LOWER METAL ENRICHMENT OF VIRIALIZED GAS IN MINIHALOS

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Received 2007 April 23; accepted 2007 October 7

ABSTRACT

We differentiate between the metal enrichment of the gas in virialized minihalos and that of the intergalactic medium at high redshift, which is pertinent to cosmological reionization, with the initial expectation that gas in the high-density regions within formed dark matter halos may be more robust and thus resistant to mixing with the lower density intergalactic medium. Using detailed hydrodynamic simulations of gas clouds in minihalos subject to destructive processes associated with the encompassing intergalactic shocks carrying metal-enriched gas, we find, as an example, that, for realistic shocks with velocities of $10^{\text{0.5}}$ km s$^{-1}$, more than (90\%, 65\%) of the high-density gas with $\rho \geq 500 \rho_b$ inside a minihalo virialized at $z = 10$ with a mass of $(10^7, 10^6) M_\odot$, respectively, remains at a metallicity lower than 3\% of that of the intergalactic medium by redshift $z = 6$. It may be expected that the high-density gas in minihalos will become fuel for subsequent star formation when they are incorporated into larger halos, where efficient atomic cooling can induce gas condensation and hence star formation. Since minihalos virialize at high redshift, when the universe is not expected to have been significantly reionized, the implication is that gas in virialized minihalos may provide an abundant reservoir of primordial gas that could possibly allow the formation of Population III metal-free stars to extend to much lower redshifts than would have been otherwise expected on the basis of the enrichment of the intergalactic medium.

Subject headings: cosmology: theory — early universe — galaxies: formation — intergalactic medium — supernovae: general

1. INTRODUCTION

Recent observations of the high-redshift ($z > 6$) quasar spectra from the Sloan Digital Sky Survey (SDSS; Fan et al. 2001; Becker et al. 2001; Barkana 2002; Cen & McDonald 2002) and the cosmic microwave background fluctuations from the Wilkinson Microwave Anisotropy Probe (WMAP; Spergel et al. 2007) combine to paint a complicated and yet uncertain reionization picture. What is strongly suggested, however, is that the reionization process began at an early time, probably at $z \geq 10$. If stars are primarily responsible for producing the ionizing photons, the overall reionization picture would depend significantly on how and when the transition from the metal-free Population III stars to the metal-poor Population II stars occurs (see, e.g., Omukai 2000; Bromm et al. 2001a; Schneider et al. 2002; Mackey et al. 2003; Schneider et al. 2003; Bromm & Loeb 2003), especially if the Population III star formation history is more extended than is usually thought (Wyithe & Cen 2007). This is based on the finding that Population III stars may be much more massive and hotter than Population II stars (Carr et al. 1984; Larson 1998; Abel et al. 2000; Hernández & Ferrara 2001; Bromm et al. 2001a, 2001b; Nakamura & Umemura 2001; Bromm et al. 2002; Omukai & Palla 2003; Mackey et al. 2003) and hence are efficient producers of photons that ionize hydrogen and helium.

The metallicity of the star-forming gas plays several important roles in the physics of the first stars. First, the transition from Population III to Population II is facilitated by the presence of a small amount of metals; in particular, oxygen and carbon (Bromm & Loeb 2003). Thus, it is the amount of C and O, not necessarily the total amount of "metals," that determines the transition (Fang & Cen 2004). The yield patterns for (nonrotating) stars with masses in the range of 140–260 $M_\odot$ that explode via pair-instability supernovae (PISNe) and regular Type II SNe are different. In the PISN case, the supernova ejecta is enriched by α-elements, whereas the major products of a Type II SN are hydrogen and helium, with a small amount of heavy elements (see, e.g., Woosley & Weaver 1995; Heger & Woosley 2002). Consequently, the transition from Population III stars to Population II stars may occur at different times, depending on the initial mass function (IMF; e.g., Fang & Cen 2004).

