DESTRUCTION OF MOLECULAR HYDROGEN DURING COSMOLOGICAL REIONIZATION

ZOLTÁN HAIMAN,1 MARTIN J. REES,2 AND ABRAHAM LOEB1

Received 1996 May 23; accepted 1996 September 6

ABSTRACT

We investigate the ability of primordial gas clouds to retain molecular hydrogen (H₂) during the initial phase of the reionization epoch. We find that before the Strömgren spheres of the individual ionizing sources overlap, the UV background below the ionization threshold is able to penetrate large clouds and suppress their H₂ abundance. The consequence lack of H₂ cooling could prevent the collapse and fragmentation of clouds with virial temperatures \( T_{\text{vir}} \lesssim 10^4 \) K (or masses \( \lesssim 10^6 M_\odot [(1 + z_{\text{vir}})/10]^{-3/2} \)). This negative feedback on structure formation arises from the very first ionizing sources and precedes the feedback due to the photoionization heating.

Subject headings: cosmology: theory — early universe — galaxies: formation — molecular processes — radiative transfer

1. INTRODUCTION

Lately, there has been a renewed interest in the reionization of the intergalactic medium (IGM) due to its signature on the cosmic microwave background anisotropies (Tegmark, Silk, & Blanchard 1994; Aghanim et al. 1996; Loeb 1996). Observations of quasi-stellar object (QSO) spectra imply that the IGM has been highly ionized already by the redshift \( z \approx 5 \) (Schneider, Schmidt, & Gunn 1991). Evidence for the presence of an ionizing background radiation with the required magnitude is provided by the proximity effect (Bajtlik, Duncan, & Ostriker 1988; Lu, Wolfe, & Turnshek 1991; Bechtold 1994; and references therein). The ionizing sources could have either been stars or quasars (Miralda-Escudé & Ostriker 1990; Cen & Ostriker 1993; Meiksin & Madau 1993; Shapiro, Giroux, & Babul 1994; Fukugita & Kawasaki 1994).

In cold dark matter (CDM) cosmologies, the first generation of objects forms at redshifts 10–50. Fragmentation of these objects into stars would result in early reionization. Couchman & Rees (1986) have argued that the heating associated with the photoionization of the IGM would inevitably raise the gas temperature to \( T \sim 10^4 \) K and increase the value of the cosmological Jeans mass well above its initial value of \( \sim 10^5 M_\odot \). This effect would tend to suppress further baryonic structure formation until the nonlinear mass scale rises above the new Jeans mass. This negative feedback serves as a buffer that monitors the fraction of baryons converted into massive stars and the metallicity of the IGM as a result of early reionization (Fukugita & Kawasaki 1994), and it could in principle explain the nearly universal metallicity value of \( \sim 1\% \) observed in Ly\( \alpha \) absorption systems (Cowie et al. 1995; Tytler et al. 1995).

The ability of primordial gas clouds with virial temperatures \( T_{\text{vir}} \approx 10^2–10^3 \) K to cool radiatively and fragment after they virialize depends strongly on the existence of molecular hydrogen in them (Haiman, Thoul, & Loeb 1996, hereafter HTL96; Tegmark et al. 1996). If the existence of H₂ is required in order to produce ionizing sources such as stars or quasars, then it is important to consider the negative feedback that a UV background has on the H₂ abundance in the universe. As hydrogen molecules can be destroyed by photons below the Lyman limit (Stecker & Williams 1967) to which the universe is transparent before being ionized, it is possible that the feedback due to H₂ destruction preceded the feedback due to photoionization heating.

In the early phase of reionization, after the first few sources of ionizing radiation appear, the Strömgren spheres of the individual sources are well separated from each other. During this phase, the ionizing photons from each source are absorbed at the boundary of their local Strömgren spheres, leaving the bulk of the universe neutral. However, the photons below the Lyman limit escape the Strömgren spheres and are able to dissociate H₂ molecules elsewhere. Therefore ionizing sources can affect the H₂ abundance at a large distance (Dekel & Rees 1987).

