 \) eV. Note that the cloud cools nearly to the CMB floor for \( U_{\text{CR}} \) values ranging over almost four orders of magnitude. CR-induced heating starts to overcome cooling effects and the minimum temperature begins to rise only once \( U_{\text{CR}} \) values are near the \(^6\)Li constraint. For shallower spectral slopes, and thus smaller numbers of low-energy CRs for a given \( U_{\text{CR}} \), CRs cause little cooling or heating and do not change the minimum temperature of the minihalo gas.
Recalling the dependence of $U_{\text{CR}}$ the same values used thus far in this study. Following condition is met: 

\[ \text{CR-induced cooling will be significant in the minihalo, if the fol-} \]

\[ \text{tion (4), we find that there will be sufficient ionization, so that the} \]

\[ \text{average heating and ionization rates for each case show that CRs can facilitate cooling in the} \]

\[ \text{minihalo to nearly the CMB floor if the ionization rate is greater} \]

\[ \text{than approximately } 10^{-15} \text{s}^{-1}. \text{Any rate below this yields negligible} \]

\[ \text{cooling. Considering the dependence of the ionization rate on} \]

\[ \text{10^{-15} \text{erg cm}^{-3}} \right) \left( \frac{1}{21} \right)^{3/2} \left( \frac{U_{\text{CR}}}{10^7 \text{eV}} \right)^{-1.3} \gtrsim 1. \quad (31) \]

Recalling the dependence of $U_{\text{CR}}$ on redshift and $\Psi_\ast$ from equation (4), we find that there will be sufficient ionization, so that the CR-induced cooling will be significant in the minihalo, if the following condition is met:

\[ \left( \frac{\Psi_\ast}{10^{-2} M_\odot \text{yr}^{-1} \text{Mpc}^{-3}} \right) \left( \frac{1+z}{21} \right)^{3/2} \left( \frac{\epsilon_{\text{min}}}{10^7 \text{eV}} \right)^{-1.3} \gtrsim 1. \quad (31) \]

Here, we assume that all other variables that determine $U_{\text{CR}}$ have the same values used thus far in this study.

In Fig. 4, we show the evolution of a gas cloud in the virialization shock case at $z = 15$. The heating rate is now evaluated at $D = 500 \text{pc}$, a typical distance from the outer edge of the halo to its densest region. We again examine the evolution for various combinations of $\Psi_\ast$ and $\epsilon_{\text{min}}$. Note that here we consider star formation rates that are an order of magnitude larger than that in the minihalo case. Such increased rates reflect the later formation times of the first dwarf galaxies, so that structure formation has already progressed further. In most of these cases, the presence of CRs has a negligible effect. This is in part due to the enhanced CR attenuation resulting from the larger cloud sizes. Furthermore, unlike in the free-fall minihalo evolution, in the case of strong shocks the gas never reaches very high densities of neutral hydrogen. Densities for the shocked case are instead typically around $\lesssim 0.2 \text{ cm}^{-3}$ early in the halo’s evolution, and $100 \text{ cm}^{-3}$ after the gas temperature has reached the CMB floor. The average heating and ionization rates for this case are thus lower, but the main explanation for the lack of CR-induced differences is that even without CRs the shocked region is highly ionized and able to cool through molecular transitions to the CMB floor. As shown in Johnson & Bromm (2006), HD transitions are able to cool the gas down to the CMB temperature within a Hubble time. Thus, with these cooling mechanisms already in place, CRs can only serve to heat the gas, and the CR heating effect can dominate over molecular cooling only for more extreme star formation rates, as can be seen in the upper right-hand panel of Fig. 4, where the gas temperature does not reach the CMB floor in the presence of strong CR heating.

### 3.3 Local CR feedback

Our study thus far has assumed that the CRs possibly created in the early Universe all become part of a homogeneous and isotropic background. However, if a particular minihalo is within sufficiently close range to a CR-accelerating PISN, the flux of CRs from the nearby PISN may have a greater effect on the evolution of the minihalo than that from the CR background.