Second, while the ionizing photon production efficiency depends only relatively weakly on the exact IMF, as long as the stars are more massive than $\sim 10 M_\odot$ (e.g., Tumlinson et al. 2004), its dependence on metallicity is strong, because the effective temperature of the stellar photosphere depends sensitively on the opacity, and hence metallicity, of the stellar atmosphere. The amount of metals produced depends on the IMF. For example, in the most extreme case in which all Population III stars are more massive than, say, $\geq 270 M_\odot$, these stars may conclude by imploding to form intermediate-mass black holes without ejecting a significant amount of metals to the surroundings. However, exactly how massive Population III stars are is uncertain. While simulations have suggested that Population III stars may be more massive than $100 M_\odot$ ("very massive stars" [VMSs]; Abel et al. 2000; Bromm et al. 2001a), Tan & McKee (2004) find that stellar feedback processes may limit the masses of the Population III stars to $30–100 M_\odot$. Observationally, the VMS picture is advocated by Oh et al. (2001) and Qian & Wasserburg (2002). On the basis of an analysis of metal yield patterns from the PISN explosions of the VMS progenitors (Heger & Woosley 2002), Tumlinson et al. (2004), Daigle et al. (2004), Umeda & Nomoto (2003, 2005), and Venkatesan & Truran (2003), on the other hand, argue that the general pattern in metal-poor halo stars, in the Ly$\alpha$ forest and cosmic star formation history, is more consistent with the yield pattern of Type II SNe, perhaps with a lower cutoff of $10 M_\odot$.

Clearly, the metallicity of the gas out of which stars are formed is critically important. The conventional picture that is often adopted goes as follows: formed stars eject metals into the intergalactic...
medium (IGM) and eventually raise the metallicity of the IGM to above the threshold for the Population III to Population II transition. A somewhat refined version of this picture takes into account that the metallicity enrichment process of the IGM is unlikely to be synchronous for different regions (e.g., Furlanetto & Loeb 2005). Here, we point out a possibly large difference between the metallicity of the IGM and the metallicity of the gas in minihalos. Since minihalos collapse at very high redshifts (e.g., Wyithe & Cen 2007), the large amount of dense gas in minihalos may thus provide a primary fuel for subsequent star formation when eventually the minihalos are incorporated into large systems, where efficient atomic cooling allows the gas to condense and form stars.

To quantify this possible difference between the metallicity of the minihalo gas and that of the IGM, we study the stability and metal enrichment of minihalos subject to metal-rich shock waves launched by supernova explosions from large galaxies. We treat an idealized situation in which a minihalo is subject to shock waves enriched with a chosen metallicity, and we investigate how gas inside may be contaminated by metals. We assume that there has been no star formation, and hence, no self–metal enrichment, in the minihalos, because of the lack of adequate coolants; molecular hydrogen is simply assumed to have been long since destroyed by Lyman-Werner photons produced by earlier stars elsewhere. However, we note that the effect of Lyman-Werner photons is more complicated, and studies have suggested that suppression of star formation in minihalos due to Lyman-Werner photons tends to be translated into a delay in star formation (Johnson et al. 2007). Because the gas in minihalos is significantly overdense compared to the IGM and is bounded by the gravitational potential wells produced by the dark matter halos, the mixing of metals into the gas in minihalos by metal-rich outflows from star-forming galaxies should be expected to be different from that of the IGM. As we will show, the process of the mixing of metal-rich outflows with the gas in minihalos is quite incomplete. Several authors (Murray et al. 1993, hereafter M93; Klein et al. 1994; Dinge 1997; Miniati et al. 1997) have addressed the problem of the stability of a non–self-gravitating gas cloud moving at the sound speed of the background medium, which is equivalent to a shock wave sweeping the gas cloud. They have found that the cloud gets disrupted after a time comparable to the dynamical time of the cloud. Here, we are interested in the self-gravitating case. In particular, we are interested in minihalos that are gravitationally dominated by their dark matter content and that have no cooling. A very similar case was already studied by M93, in the context of a two-phase medium, using two-dimensional simulations. In this work we employ three-dimensional hydrodynamical simulations to study this problem. We simulated halos of masses $10^6$ and $10^7 M_\odot$ that were subject to shock waves with velocities of 10, 30, 100, and 300 km s$^{-1}$. For the slowest cases of 10 and 30 km s$^{-1}$, the halos are quite stable, and the gas inside the virial radius of the halos remains fairly uncontaminated after many dynamical times. Only for the shock velocities of 100 and 300 km s$^{-1}$ do the halos start to become unstable, losing a significant fraction of their gas and getting substantially enriched in their inner regions.

