Ultraviolet Background Radiation from Cosmic Structure Formation

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5 September 2003

ABSTRACT

We calculate the contribution to the ultraviolet background (UVB) from thermal emission from gas shock heated by cosmic structure formation. Our main calculation is based on an updated version of Press-Schechter theory. It is consistent with a more empirical estimate based on the observed properties of galaxies and the observed cosmic star formation history. Thermal UVB emission is characterized by a hard spectrum extending well beyond 4 Ry. The bulk of the radiation is produced by objects in the mass range $10^{11} - 10^{13} M_\odot$, i.e. large galaxies and small groups. We compute a composite UVB spectrum due to QSO, stellar and thermal components. The ratio of the UVB intensities at the H and He Lyman limits increases from 60 at $z = 2$ to more than 300 at $z = 6$. A comparison of the resulting photoionization rates to the observed Gunn-Peterson effect at high redshifts constrains the escape fraction of ionizing photons from galaxies to be less than a few percent. Near 1 Ry, thermal and stellar emission are comparable amounting to about 10\%, 20\% and 35\% of the total flux at redshifts of 3, 4.5 and higher, respectively. However, near the ionization threshold for He\textsuperscript{II}, the thermal contribution is much stronger. It is comparable to the QSO intensity already at redshift $\sim 3$ and dominates at redshifts above 4. Thermal photons alone are enough to produce and sustain He\textsuperscript{II} reionization already at $z \approx 6$. We discuss the possible implications of our results for the thermal history of the intergalactic medium, in particular for He\textsuperscript{II} reionization.

Key words: cosmology: large-scale structure of universe — radiation mechanism: thermal — shock waves

1 INTRODUCTION

Neutral hydrogen in the intergalactic medium (IGM) produces a forest of resonant Ly\textalpha absorption lines in the spectra of high-redshift quasars. The connection of these features to the structure formation process has now been firmly established using N-body/hydrodynamic numerical simulations (\textsuperscript{??}). There is consensus that the observed IGM temperature results from a balance between photo-ionization heating and adiabatic cooling due to the Hubble expansion. Such photo-heating is provided by the extragalactic ultraviolet background (UVB) whose nature, origin and evolution have, therefore, been subject to considerable investigation.

At low redshifts, the ionization balance is consistent with a pure power-law ionizing spectrum. Traditionally QSOs have considered as the main sources of ionizing photon. However, a number of recent theoretical and observational developments are at odds with this.\textsuperscript{?)} find that the break in the redshift evolution of absorbers occurs at lower redshifts than predicted by numerical simulations using a standard QSO ionizing background (\textsuperscript{?), hinting at an incomplete description of the UVB. This led \textsuperscript{?)}}[B01 hereafter]BianchiAA2001 to recompute the UVB as a superposition of contributions from QSOs and galaxies. Those authors adopted an escape fraction of ionizing photons from galaxies, $f_{\text{esc}} \approx 10\%$. In spite of many theoretical (\textsuperscript{??}) and observational studies (\textsuperscript{??}) the determination of this parameter remains highly uncertain. In §3.3 below we use recent
measurements of the Gunn-Peterson effect in high redshift quasars (?) to argue for values of $f_{esc}$ as small as a few % for high redshift galaxies (§3.3), in agreement with independent estimates by ?). These authors set a 3σ (statistical) upper limit $f_{esc} \lesssim 4\%$ using photometry of 27 spectroscopically identified galaxies with $1.9 < z < 3.5$ in the Hubble Deep Field.

Similarly, high resolution simulations using a standard cold dark matter model and a standard QSOs ionizing background (?), produce a Ly-α forest lines with a minimum width significantly below that observed (??). This appears to require additional heat sources, for example photo-electric dust heating (?), radiative transfer effects (?), or Compton heating by X-ray background photons (?).

The situation around redshift $z \approx 3$ is more complicated. Apparently the spectrum must be very soft, with a large break at the He$^+$ edge, ?) reports an abrupt change of the C iv /Si iv ratio at $z \approx 3$. This may indicate a rapid and significant change in the shape of the ionizing spectrum, perhaps due to He ii reionization. This is corroborated by the detection of patchy He ii Ly-α absorption at similar redshifts. In addition, ?) analyzed the Gunn-Peterson He II absorption trough at $z = 3.05$ found the UVB to be characterized by a high softness parameter $S \approx 800$, in contrast to the harder spectrum deduced at $z = 2.87$. However, a more recent VLT/UVES study (?) using 7 QSOs finds no strong discontinuity for the quantity C iv /Si iv around $z = 3$ and suggests that it might not be a good indicator of He ii reionization.