In a previous paper (Haiman, Rees, & Loeb 1996; hereafter HRL96), we studied the effect of quasar-like ionizing sources (with a power-law spectrum) on the H₂ cooling of gas clouds after the universe is reionized. We have found that in very dense clouds \( n \gtrsim 10^3 \text{ cm}^{-3} \) the ionizing radiation enhances the formation of molecules and results in a net cooling of the gas. Here we examine the effect of the background radiation on the H₂ molecules at lower densities, inside clouds that are collapsing when most of the universe is still neutral. In contrast with HRL96, we calculate the spectrum of the background radiation before the reionization epoch. The processed photon spectrum above the Lyman limit is suppressed strongly due to H ± continuum absorption throughout the universe. Below the Lyman limit, the spectrum is modulated due to H ± line absorption. By estimating the required sizes for clouds to self-shield and retain their H₂ molecules, we demonstrate that the H₂ molecules inside virialized clouds are destroyed easily before the universe becomes reionized. The photodissociation is due to photons softer than the Lyman limit; our results, therefore, are insensitive to whether the first objects emit a thermal...
MOLECULAR HYDROGEN DESTRUCTION 459

(O-B star) or nonthermal spectrum. The early destruction of
H₂ may provide a negative feedback that suppresses col-
lapse and fragmentation prior to the feedback due to photo-
ionization heating of neutral hydrogen. These two feedback
scenarios could differ appreciably in the redshift at which
the IGM became ionized, the predicted number of Popu-
lation III stars at the end of the reionization epoch, and
the scales of the aggregates in which they, and hence the first
nucleosynthesis products, are formed.

The paper is organized as follows. In §2 we calculate the
characteristic spectrum of the ionizing radiation during the
early phase of the reionization, when the Strömgren spheres
around individual sources do not overlap. In §3 we use this
spectrum to examine the question of whether relatively
weak fluxes of photons can penetrate large clouds and
destroy in them. Finally, §4 summarizes the cosmo-
logical implications of the feedback due to the destruction
of H₂ molecules.

2. SPECTRUM OF UV BACKGROUND DURING THE EARLY
PHASE OF REIONIZATION

We consider the first stage of the reionization epoch,
when most of the universe is still neutral and the Strömgren
spheres around individual sources do not overlap. In order
to calculate the spectrum of the UV background outside the
Strömgren spheres, we make four assumptions: (i) the ion-
zizing sources turn on simultaneously at the redshift z_{on}; (ii)
the sources are distributed uniformly in space; (iii) the
emitted spectral intensity is a power law with an index α;
and (iv) the source population maintains a constant co-
moving density. Assumption (iii) would be appropriate if
the ionizing sources were mainly accreting black holes; our
conclusions, however, depend primarily on the emission of
photons in the 11.18–13.6 eV range and therefore apply
with little change if the radiation comes from O and B stars.

The spectrum of the background radiation is suppressed
above the Lyman limit through absorption by the column
of neutral hydrogen and helium between the source and the
observer. Photons below the Lyman limit will freely escape
the Strömgren spheres around their sources. Since H₂ is
dissociated by line absorption into the Lyman and Werner
lines of molecular hydrogen between 11.18–14.67 eV, special
care is needed in calculating the spectrum in this range.