The paper is organized as follows. In § 2 we specify the physical model for the minihalos and shock waves and describe some technical specifications for the code we use. Section 3 presents our results, followed by conclusions in § 4.

2. DESCRIPTION OF THE MODEL

We analyze the metal enrichment of gas in spherical minihalos with total virial masses of $10^6$ and $10^7 M_\odot$, whose virial temperatures are 710 and 3295 K, respectively, at $z = 10$. Initially, the gas in the minihalos is assumed to have zero metallicity. Then the minihalo is exposed to an IGM sweeping through at a velocity of $V_s$ and metallicity $Z_{IGM}$, and we quantify the evolution of the metallicity of the gas inside the minihalo. We study four cases, with $V_s = 10, 30, 100,$ and 300 km s$^{-1}$, for each of the two choices of the minihalo masses.

The gravitational potential of a halo is determined by its dark matter and is assumed not to change. The density of a virialized dark matter halo as a function of radius, $r$, is given by the Navarro-Frenk-White (Navarro et al. 1997) density profile:

$$\rho_{DM}(r) = \frac{\rho_{crit} c}{u(1 + u)^{2}},$$

where $\rho_{crit} = 3H(z)^2/8\pi G$ is the critical density of the universe at redshift $z$, $c = 200^{3/2} m(c)$, and $u = r/r_v$. The characteristic radius $r_v$ is defined in terms of the concentration parameter of the halo, $c$, which is a function of the halo mass and the redshift, and the virial radius, $r_{vir}$. The virial radius is defined in terms of the halo mass, $M_H$, by $(4\pi/3)r_v^3 200\rho_{crit} = M_H$, and the function $m(u) = \ln(1 + u) - u/(1 + u)$. For the concentration parameter, we adopt the fitting formula provided by Dolag et al. (2004):

$$c = \frac{9.59}{1 + z} \left( \frac{M_H}{10^{14} h^{-1} M_\odot} \right)^{-0.102},$$

which is based on computations of a ΛCDM cosmological model with $\Omega_m = 0.3$, $\Omega_L = 0.7$, $\Omega_b = 0.045$, and $\sigma_8 = 0.9$.

Since the gravitational potential, $\phi$, is determined by the dark matter content of the minihalos, it is given by

$$\phi(r) = \begin{cases} -4\pi G \rho_c r^2 \ln(1 + u) \quad & \text{for } u \leq d, \\ -4\pi G \rho_c r^2 \ln \left[ \frac{1 + d}{d} + m(d) \left( \frac{1}{u} - \frac{1}{d} \right) \right] \quad & \text{for } u > d, \end{cases}$$

where $r_d$ is the radius at which the dark matter density of the halo equals the mean density of the universe.