The average CR luminosity associated with the PISN is given by

\[ L_{\text{CR}} = 2 \times 10^{38} \text{ erg s}^{-1} \left( \frac{f_{\text{CR}}}{0.1} \right) \left( \frac{E_{\text{SN}}}{10^{52} \text{ erg}} \right) \left( \frac{\Delta_{\text{SN}}}{2 \times 10^3 \text{ yr}} \right)^{-1}, \quad (32) \]

where $\Delta_{\text{SN}}$ is the time over which the PISN emits CRs, here assumed to be approximately the time that passes from the beginning of the PISN to the end of its Sedov–Taylor phase of expansion (see Lagage & Cesarsky 1983). We can then estimate the CR flux, $f_{\text{CR}}$, and energy density, $U_{\text{CR}}$, emitted by the PISN using

\[ f_{\text{CR}} = \frac{L_{\text{CR}}}{4\pi d^2}, \quad (33) \]
by photoionization-heating (e.g. Alvarez et al. 2006; Greif et al. 2007). When considering realistic cases where the gas evolution in minihaloes could have been significantly impacted by CRs, one therefore needs to invoke a universal background that consists of the global contributions to the CR flux, as opposed to the burst-like emission from a single nearby PISN. However, at close distances the local CR feedback may still provide another source of ionization in nearby gas clouds, indirectly leading to increased molecular cooling, thus helping to facilitate collapse and possibly formation of lower mass stars, as discussed below in Section 3.5.

3.4 Dependence on minimum CR energy

One of the crucial uncertainties concerning high-redshift CRs is the minimum energy \( \epsilon_{\min} \) of the CRs that impinge on a primordial cloud. Our default value of \( \epsilon_{\min} = 10^6 \) eV derives from a simple estimate for the lowest possible energy that a CR proton could gain in a SN shock, though other processes may influence this value, possibly increasing or decreasing it. The minimum CR kinetic energy, however, is crucial, because the ionization cross-section varies roughly as \( \epsilon_{\CR} \) for non-relativistic CRs with kinetic energy less than their rest mass energy but greater than \( \sim 10^7 \) eV, the energy corresponding to a momentum of \( \beta_0 \approx 0.01 \). Thus, higher energy CRs will travel farther into the cloud before first ionizing a particle. Only lower energy CRs will release a large portion of their energy into the cloud, and their absence can significantly lower the overall heating and ionization inside the cloud. This is especially apparent when looking at how the thermal evolution of the minihalo (Fig. 2) changes when \( \epsilon_{\min} \) is increased. For typical star formation rates and \( \epsilon_{\min} = 10^6 \) eV, the impact of CRs becomes negligible. It is furthermore interesting to note that for a given \( U_{\CR} \) and low \( \epsilon_{\min} \) values (<\( 10^6 \) eV), using a power law in momentum instead of a power law in energy will give CR heating and ionization rates about an order of magnitude smaller than those found in the original calculations. This is because using a momentum power law is equivalent to using a shallower spectral index at low CR energies, resulting in fewer low-energy CRs, and this again illustrates the importance of having a sufficient number of low-energy CRs to contribute to CR effects. Decreasing \( \epsilon_{\min} \) to arbitrarily low values, on the other hand, will not yield ever increasing CR ionization rates, as the cross-section starts to quickly fall off below \( \sim 10^5 \) eV, and our lowest \( \epsilon_{\min} \) of \( 10^5 \) eV is already sufficiently close to the peak in the ionization cross-section to obtain the maximum CR effects. This was confirmed by actually lowering the value of \( \epsilon_{\min} \) to 1 and 10 keV to see how the heating and ionization rates changed. In both the minihalo and virialization shock cases, these rates were changed by less than a factor of 2. One of the reasons for the uncertainty in \( \epsilon_{\min} \) is that the low-energy part of the CR spectrum is difficult to observe even inside the MW. The interaction of CRs with the solar wind and the resulting deflection by solar magnetic fields prevents any CRs with energies lower than around \( 10^8 \) eV from reaching the Earth, and thus MW CRs below this energy cannot be directly detected. The Galactic CR spectrum is thought to extend to lower energies between \( 10^7 \) and \( 10^9 \) eV (e.g. Webber 1998). However, at non-relativistic energies the CR spectrum in the MW is expected to become much flatter due to ionization energy losses that would be much less relevant in the high-\( z \) IGM, as even at the redshifts we consider the density of the IGM is around three orders of magnitude lower than the average MW density.