The cumulative flux (in units of ergs cm⁻² s⁻¹ sr⁻¹ Hz⁻¹) observed at the frequency ν and redshift z_{obs} due to
sources with a comoving emission coefficient j(ν, z) (in units of ergs cm⁻³ s⁻¹ sr⁻¹ Hz⁻¹) between the redshifts z_{obs}
and z_{on} can be written as

\[
J(ν, z_{obs}) = \frac{c}{H_0} \int_{z_{obs}}^{z_{on}} j(ν, z) \exp \left[ -τ(z) \right] \frac{dz}{(1+z)^2(1+q_0 z)^{1/2}},
\]  

where \(ν_1 = ν(1+z)/(1+z_{obs})\), and \(τ(z)\) is the optical depth due to neutral hydrogen and helium between the observer’s
redshift z_{obs} and the redshift z,

\[
τ(z) = \frac{c}{H_0} \int_{z_{obs}}^{z} \frac{κ(ν, z) dz}{(1+z)^2(1+q_0 z)^{1/2}},
\]

where

\[
κ(ν, z) = n_H(z)σ_{HH}(ν) + n_H(z)σ_{He}(ν).
\]

Above the hydrogen ionization threshold, equation (1)
must be evaluated numerically. Below this threshold, the
flux is given by the integral in equation (1), with \(τ = 0\) if we
ignore absorption due to lines. Assuming \(q_0 = \frac{1}{2}\) and
writing the emission coefficient as \(j(ν, z) = j_0(ν/ν_H)^{-α}\), where \(ν_H\) is the hydrogen ionization energy, we obtain in this case

\[
J(ν, z_{obs}) = \frac{c j_0}{H_0 ν_H} \left( \frac{ν}{ν_H} \right)^{-α} \left( 1 + z_{obs} \right)^{1.5+α} \times \left[ \frac{1}{1 + z_{obs}^{1.5+α}} - \frac{1}{1 + z_{on}^{1.5+α}} \right].
\]

In reality, the spectrum below the ionization threshold is
processed due to line absorption in the Lyman bands of
neutral hydrogen. When the IGM is still neutral, the optical
depth across the universe in the resonant Lyman lines is
\(\sim 10^6\), and so each photon becomes absorbed at the redshift
at which its frequency equals one of the Lyman lines. Hence,
photons observed to have a frequency ν at a redshift z_{obs}
could have been emitted only from redshifts that are
between the observer and an effective screen located some
distance away. The redshift of the screen z_{screen} depends on
ν_{L₁Ly}i, the frequency of the nearest Lyman line above ν,
through the resonance condition

\[
ν = \frac{(1 + z_{screen})}{(1 + z_{obs})} = ν_{L₁Ly}.
\]

We can incorporate the effect of line absorption by replacing
the turn-on redshift z_{on} with the screen redshift z_{screen} in
equation (4). This is equivalent to multiplying the flux in
equation (4) by the modulating factor

\[
f = 1 - \left[ (1 + z_{obs})/(1 + z_{screen}) \right]^{1.5+α}.
\]

So far we have ignored the radiative decay of the excited
hydrogen atoms. In reality, each absorbing screen will
reemit line photons. Lyx photons are just recycled until they
redshift out of resonance, because their destruction rate by
the 2s \(\rightarrow 1s\) two-photon decay is much slower than the
cosmological expansion (see eq. [6114] in Peebles 1993).
Therefore, there is no net absorption at the redshift corre-
sponding to the Lyx resonance. On the other hand, Lyβ and
higher Lyman line photons are reemitted as a sum of Lyx
and other Balmer or lower line photons. Consequently,
the screens corresponding to Lyβ and higher Lyman lines
produce two effects: (i) they block the view of all sources
beyond the screen redshift, and (ii) they cascade to the
ground state by emitting Lyx and some other, lower fre-
cquency, photons. This results in a sawtooth modulation of
the spectrum in the frequency range from Lyβ up to the
Lyman limit and in a discontinuous step in the spectrum
below the Lyx frequency.