For the gas, we used an $X = 0.76$, $Y = 0.24$, and $Z = 0$ composition. Its density and temperature profiles are determined by assuming that $P_g = \rho_g k_{B} T_g \gamma \rho_b^{\gamma} (m_{p} \mu)^{\gamma}$, where $\gamma$ is the polytropic index. Then we can write

$$\rho_g(r) = \rho_0 (u)^{\gamma-1} \quad \text{and} \quad T_g(r) = T_0 (u)^{\gamma-1},$$

where $\rho_0$ and $T_0$ are the density and temperature at the center of the halo. Since the halos are in hydrodynamic equilibrium, we find that

$$y(u)^{\gamma-1} = \begin{cases} \frac{1}{\gamma} \frac{G M_H m_{p} \mu}{k T c r_m(c)} \left[ \ln(1 + u) - 1 \right] \quad & \text{for } u \leq d, \\ \frac{1}{\gamma} \frac{G M_H m_{p} \mu}{k T c r_m(c)} \left[ \ln(1 + d) - 1 + m(d) \left( \frac{1}{u} - \frac{1}{d} \right) \right] \quad & \text{for } u > d. \end{cases}$$
We have three free parameters in our gas profile: \( C_26 c \), \( T_c \), and \( C_13 \). We choose the central density such that at \( r = r_{\text{vir}} \) the ratio between the dark and baryonic matter densities is equal to \( C_10 m / C_10 b \).

The values for \( T_c \) and \( C_13 \) are determined by considering that, according to many hydrodynamic simulations, in the outer part of the halos, the gas density profile traces the dark matter density profile (Komatsu & Seljak 2001); i.e.,

\[
\frac{d \ln (\rho_{\text{DM}})}{d \ln (\rho_g)} = 1. \tag{6}
\]

Thus, \( T_c \) and \( \gamma \) were determined using the fitting formulae provided by Komatsu & Seljak (2001) that satisfy equation (6) within the range \( c/2 \leq u \leq 2c \):

\[
T_c = \frac{G M_{\odot} m_p \mu}{3 k r_s c} \left[ 0.00676(c - 6.5)^2 + 0.206(c - 6.5) + 2.48 \right],
\]

\[
\gamma = 1.15 + 0.01(c - 6.5). \tag{7}
\]

As mentioned above, the minihalos are shocked by a continuous, hot, metal-rich shock wave that has a temperature of \( T_s = 3 V_s^2 m_p / 16 k_B \), a mean density of \( \rho_s = 4 \rho_b \), and a metallicity of \( Z_{\text{IGM}} \). We also introduce fluctuations in space and time in the density of the shock wave; i.e.,

\[
\rho_s = 4 \rho_b \left[ 1 + A \sin \left( \frac{2 \pi y}{\lambda} + \phi_y \right) \sin \left( \frac{2 \pi z}{\lambda} + \phi_z \right) \sin \left( \frac{2 \pi t}{T} + \phi_t \right) \right], \tag{8}
\]

where \( y \) and \( z \) represent the two spatial coordinates perpendicular to \( x \), which is the direction of propagation of the shock wave; \( t \) is the time; \( A \) and \( \lambda \) represent the amplitude and the length of the fluctuation; and \( \phi_i \) corresponds to an arbitrary phase in the coordinate \( i \). The spatial phases were randomly chosen every time that the value \( t/T \) became an integer, where \( T = \lambda / V_s \). What are appropriate values for \( A \) and \( \lambda \)? At the redshifts of interest \((z \sim 6-10)\), large atomic-cooling halos start to become nonlinear, which
means that the density variance is of order unity on the mass scales of $10^8-10^9 M_\odot$, with corresponding length scales of comoving $\sim 0.1 \text{ Mpc}$. By definition, when a certain mass scale $M$ becomes nonlinear, $A(M) \sim 1$. We have experimented with values of $A = 0.3-0.9$ and $\lambda = 0.003, 0.01, \text{ and } 0.03 \text{ h}^{-1} \text{Mpc}$ in comoving units. Our results turned out to be nearly independent of the values for $A$ and $\lambda$ in the ranges of relevance.

Each simulation starts at $z = 10$, when the IGM shock wave enters the left face of our simulation cube. We do not attempt to vary the background density with time, aside from the variation imposed (see eq. [8] above). We expect that if the background density were allowed to decrease with time, the metal enrichment of the gas in minihalos might be reduced. When there is a need to indicate a redshift during the evolutionary phase of a minihalo, we translate the elapsed time since $z = 10$ to a certain redshift, using the standard cosmological model parameters (Spergel et al. 2007).