?) concluded that the IGM temperature evolution differs considerably from simple expectations based on QSOs as sole ionizing power input. ?) found a similar result with an analogous but independent approach. Both analyses suggest at $z \approx 3$ the IGM temperature (of the gas at mean density) undergoes a “sudden” jump which is interpreted as associated with He ii reionization. It is worth pointing out that ?) recover higher temperature values and a smoother temperature jump than ?), providing a measure of the possible uncertainties.

In this paper we explore another source of UVB ionizing photons, namely thermal emission from shock-heated gas in collapsed cosmic structures. We find that thermal radiation, produced mainly in halos with temperatures between $10^6$ K and a few $\times 10^7$ K, is characterized by a hard spectrum with many photons above the H i and He ii ionization thresholds. As noted above we limit the escape fraction of UV ionizing photons from high redshift galaxies, $f_{esc}$, to a few %. Thermal emission provides a significant fraction of H i ionizing photons at redshift $\gtrsim 3$. In addition, we find that thermal emission plays a major role in the reionization of He ii, being comparable to the QSOs at redshift $\sim 3$ and dominating the flux at $z > 4$. In fact it turns out that thermal photons alone are enough to cause He ii reionization at $z \approx 6$. Our study is based on well understood emission mechanisms, bremsstrahlung and line emission from optically thin thermal plasma, so we expect our results to be robust.

Thermal emission is not the only possible source of ionizing photons in collapsing structures. An alternative is inverse Compton emission by relativistic electrons accelerated in large scale structure shocks (?). However, based on methods similar to those presented in ?) we find such a component to be negligible compared to that from QSOs, a result in agreement with ?).

The paper is organized as follows. In §2 we present the details of our model, that is we describe how we solve the radiative transfer equation and estimate the radiation due to QSOs, stars and, particularly, thermal emission. Our results are presented in §3 where the relevant features of the transmitted radiation flux are described. Finally, in §4 we discuss the thermal evolution of the IGM and in §5, we summarize our main results.

2 Model

2.1 Radiative Transfer

The mean specific ionizing flux, $J(\nu_o, z_o)$, observed at frequency, $\nu_o$, and redshift, $z_o$, is the solution to the cosmological radiative transfer equation which reads (?)

$$J(\nu_o, z_o) = \frac{c}{4\pi H_0} \int_{z_o}^\infty e^{-\tau_{eff}(\nu_o, z_o, z)} \frac{\Omega_m(1+z)^3 + \Omega_\Lambda}{1+z} \left(\frac{1+z_o}{1+z}\right)^3 j(\nu, z) \, dz$$

(1)

where $c$ is the speed of light, $H_0$ is the Hubble parameter, $\Omega_m$ and $\Omega_\Lambda$ are the matter and vacuum energy densities respectively, $z$ and $j(\nu, z)$ is the volume-averaged proper spectral emissivity computed at emission redshift, $z$, and at the appropriately blueshifted photon frequency $\nu = \nu_o (1 + z_o)/(1 + z)$. In addition, $\tau_{eff}(\nu_o, z_o, z)$ is the effective optical depth at frequency $\nu_o$ due to absorption of residual neutral gas in the IGM between $z_o$ and $z$. Following ?), for a distribution of discrete absorbers in an otherwise transparent (i.e. ionized) medium we write

$$\tau_{eff}(\nu_o, z_o, z) = \int_{z_o}^z d\zeta \int_0^\zeta dN_{HI} f(N_{HI}^1, \zeta')(1 - e^{-\tau(\nu')}).$$

(2)

The distribution of absorbers as a function of the H i column density and redshift, $f(N_{HI}, z') = \partial^2 N/\partial N_{HI} \partial z'$, is typically derived from counts of Ly$\alpha$ lines in QSOs absorption spectra. It is assumed that the number density of absorbers evolves with

1 Note that $S$ is defined as the ratio of the photoionization rate in H i over the one in He ii (§3.2).

2 We assume a flat cosmological model with normalized Hubble constant $h_{70} = H_0/70$ km s$^{-1}$ Mpc$^{-1} = 1$, $\Omega_m = 0.3$, $\Omega_\Lambda = 0.7$, $\sigma_8 = 0.9$. 