Although it is outside the range of the Lyman and
Werner bands of molecular hydrogen, the flux just below
the Lyx frequency effects the dissociation rate of H₂. To calculate this flux, we assume that each radiative
excitation to the 3p state or higher is followed instantane-
ously by a decay to the ground state with the emission of
exactly one Lyx photon (plus other lower energy photons
that are not important for our analysis). In addition, we
assume that the reemitted Lyx photons shift immediately
out of resonance due to cosmological expansion and
proceed to the observer without scattering further. Under
these assumptions, the additional flux at the frequency ν
below Ly\(\alpha\) due to Ly\(\alpha\) emission from all screens can be written as a sum of terms, each of which is similar to equation (4),

\[
\Delta J(\nu, z_{\text{obs}}) = \sum_{i = \beta, \gamma, \ldots} \frac{c f_i}{H_0} \left( \frac{\nu}{\nu_\beta} \right)^{-z} \left( \frac{v_{\text{Ly}\alpha}}{v_{\text{Ly}i}} \right)^{1-a} \left( 1 + z_{\text{obs}} \right)^{1.5 + z} 
\times \frac{1}{(1 + z_i)^{1.5 + z}} - \frac{1}{(1 + z_{\text{obs}})^{1.5 + z}},
\]

where \(z_a\) is the redshift at which the absorbing screens need to be located in order that their Ly\(\alpha\) emission would appear at a frequency \(\nu\) at the redshift \(z_{\text{obs}}\),

\[
z_a = \frac{v_{\text{Ly}\alpha}}{v} (1 + z_{\text{obs}}) - 1,
\]

and \(z_i\) is the highest redshift at which the Ly\(i\) photons absorbed at \(z_a\) could have been emitted, i.e., the location of a Ly\((i+1)\)-absorbing screen,

\[
z_i = \frac{v_{\text{Ly}(i+1)}}{v_{\text{Ly}i}} (1 + z_{\text{obs}}) - 1.
\]

The factor \((v_{\text{Ly}/v_{\text{Ly}i}})^{1-a}\) in equation (7) accounts for the conversion of a Ly\(i\) to a Ly\(\alpha\) photon during absorption and reemission.

Figure 1 shows, as an example, the processed spectrum at \(z_{\text{obs}} = 25\) assuming that the sources turn on at \(z_{\text{on}} = 35\). The top panel shows that above the Lyman limit the flux is suppressed up to the keV energy range, and the bottom panel presents the sawtooth modulation below the Lyman limit, as well as the discontinuous jump at the Ly\(\alpha\) frequency.

The assumption that the UV source distribution is homogeneous remains valid as long as the radius around the observer out to the nearest observing screen is larger than the clustering length of the sources. Clearly the assumption breaks down for the highest transitions in the Lyman series. However, our calculation is concerned with the molecular Lyman or Werner lines that are separated from the nearest Lyman line by redshift intervals that are typically \([\Delta z/(1 + z)] \approx 10^{-2}\), i.e., much larger than the expected clustering lengths at high redshifts.

One may wonder whether it is possible to detect the reionization spectrum directly. The sawtooth template presented in Figure 1 has a characteristic structure that should survive until the present time, provided the breakthrough epoch during which the universe becomes transparent to lines is sufficiently sudden. Following an abrupt breakthrough epoch (as expected for an unclustered source population), the spectrum below the Lyman limit would be preserved and would show up as a slight modulation on top of the contribution from more recent sources that dominate the observed infrared background. The frequencies at which this sawtooth modulation appears depends on the reionization redshift, and its detection would provide valuable information about the reionization epoch. Unfortunately, the amplitude of the sawtooth modulation is small. We estimate this amplitude based on the DIRBE observations (see Bond 1995 for a review) and assuming a constant comoving density of ionizing sources that ionize the universe abruptly at \(z = 5\) and produce \(J_{21} = 0.1\) at \(z = 2\). For these conditions, the fractional modulation amplitude is smaller than \(10^{-6}\), implying that direct detection of the reionization signal is not feasible.

3. DESTRUCTION OF MOLECULAR HYDROGEN

We may now use the spectrum illustrated in Figure 1 to derive the H\(_2\) fraction inside overdense regions in the universe. In particular, we would like to find out how massive should a gas cloud be in order to form sufficient H\(_2\) to cool on its dynamical timescale.