We use the total variation diminishing (TVD) hydrodynamics code (Cen et al. 2003) to perform the described simulations. The size of the boxes is chosen such that at the border of the box, the gas density of the halo is equal to the mean baryonic density of the universe. Thus, the comoving size of the boxes is 0.0191 and 0.0457 $h^{-1}$ Mpc for $M_H = 10^6$ and $10^7 M_\odot$, respectively. For most of the simulations, we use $256^3$ cells for each simulation. Our results seem to be convergent to a few percent accuracy, as will be shown at the end of the next section.

3. RESULTS

In this section we analyze our results for the stability and chemical evolution from $z = 10$ to $z = 6$ for the two halos considered. Although for the parameters of the range considered, we observe different levels of instability and mixing, in all cases the gas in the inner region of the halos remains substantially less metalic than the IGM. For $M_H = 10^7 M_\odot$ and $V_s = 10 \text{ km s}^{-1}$, almost all the mass at a density higher than the virial density, $\rho_{\text{vir}} (\approx 49.3 \rho_b$ for $M_H = 10^7 M_\odot$), at $z = 6$ has $Z < 0.03Z_{\text{IGM}}$, whereas for $M_H = 10^6 M_\odot$ and $V_s = 300 \text{ km s}^{-1}$, most of the mass at $\rho \gtrsim \rho_{\text{vir}} (\approx 40.3 \rho_b$ for $M_H = 10^6 M_\odot$) has $Z < 0.3Z_{\text{IGM}}$.

Figures 1 and 2 show the density, metallicity, and velocity of the gas in a slice through the center of the halo that is perpendicular to the shock wave front, at $z = 9$ and $z = 6$. Perhaps the most noticeable result is that the gas cloud inside the minihalo is
able to withstand significant shock waves and can reside inside the halo gravitational potential well for an extended period of time. The gravitational potential well of the dark matter halo is more "steady" than a pure self-gravitating gas cloud, since the dark matter is unaffected by gas dynamical processes at the zeroth order, in agreement with M93.

However, the mixing between the primordial minihalo gas and the metal-enriched IGM due to hydrodynamical instabilities is apparent. First, as seen in the left panels in Figure 2, the Richtmyer-Meshkov instability seems to be most apparent when the interface between the sweeping IGM and the minihalo gas cloud is first being accelerated by the shock moving from left to right. Subsequently, with the buildup of a smoother and larger density transition region on the left side of the halo gas cloud and reduced shock strengths, the Richtmyer-Meshkov instability progressively abates. Second, as can be seen in the right panels in Figures 1 and 2, the Kelvin-Helmholtz instability provides an efficient mechanism for mixing gas in the shearing regions at the outer part of the minihalos. The fact that the density peak largely remains at the center of the dark matter halo over the extended period of time, whereas the outer layers become mixed with the IGM, suggests that mixing due to hydrodynamic instabilities plays the dominant role, whereas ram pressure stripping is subdominant, at least for these two cases and during the displayed time interval. Nevertheless,