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redshift as $\partial N/\partial z' \propto (1 + z')^{-1}$, implying an evolution of the effective optical depth of the Lyα forest as $\tau_{\text{eff}} \sim (1 + z)^\gamma$ (?). This was the approach taken in B01 who, based on observations of the Lyα forest in the redshift range $1.5 < z < 4$ (?), determined a value of $\gamma \simeq 3.4$. For redshifts higher than 4, however, recent observations, including the detection of a nearly complete absorption through in the QSOs spectra discovered by the SDSS at $z \sim 6$, imply a much stronger evolution of that function (?). Therefore, for $z > 4$ we retain the power-law behavior for $\tau_{\text{eff}}$ but determine the power-law index $\gamma$ through a comparison of the power-law models with the new measurements of ?. In Fig. 1 we report the power-law models for different values of $\gamma$ together with observed data points. Actually the data do not follow a smooth behavior and, therefore, cannot be described in full detail by any of the smooth power-laws. The features in the data are likely due to inhomogeneities in the reionization process (patchiness). In addition, as we approach $z \sim 6$, the data indicate a sudden increase of the IGM optical depth. Both of these features in the redshift dependence of $\tau_{\text{eff}}$ are reflected in the redshift evolution of the ionization rates that ? inferred by modeling the distribution of absorbers in the IGM in a fashion similar to ours. Here we adopt a simple power-law at $z > 4$, which approximately describes the behavior of $\tau_{\text{eff}}(z)$. However, it is part of our objectives to compare the ionization rates measured by ?) with those produced by the radiation processes explored in this paper. We will heretofore need to bear these approximations in mind in order to properly interpret the results.

Each absorber is characterized by an optical depth $\tau(\nu') = N_{\text{HI}}(\nu') + N_{\text{HeII}}(\nu')$, due to H I and He II ionization (He I ionization is negligible). The value of $N_{\text{HeII}}$ can be derived from $N_{\text{HI}}$ by studying the radiative transfer within each cloud (?). For simplicity, we have assumed that all clouds are optically thin at the He II ionization threshold, which yields $N_{\text{HeII}}/N_{\text{HI}} \approx 1.8 J(13.6 \text{eV})/J(54.4 \text{eV})$ (?). While this approximation slightly underestimates the optical depth at $h\nu > 54.4$ eV, it does not severely affect the opacity of the most abundant ($N_{\text{HI}} < 10^{17} \text{cm}^{-2}$) clouds for $h\nu < 100 \text{ eV}$ (?). In Fig. 1 we report the power-law models for different values of $\gamma$ together with observed data points. Actually the data do not follow a smooth behavior and, therefore, cannot be described in full detail by any of the smooth power-laws. The features in the data are likely due to inhomogeneities in the reionization process (patchiness). In addition, as we approach $z \sim 6$, the data indicate a sudden increase of the IGM optical depth. Both of these features in the redshift dependence of $\tau_{\text{eff}}$ are reflected in the redshift evolution of the ionization rates that ? inferred by modeling the distribution of absorbers in the IGM in a fashion similar to ours. Here we adopt a simple power-law at $z > 4$, which approximately describes the behavior of $\tau_{\text{eff}}(z)$. However, it is part of our objectives to compare the ionization rates measured by ?) with those produced by the radiation processes explored in this paper. We will heretofore need to bear these approximations in mind in order to properly interpret the results.

2.2 Radiation from QSOs and stars in galaxies

The contributions of QSOs and stars to the ionizing UV background adopted in this paper are similar to those derived in B01. We briefly summarize here the main assumptions of the model. The QSO emissivity assumes a luminosity function that follows the double power-law model of ?. For $z < 3$, we adopt the parameters given in ?, obtained by fitting a sample of over 6000 QSOs with $0.35 < z < 2.3$. At $z > 3$, we include the exponential decline suggested by ?, describing the dramatic reduction in QSOs number density at high redshift. Finally the QSO spectrum in the ionizing UV range is modeled as a simple power law, $j(\nu) \propto \nu^{-1.8}$ (?).