In calculating the H\(_2\) fraction, we consider a homogeneous cloud of primordial H/He ratio with a density \(\rho\) and temperature \(T\), which is composed of the nine species H, H\(^+\), H\(^-\), He, He\(^+\), He\(^++\), H\(_2\), H\(_2^+\), and \(e^-\). The chemical composition of the gas within an overdense region is determined by the interaction between these species and the background radiation. For the present purpose, we assume chemical and ionization equilibrium, which is reasonable for the densities and temperatures of interest. The equilibrium H\(_2\) fraction is expected to provide an upper limit on the actual H\(_2\) fraction, and so our constraints on the destruction of H\(_2\) are conservative. The assumption of equilibrium is powerful in that it makes our conclusions independent of the details of hydrodynamic and chemical history of the gas.

As emphasized before, the most important element of the chemical network in the present context is the dissociation of H\(_2\) by photons with energies below 13.6 eV. The significance of this two-step process was highlighted first by Solomon in 1965 (see Field, Somerville, & Dressler 1966) and subsequently studied quantitatively by Stecher & Williams (1967). Here we calculate the dissociation rate for this process by summing the oscillator strengths \(f_i\) for the Lyman and Werner bands of H\(_2\) from Allison & Dalgarno (1970) and multiplying each \(f_i\) by the dissociation fractions (for the second step in the two-step process) from Dalgarno

\[
\begin{array}{|c|c|c|c|c|}
\hline
\text{Ly}\alpha & \text{Ly}\beta & \text{Ly}\gamma & \text{Ly}\delta & \text{Ly}\epsilon \\
\text{H}\beta & \text{H}\gamma & \text{H}\delta & \text{H}\epsilon & \text{H}\zeta \\
\hline
\end{array}
\]
& Stephens (1970), and by the flux at each molecular line. HTL96 have shown that the collapse of a primordial gas cloud beyond virialization depends on its ability to form H\textsubscript{2} and cool radiatively. The size of a cloud that is able to form H\textsubscript{2} depends, in turn, on its self-shielding in the Lyman and Werner lines. In order to correct the dissociation rate for self-shielding as a function of H\textsubscript{2} column density, we used the self-shielding function from de Jong, Dalgarno, & Boland (1980). Other details of the chemistry, including the full list of reactions and rate coefficients that we used, are described in HRL96.

For the power-law spectrum in equations (4) and (7), we choose a power-law index \( \alpha = 0.7 \), and an upper cutoff of \( E = 40 \text{ keV} \) for the background flux, as expected from quasar spectra and fits to the X-ray background (Fabian & Barcons 1992). We leave the normalization of the flux as a free parameter and specify it as \( J_{21} \), the flux (without absorption) at the Lyman limit in units of \( 10^{-21} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ Hz}^{-1} \text{ sr}^{-1} \).

We calculate the equilibrium fraction of H\textsubscript{2} within a UV-irradiated homogenous gas cloud by integrating the chemical rate equations. In order to construct the equilibrium H\textsubscript{2} profile as a function of radius, we subdivide the homogenous cloud into a large number of plane-parallel slabs, and in each slab we integrate the rate equations until equilibrium is reached, taking account of the cumulative optical depth due to the H\textsubscript{2}, He, and H\textsubscript{2} lines. We consider a top-hat overdensity of 200 relative to the background universe and five different combinations of redshift and temperature. As a plausible example, first we set \( z = 20 \) and \( T = 10^4 \text{ K} \). Then we repeat the calculation for \( z = 10 \) and 50 keeping \( T = 10^3 \text{ K} \), and for \( T = 10^{2.2} \text{ and } 10^{3.5} \text{ K} \) keeping \( z = 20 \). We selected these temperatures to bracket the range favorable for H\textsubscript{2} cooling; H\textsubscript{2} is not formed below \( 10^{2.2} \text{ K} \), while it is dissociated collisionally above \( \sim 10^{3.5} \text{ K} \) (see HRL96). The background universe is assumed to have \( Q_0 = 0.1 \) and \( H_0 = 50 \text{ km s}^{-1} \text{ Mpc}^{-1} \). As both the temperature and the density of the gas are fixed, the only free parameter left in this calculation is the normalization of the flux, \( J_{21} \).