![Fig. 3. — Left: Gas masses with different metallicities, shown as a function of time, for the case with a shock velocity of 10 km s\(^{-1}\). Only gas with a density higher than 500\(\rho_b\) is tracked, where \(\rho_b\) is the mean baryon density of the universe at \(z = 10\). The two halo masses considered are \(10^6\) and \(10^7\)\(M_\odot\), each of which is represented by a set of thin and thick lines, respectively. The metallicity ranges considered are \(Z/Z_{\text{IGM}} < \alpha\), where the values \(\alpha = 1, 0.3, 0.1,\) and 0.03 are indicated by solid, dotted, dashed, and dot-dashed lines, respectively. The vertical axis is normalized using the initial mass with density higher than 500\(\rho_b\). The horizontal axis is in units of the dynamical time, \(t_{\text{dyn}} = (800 \pi G \rho_c \sigma^2)^{-1/2}\), of the halo at the virial radius at \(z = 10\). Right: Same as the left panel, except that gas with density higher than the virial density \(\rho_{\text{vir}}\) is followed, where \(\rho_{\text{vir}}\) is the density at the virial radius of the halo (\(\approx 40.3\rho_b\) and 49.3\(\rho_b\) for \(10^6\) and \(10^7\)\(M_\odot\) halos, respectively). In both panels we note that, since all the gas within the two ranges of density starts with \(Z/Z_{\text{IGM}} < 0.03\) (the minimum metallicity considered), all the lines initially overlap, starting with normalized mass \(M = 1\). In some cases, the metal enrichment is not efficient enough and all the gas in a given density range can have \(Z/Z_{\text{IGM}} < 0.03\) all the way until \(z = 6\). For example, in the left panel, no discernible dotted, dashed, and dot-dashed lines can be seen.

![Fig. 4. — Same as Fig. 3, but for the case of a shock velocity of 30 km s\(^{-1}\).]
the central regions of the minihalo gas clouds are significantly contaminated with metals at later times (right panels in Figs. 1 and 2). We will later show in Figure 7 our convergence test of the results, which suggests that our numerical resolution appears to be adequate to properly simulate the hydrodynamic instabilities involved.

We now turn to more quantitative results, focusing on the metal enrichment of the gas inside minihalos that is caused by the IGM shocks. Figures 3–6 show the evolution of the amount of mass at both \( \rho > 500 \rho_b \) and \( \rho > \rho_{vir} \), in units of its corresponding value at the beginning of the simulation at \( z = 10 \), that is metal-enriched to various levels with \( Z < \alpha Z_{IGM} \), with \( \alpha = 1, 0.3, 0.1, \) and 0.03. Figures 3–6 show the cases with \( V_s = 10, 30, 100, \) and 300 km s\(^{-1}\), respectively.

From Figure 3, in which \( V_s = 10 \) km s\(^{-1}\), we see that for \( M_H = 10^6 M_\odot \), only 5% of the gas is contaminated to \( Z \geq 0.03 Z_{IGM} \) for \( \rho \geq \rho_{vir} \), and for a \( 10^7 M_\odot \) halo there is practically no gas with \( Z \geq 0.03 Z_{IGM} \) in that density range after about 11 dynamic times (by \( z \sim 6 \)). For \( \rho > 500 \rho_b \), there is no gas with \( Z \)-values larger than 0.03\( Z_{IGM} \) even for the \( M_H = 10^6 M_\odot \) halo. It is interesting to note that for this velocity, the amount of gas at the two ranges of density considered actually increases instead of decreasing. This is due to the compression produced on the halo by the gas from the shock. We also observe the acoustic oscillations in the amount of gas due to this compression. (The fact that acoustic oscillations for the \( 10^6 M_\odot \) halo start earlier is just due to the smaller simulation box size.) Figure 4 shows the case in which \( V_s = 30 \) km s\(^{-1}\). Here we see that for \( \rho > 500 \rho_b \), again, there is no gas that becomes more metal-rich than \( Z = 0.03 Z_{IGM} \) by \( z \sim 6 \). For \( \rho \geq \rho_{vir} \), only \( \sim 5\% \) of the gas ends up with \( Z \geq 0.03 Z_{IGM} \) for a \( 10^7 M_\odot \) halo. However, for \( M_H = 10^6 M_\odot \), \( \sim 5\% \) of the gas mass reaches \( Z \geq 0.1 Z_{IGM} \).