3.5 Fragmentation scale

As is evident in Fig. 2, given a strong enough energy density, CRs can serve to lower the minimum temperature that a collapsing cloud
inside a minihalo is able to reach. This could have important implications for the fragmentation scale of such gas clouds. We estimate the possible change in fragmentation scale due to CRs by assuming that the immediate progenitor of a protostar will have a mass approximately given by the Bonnor–Ebert (BE) mass (e.g. Johnson & Bromm 2006):

$$M_{\text{BE}} \simeq 700 \, M_\odot \left( \frac{T_f}{200 \, \text{K}} \right)^{3/2} \left( \frac{n_f}{10^4 \, \text{cm}^{-3}} \right)^{-1/2},$$

where $n_f$ and $T_f$ are the density and temperature of the primordial gas at the point when fragmentation occurs. For each of the cases shown in Fig. 2, the evolution of the BE mass was calculated as the cloud collapsed in free-fall. This evolution is shown in Fig. 6. For the three cases that had the most significant CR effects, the respective cooling and free-fall times were also examined. When the cooling time, $t_{\text{cool}}$, first becomes shorter than the free-fall time, the gas can begin to cool and fall to the centre of the DM halo unimpeded by gas pressure (Rees & Ostriker 1977; White & Rees 1978). As the cloud contracts, however, the density increases and the free-fall time becomes shorter, eventually falling back below the cooling time. At this point, when $t_{\text{cool}} \sim t_{\text{free}}$, we evaluate the BE mass. This is the instance where the gas undergoes a phase of slow, quasi-hydrostatic contraction, what has been termed ‘the loitering phase’ (see Bromm & Larson 2004). Slow contraction continues, and the gas will leave this loitering regime when its mass is high enough to become gravitationally unstable, thus triggering runaway collapse. Thus, the ‘loitering phase’ is the characteristic point in the evolution when fragmentation occurs and the BE mass is relevant, at least to zeroth order, given that star formation is too complex to allow reliable predictions from simple one-zone models such as considered here (e.g. Larson 2003).

Bearing this caveat in mind, we find that, with a sufficiently strong CR flux present, the gas density at which the ‘loitering phase’ occurs is increased by a factor of $\sim 10–100$, and the corresponding temperature decreased by a factor of 2–3. Evaluating equation (37) shows that the BE mass thus decreases by an order of magnitude down to around $10 \, M_\odot$. This is the mass scale of what has been termed ‘Pop II.5’ stars (e.g. Mackey et al. 2003; Johnson & Bromm 2006). This decrease in fragmentation scale occurs for CR ionization rates greater than $\sim 10^{-19} \, \text{s}^{-1}$. A sufficiently large flux of low-energy CRs could thus have facilitated the fragmentation and collapse of primordial minihalo gas into Pop II.5 stars. Again, we emphasize the somewhat speculative nature of our argument regarding stellar mass scales. However, the general trend suggested here, that the presence of CRs in the early Universe tends to enable lower mass star formation, might well survive closer scrutiny with numerical simulations.

### 4 SUMMARY AND DISCUSSION

We have investigated the effect of CRs on the thermal and chemical evolution of primordial gas clouds in two important sites for Pop III star formation: minihaloes and higher mass haloes undergoing strong virialization shocks. We show that the presence of CRs has a negligible impact on the evolution in the latter case, since the primordial gas is able to cool to the CMB floor even without the help of CR ionization, and the direct CR heating is also unimportant unless extremely high star formation rates are assumed. Thus, the thermal and chemical evolution of haloes corresponding to the first dwarf galaxies is rather robust, and the CR emission from the deaths of previously formed Pop III stars will not initiate any feedback in these star formation sites. The impact of CRs on gas in minihaloes, which would typically have formed before the more massive dwarf systems, could have been much more pronounced, given a sufficiently low $\epsilon_{\text{min}}$ and Pop III star formation rates that are not too low. In each case, the effect of direct CR heating is weaker than the indirect molecular cooling that follows from the increased ionization due to CRs. The additional molecular cooling induced by the CRs allows the gas in a minihalo to cool to lower temperatures. For CR ionization rates above a critical value of $\sim 10^{-19} \, \text{s}^{-1}$, we find that the mass scale of metal-free stars might be reduced to $\sim 10 \, M_\odot$, corresponding to what has been termed ‘Pop II.5’.
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