An example for a typical H\textsubscript{2} fraction profile in a cloud is shown in Figure 2. This figure shows the variation of the equilibrium H\textsubscript{2} fraction \( f_{\text{H}_2} \) as a fraction of neutral hydrogen column density \( n_{\text{H}_2} \) for a hydrogen density \( n_{\text{H}} = 1 \text{ cm}^{-3} \), a temperature \( T = 10^4 \text{ K} \), and a flux \( J_{21} = 0.1 \). Based on the processes that determine the H\textsubscript{2} fraction, three characteristic regimes may be distinguished. Up to hydrogen column densities of \( N_{\text{H}_2} \sim 10^{19} \text{ cm}^{-2} \), the most important mechanism for the destruction of H\textsubscript{2} is direct photodissociation by photons with energies above 15.4 eV. At larger \( N_{\text{H}_2} \), the destruction rate is dominated by the Solomon process, and \( f_{\text{H}_2} \) drops. At still larger depths, \( N_{\text{H}_2} \sim 10^{23} \text{ cm}^{-2} \), the Lyman and Werner lines of H\textsubscript{2} become self-shielding, and collisional destruction by fast electrons (supplied mainly by the ionization of helium) dominates. The equilibrium state deep inside the cloud in which \( N_{\text{H}_2} \sim 10^{24} \text{ cm}^{-2} \) is fully molecular, but this state is not likely to be reached within a Hubble time (see discussion in HRL96).

The effect of the sawtooth modulation on the H\textsubscript{2} dissociation rate can be fairly large. In the limit \( z_{\text{obs}}/z_{\text{vir}} \to 0 \), the modulation decreases the total dissociation rate by a factor of 32.5. This is the maximum effect the sawtooth modulation can introduce, and as \( z_{\text{obs}} \) approaches \( z_{\text{vir}} \), the decrease in the rate is smaller (see eq. [6]).
confirms the existence of a background flux of at least $J_{21} \approx 0.1$ that keeps the IGM fully ionized at $z \approx 5$. A similar or even higher value of $J_{21}$ is expected at the time of breakthrough during the reionization epoch at higher redshifts, since the density of the IGM is then higher. Figures 3a and 3b show that even at fluxes that would be able to ionize the IGM only partially, $J_{21} \lesssim 0.01$, the mass required to shield $H_2$ against the radiation is large, $M_{H_2} \gtrsim 10^9 M_\odot$. This conclusion can be inferred from both Figures 3a and 3b, implying that it is, in fact, insensitive to the details of the sawtooth modulation of the spectrum.

4. COSMOLOGICAL IMPLICATIONS

Our main result is that the mass of a virialized primordial cloud has to be exceedingly high in order to self-shield against even a low level of background radiation. Figures 3a and 3b show that even at fluxes that would be able to ionize the IGM only partially, $J_{21} \lesssim 0.01$, the mass required to shield $H_2$ against the radiation is large, $M_{H_2} \gtrsim 10^9 M_\odot$. This conclusion can be inferred from both Figures 3a and 3b, implying that it is, in fact, insensitive to the details of the sawtooth modulation of the spectrum.

The cooling masses we obtained depend primarily on the number of photons in the narrow spectral range of 11.18–13.6 eV. Our conclusions are therefore insensitive to the detailed shape of the spectrum and should hold for either stellar (blackbody) or quasar (power-law) spectra. We note also that in our calculation of the cooling mass, we assumed a homogeneous spherical cloud. Although centrally condensed clouds are, in general, more effective than homogeneous clouds in self-shielding their molecules, a steepening of the density profile is unlikely to decrease the value of the required cooling mass by orders of magnitude.