For \( V_s = 100 \) and 300 km s\(^{-1}\) (Figs. 5 and 6), the stripping of the outer parts of the halo becomes more important, and we start to see that the amount of metal-free gas for the density ranges that are considered starts to decrease significantly, for two reasons.
First, the halo is losing a significant amount of its mass, and therefore, its global structure is being modified. So we observe a decrease in the total amount of mass for \( \rho > 500 \rho_\text{b} \) and \( \rho > \rho_\text{vir} \).

Second, this stripping puts the IGM gas into contact with the innermost part of the halo, moving the mixing layer inward and increasing the efficiency of the mixing in higher density regions in the halo. For \( V_s = 100 \text{ km s}^{-1} \), at \( \rho > 500 \rho_\text{b} \) there is no significant mixing, but just a small overall reduction of the mass. On the other hand, for \( \rho \geq \rho_\text{vir} \), the total decrease of the mass starts to be significant, reaching \( \approx 50\% \) for \( M_H = 10^6 M_\odot \), and the amount of gas that is purer than \( 0.03 Z_{\text{IGM}} \) is only \( \approx 30\% \) and \( \approx 50\% \) of the original counterparts at \( z = 10 \) for \( M_H = 10^6 \) and \( 10^7 M_\odot \), respectively. For \( V_s = 300 \text{ km s}^{-1} \), at \( \rho > 500 \rho_\text{b} \) we observe a significant reduction of the overall mass, especially for \( M_H = 10^7 M_\odot \), where the mass is reduced to \( \approx 30\% \) of its original value. We can see that the mixing itself does not play a very significant role at these densities, with practically no difference between the total mass and the mass of the gas with \( Z < 0.03 Z_{\text{IGM}} \) for \( M_H = 10^7 M_\odot \). The same thing happens with \( M_H = 10^6 M_\odot \), but in this case for \( Z < 0.1 Z_{\text{IGM}} \). For \( \rho > \rho_\text{vir} \) we see that, along with the total reduction of mass, there is substantial enrichment of the gas. The mass of gas with \( Z < 0.03 Z_{\text{IGM}} \) is only \( \approx 3\% \) and \( \approx 25\% \) of its value at \( z = 10 \) for \( M_H = 10^6 \) and \( 10^7 M_\odot \), respectively. To provide a convenient numerical form, we summarize in Tables 1 and 2 the masses of the gas in the different ranges of density and metallicity for the halo masses and shock velocities considered at the relevant redshifts \( z = 7 \) and \( z = 6 \), respectively.

The sensitive dependence of the gas cloud disruption on the shock velocity can be understood in the context of the instability analysis by M93. M93 show that when the parameter \( \eta \), defined as

\[
\eta = \frac{g D R_c}{2 \pi V_s^2},
\]

is above unity, the cloud is stable for up to many dynamical times, where \( D \) is the density ratio of the gas cloud to the background gas, \( R_c \) is the radius of the gas cloud, and \( g \) is the surface gravity. Numerically,

\[
\eta(r) = \frac{1 + z}{11} \left( \frac{V_s}{22 \text{ km s}^{-1}} \right)^{-2} \left( \frac{M_H}{10^6 M_\odot} \right)^{2/3} \left( \frac{M_\odot}{M_H} \right)^{4.7},
\]

where we have assumed that the density slope near the virial radius is \(-2.4 \) (Navarro et al. 1997), \( M_\odot \) is the mass with radius \( r \), and \( z \) is the redshift. Equation (10) suggests that for \( V_s = 25 \text{ km s}^{-1} \), the gas cloud in a minihalo of mass \( M_H = 10^6 M_\odot \) at \( z \approx 10 \) is generally quite stable, in agreement with our results. For \( V_s = 300 \text{ km s}^{-1} \) and \( M_H = 10^7 M_\odot \), one obtains \( \eta \approx 0.005 \) at \( z = 10 \) and \( M_\odot = M_H \), which suggests that the outskirts of the minihalo gas cloud would be disrupted on the order of a dynamic time, which is also consistent with our results. For \( V_s = 100 \text{ km s}^{-1} \) and \( M_H = 10^7 M_\odot \), we find \( \eta \approx 0.03 \) at \( z = 10 \) and \( M_\odot = M_H \); M93 find that at \( \eta = 0.25 \), the gas mass loss is still relatively small over many dynamical times, again consistent with our simulations.