For the stellar component, we assume a star formation rate that is constant from high redshifts to $z \approx 1$, then rapidly decreases to local values, as indicated by several galaxy surveys in the rest-frame non-ionizing UV (??). Synthetic galactic spectra (produced with the 2001 version of the ? code) have then been used to calculate the emissivity for the ionizing UV as a function of $z$. Finally, the internal absorption of radiation by the galaxy interstellar medium was modeled by a redshift-independent value for $f_{\text{esc}}$, the fraction of Ly-continuum photons that can escape into the IGM. As we shall show in §3 the values for the ionization rates recently measured by ?) constrain $f_{\text{esc}}$ to a level of a few %. In the following we will usually set $f_{\text{esc}} = 1\%$. We should also mention for completeness that we have neglected contributions from radiative recombinations in the absorbers (?). It would be possible to correct for this omission, but it is not strictly necessary for comparing the contributions from the various emission processes explored in this paper.

2.3 Thermal Emission

In this section we compute the contribution to the ionizing background flux from thermal emission from gas accreting onto dark matter halos. Since we calculate this for the first time here, we outline the assumptions of the model in more detail than for the QSOs and stellar components (§2.2).

The gravitational collapse of matter density enhancements during structure formation drives supersonic flows which eventually shock (?), e.g., ). The shocked gas is collisionally ionized and heated to temperatures $T \approx 10^6$ K ($\nu/10 \text{km s}^{-1})^2$ so that dense regions can efficiently radiate away thermal energy through bremsstrahlung and line emission (?). Where the gas overdensity is above a few hundred and the temperature $T \geq 10^5$ K, the ionization fraction of the shocked gas is Therefore, we expect that in general only a small fraction of the emitted radiation will be absorbed by local neutral gas and most of it will “escape”, adding to the background flux.

In the following, the description of the evolution of the baryonic gas, responsible for the emission of thermal radiation, is rather simplified. Although in principle it could be pursued in some detail through semi-analytic schemes (??), the latter provide considerably more information that we require and their level of sophistication is well beyond the scope of our current study. However, we do test our model predictions for the amount of cooling that must have occurred against independent empirical estimates. These are based on the observed cosmic star formation history and the distribution of stellar mass as a function of halo masses, that we infer from recent SDSS data. The comparison is detailed in ??2.4. In addition, we comment on possible additional effects neglected in our simplified approach. In general, however, the agreement that we find between our predictions and our empirical estimates is encouraging and supports our adopted approach.
Figure 1. Multi power-law fits of the redshift evolution of $\tau^{\text{eff}} \propto (1 + z)^\gamma$. Up to $z \leq 4$ a single curve corresponding to $\gamma \approx 3.4$ fits the data. For higher redshifts we show the cases of $\gamma = 5$ (dot), 5.5 (dash) and 6 (long dash) respectively. Data points are taken from ?).

2.3.1 Collapsed Halos

The volume averaged comoving thermal emissivity in units of erg s$^{-1}$ cm$^{-3}$ Hz$^{-1}$ can be written as

$$j_T(\nu, z) = \int dM \frac{dn}{dM}(M, z) M \epsilon[T(M), \nu] \min \left[ \frac{\tau_{\text{cool}}(M, z)}{\tau_{\text{Hubble}}(z)}, 1 \right],$$

where $dn/dM$ is the comoving number density of collapsed dark matter halos of mass $M$ at redshift, $z$, and $\epsilon(T, \nu)$ is the spectral emissivity per unit mass. The last term is the ratio of the cooling time to the Hubble time at redshift $z$ and accounts for the fact that any parcel of gas emits only as long as is allowed by its supply of thermal energy.