We have presented our results for a range of temperatures and redshifts. Our results are relevant to any model in which the initial fluctuations have amplitudes decreasing with scale, so that cosmic structures form hierarchically. Such models predict different turn-on redshifts for the first sources of radiation; in Peebles isocurvature baryon (PIB) models, turn-on may occur at $z \approx 100$; in CDM it may occur in the range $z \approx 10$–$50$; for mixed dark matter models the first structures may form still more recently. Molecular cooling is more efficient at high densities and therefore at larger redshifts; however, in all these models it determines the scale of the first objects that condense out and contribute the first injection of heat into the universe. The amount of UV background generated by each such object is very uncertain; it depends on the efficiency of star formation and whether initial mass function (IMF) favors massive stars, or even supermassive objects or black holes. However, we have shown that $H_2$ cooling is quenched even in massive objects before the UV background has reached the level needed to photoionize the entire IGM. Therefore, the radiation of the first sources exerts an important feedback on the cosmogonic process: the only clouds able to fragment into new stars are those with virial temperature above $\approx 10^4 K$, i.e., with masses $\lesssim 10^8 M_\odot [(1 + z_{vir})/10]^{-3/2}$, so that isothermal contrac-
tion is possible due to atomic line cooling. Clouds that virialize at lower temperatures will, in the absence of a sufficient amount of H$_2$, maintain virial equilibrium and not fragment into stars.

We therefore draw the robust conclusion that the IGM remained predominantly neutral until a sufficient number of objects with virial temperatures above $10^4$ K (masses above $D_10^8$ had gone nonlinear. Most of the O-B stars or accreting black holes that photoionized the IGM had to form in systems at least as large as this. Formation of such systems would have continued unimpeded until ionization was complete; subsequently $J_{21}$ rises sharply to a value of order unity, when the universe becomes, in effect, a single H II region. The only net cooling of a fully photoionized gas comes from bremsstrahlung, which is less effective than the collisionally excited line emission from a partially ionized gas (Efstathiou 1992; Weinberg, Hernquist, & Katz 1997; Navarro & Steinmetz 1996). The completion of photoionization may therefore signal another pause in the cosmogenic process, associated with the further decline in the efficiency of cooling, and the consequent further increase in the minimum scale that can collapse.

We thank Alex Dalgarno for discussions on the chemistry, and the referee, Alain Blanchard, for useful comments. A. L. acknowledges support from the NASA ATP grant NAG 5-3085.