Mori et al. (2002) show, in simulations of the propagation of supernova blast waves from \( 10^5 h^{-1} M_\odot \) galaxies at \( z = 9 \), that after more than a hundred million years the relative filling factor for regions being swept by shocks of velocities larger than \( U = 10, 30, \) and \( 100 \) km s\(^{-1}\) is roughly \( 100\% \), \( 35\% \), and \( 10\% \), respectively. We expect the velocities to be still smaller at the higher redshifts of concern here, due to enhanced cooling and a larger Hubble velocity. Therefore, in a real cosmological setting, in combination with our findings, we expect that a large fraction of the gas already virialized in minihalos will be largely unaffected by metal-carrying blast waves and will remain metal-free to modern redshifts, possibly as low as \( z = 5 \) -- \( 6 \), when the gas in minihalos may be photoevaporated globally.

It is prudent to check the convergence of computed results. We have performed additional simulations with \( 128^3 \) and \( 512^3 \) grid points. We show in Figure 7 an example of the convergence tests that we have done. We see that while the difference between the \( 128^3 \) and \( 256^3 \) cases can amount to tens of percent at late times (say, \( t/t_{\text{dyn}} > 5 \)), the difference between the \( 256^3 \) and \( 512^3 \) cases.
is dramatically reduced and is at the level of a few percent even at very late times \((t/t_{dyn}>10)\). It is instructive to note that the tendency is to decrease the level of mixing as we increase the resolution. Thus, our results must be interpreted as an upper limit to the metal enrichment of minihalos by shock waves, with an accuracy of a few percent.

4. CONCLUSIONS

It is frequently assumed that the metallicity of the intergalactic medium is the primary determinant of the epoch of the transition from Population III stars to Population II stars. We wish to point out a potentially large difference between the metallicity of the intergalactic medium and the metallicity of the gas in minihalos. Utilizing hydrodynamic simulations of gas clouds in minihalos subject to destructive processes associated with the encompassing intergalactic shocks carrying metal-enriched gas, we find that a large fraction of the gas in virialized minihalos remains at a metallicity that is much lower than that of the intergalactic medium. For example, for realistic shocks with velocities of 10–100 km s\(^{-1}\), more than (90\%, 65\%) of the high-density gas with \(\rho > 500\rho_b\) inside a minihalo virialized at \(z = 10\) of mass \((10^7, 10^8) M_\odot\) remains at a metallicity lower than 3% of that of the intergalactic medium by redshift \(z = 6\), under the harsh condition that the minihalo is exposed to shock waves continuously from \(z = 10\) to \(z = 6\).

In the standard cosmological model, if large halos with efficient atomic cooling are responsible for producing most of the reionizing photons, smaller minihalos will virialize before the universe is significantly reionized. Thus, gas in virialized minihalos may provide an abundant reservoir of primordial gas to possibly allow the formation of Population III metal-free stars to extend to much lower redshifts than otherwise expected on the basis of the enrichment of the intergalactic medium.

A related issue that is not addressed here concerns the fate of the gas inside minihalos when it is exposed to reionizing photons. The situation is complicated because the timescale of the photoevaporation of gas in minihalos (Barkana & Loeb 2002; Iliev et al. 2005; Ciardi et al. 2006) may be \(100\) Myr (Shapiro et al. 2004); the timescale may still be longer at higher redshifts \((z > 10)\) and/or at lower ionizing fluxes than those used in the work of Shapiro et al. (2004). It may be that full understanding will require detailed calculations that incorporate both radiative transfer and metal-enrichment processes.

We gratefully acknowledge financial support by NSF grant AST0407176, NASA grant NNG06GI09G, and a grant from Princeton University.

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