REFERENCES

Aghanim, N., Désert, F. X., Puget, J. L., & Gispert, R. 1996, A&A, 311, 1
Allison, A. C., & Dalgarno, A. 1970, At. Data, 1, 289
Bajtlik, S., Duncan, R. C., & Ostriker, J. P. 1988, ApJ, 327, 570
Bardeen, J. M., Bond, J. R., Kaiser, N., & Szalay, A. S. 1986, ApJ, 304, 15
(BBKSW)
Bechtold, J. 1994, ApJS, 91, 1
Bond, J. R. 1995, in Cosmology and Large Scale Structure, ed. R. Schaeffer (Amsterdam: Elsevier), 1
Cen, R., & Ostriker, J. P. 1993, ApJ, 414, 507
Couchman, H. M. P., & Rees, M. J. 1986, MNRAS, 221, 53
Cowie, L. L., Songaila, A., Kim, T.-S., & Hu, E. M. 1995, AJ, 109, 1522
Dalgarno, A., & Stephens, T. L. 1970, ApJ, 160, L107
de Jong, T., Dalgarno, A., & Boland, W. 1980, A&A, 91, 68
Dekel, A., & Rees, M. J. 1987, Nature, 326, 455
Efstathiou, G. P. 1992, MNRAS, 256, L43
Fabian, A. C., & Barcons, X. 1992, ARA&A, 30, 429
Field, G. B., Somerville, W. B., & Dressler, K. 1966, ARA&A, 4, 207
Fukugita, M., & Kawasaki, M. 1994, MNRAS, 269, 563
Haiman, Z., Rees, M. J., & Loeb, A. 1996, ApJ, 467, 522 (HRL96)
Haiman, Z., Thoul, A., & Loeb, A. 1996, ApJ, 464, 523 (HTL96)
Lepp, S., & Shull, J. M. 1984, ApJ, 280, 465
Loeb, A. 1996, ApJ, 459, L5
Lu, L., Wolfe, A. M., & Turnshek, D. A. 1991, ApJ, 367, 19
Madau, P., & Shull, J. M. 1996, ApJ, 457, 551
Meiksin, A., & Madau, P. 1993, ApJ, 412, 34
Miralda-Escudé, J., & Ostriker, J. P. 1990, ApJ, 350, 1
Navarro, J., & Steinmetz, M. 1996, ApJ, submitted, preprint astro-ph/9605043
Peebles, P. J. E. 1993, Principles of Physical Cosmology (Princeton: Princeton Univ. Press)
Schneider, D. P., Schmidt, M., & Gunn, J. E. 1991, AJ, 102, 837
Shapiro, P. R., Giroux, M. L., & Babul, A. 1994, ApJ, 427, 25
Stecker, T. P., & Williams, D. A. 1967, ApJ, 149, L29
Tegmark, M., Silk, J., & Blanchard, A. 1994, ApJ, 420, 484
Tegmark, M., Silk, J., Rees, M. J., Abel, T., & Blanchard, A. 1996, ApJ, in press, preprint astro-ph/9603007
Tytler, D., et al. 1995, in QSO Absorption Lines, ed. G. Meylan (Heidelberg: Springer), 289
Weinberg, D. H., Hernquist, L., & Katz, N. 1997, ApJ, in press, preprint astro-ph/9604175
ERRATUM

In the paper “Destruction of Molecular Hydrogen during Cosmological Reionization” by Zoltán Haiman, Martin J. Rees, and Abraham Loeb (ApJ, 476, 458 [1997]), Figure 3 is incorrect. Because of a typographical error in the chemistry routine, the ionization profiles and the derived shielding masses were calculated for densities lower than the stated values. As a result, the shielding masses shown in Figures 3a and 3b are artificially high. Figure 1 below shows the corrected version of Figure 3a; the extent of the correction to Figure 3b is similar. Although the correction substantially reduces the shielding masses, the original claim in the paper—that H$_2$ is universally destroyed in objects with virial temperatures below $T_{\text{vir}} = 10^4$ K long before the universe is reionized—remains valid. The baryonic mass corresponding to $10^4$ K in the spherical top-hat collapse model is $M_\ast = 10^7 M_\odot (\Omega_b/0.1)(h/0.5)^{-1}[(1+z_{\text{vir}})/10]^{-3/2}$, while the value of $J_{21}$ necessary to ionize the universe (i.e., the radiation density corresponding to one photon per baryon) is $J_{21}^\ast = 80(\Omega_b h^2/0.025)[(1 + z)/25]^3$. In all cases shown in the corrected figure, the shielding masses at $J_{21}^\ast$ still exceed the value of $M_\ast$ by several orders of magnitude.

The authors would like to thank Richard B. Larson for his comment that led to the correction of this error.

![Figure 1](image_url)

**Fig. 1.**—The minimum mass $M_{\text{sh}}$ of a homogeneous, spherically symmetric cloud which is able to self-shield and form a sufficient amount of H$_2$ at its center so that it cools on a dynamical timescale. Solid lines show $M_{\text{sh}}$ for the temperature $T = 1000$ K for the five different redshifts $z = 5, 10, 25, 35, \text{and } 50$. The two dashed lines show the effect of lowering or raising the temperature to log $T[K] = 2.2$ and 3.5 for the $z = 10$ case.