To describe $dn/dM$ we adopt the ?) formalism as updated by ?). This gives the mass function

$$\frac{dn}{dM}(M, z) = A \left( \frac{2}{\pi} \right)^{1/2} \rho_0 M^2 \left( 1 + \nu^{-2}\dot{\nu}^2 \right) \dot{\nu} \left| \frac{d \ln \sigma(M)}{d \ln M} \right| \exp \left( -\frac{\dot{\nu}^2}{2} \right)$$

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where \( \rho_0 \) is the current \((z = 0)\) average mass density of the universe, \( \dot{\nu} = a^{1/2} \delta_c(z)/\sigma(M) \), \( \sigma(M) \) is the mass variance of linear density perturbation corresponding to mass \( M \) at the current epoch and \( \delta_c(z) \) is the critical linear overdensity evaluated at present for a spherical perturbation that collapses at redshift \( z \). Both \( \sigma(M) \) and \( \delta_c(z) \) are taken from \( ? \) and the parameters \((A, a, p) = (0.322, 0.707, 0.3) \) are as in \( ? \). According to the spherical collapse model, for \( z \gg 1 \) the virialized dark matter halos are characterized by an overdensity \( \Delta_c \approx 18\pi^2 \) with respect to the critical value, \( \rho_c \). The baryonic gas is assumed to settle in approximate hydrostatic equilibrium within each halo at its virial temperature \((?\text{, e.g. })\) [brno98]

\[
T(M) = 1.4 \times 10^7 K \ h_{70}^{2/3} \left( \frac{M}{10^{12} M_\odot} \right)^{2/3} \left( \frac{\Omega_m}{0.3} \right)^{1/3} \left( \frac{\Delta_c}{18\pi^2} \right)^{1/3} \left( \frac{1+z}{5} \right)
\]

(5)

and with a density \( \rho_b = f_b \Delta_c \rho_{cr} \), where \( f_b \) is the halo baryonic fraction.

The spectral thermal emissivity, \( \varepsilon(T, \nu) \), is computed through the code by \( ? \)[version 1992] [rasm77] whereas the cooling function, \( \Lambda(T, Z) \), necessary to compute the cooling time, was taken from the work of \( ? \). Our choices are appropriate because, for regions characterized by an overdensity of order 100 or greater at redshift \( \sim 2 - 3 \) and higher, the ionization equilibrium is primarily determined by plasma collisions. We adopt an average metallicity within the collapsed halos \( Z = Z_\odot/20 \), thought to be typical for high redshift objects \((?)\).

Contributions to the integral in eq. (3) are limited below a certain mass threshold. For halo masses below \( M_{1c} = 2.6 \times 10^{11} h_{70}^{-1} (\Omega_m/0.3)^{-1/2} \left[ \varepsilon_{13.6eV}/(1+z) \right]^{3/2} M_\odot \), where \( \varepsilon_{13.6eV} \) is the photon energy in units of 13.6 eV, emission of thermal radiation is depressed exponentially because the low virial temperature is insufficient to ionize hydrogen. The most stringent constraint for low mass objects is due, however, to cooling. In fact, we find that the cooling correction in eq. (3) is

\[
\frac{\tau_{\text{cool}}}{\tau_{\text{Hubble}}} \equiv \frac{\kappa_B T \tau_{\gamma}^{-1}}{n\Lambda(T, Z)} \approx \alpha 0.4 h_{70}^{-1/3} \left( \frac{M}{10^{12} M_\odot} \right)^{2/3} \left( \frac{1+z}{5} \right)^{-1/2} \left( \frac{\Delta_c}{18\pi^2} \right)^{-2/3} \left[ \Lambda(T, Z = Z_\odot/20) \right]^{-1} \left( \frac{\Omega_m}{0.3} \right)^{5/6} \left( \frac{f_b}{0.15} \right)^{-1}.
\]

(6)

where \( n \) is the gas number density, \( Z \) is the gas metallicity, \( \Lambda \) is the cooling function and \( k_B \) is Boltzmann's constant. For \( Z \approx Z_\odot/20 \), the normalization for \( \Lambda \) given in eq. (6) holds to a good approximation for \( 10^{10.5} \leq T \leq T \approx 10^7 \text{,(cf. })\) [sdno93]. If the gas contraction induced by radiative cooling occurs isobarically the amount of radiated energy equals the enthalpy of the system and, \( \alpha = \gamma_{\text{gas}}/(\gamma_{\text{gas}} - 1) = 5/2 \). We point out that the halo life-time, instead of the Hubble time, should perhaps be used in eq. (3), but we will neglect these details in the current investigation. Thus, the above scenario should provide at least a lower limit to the potential amount of energy to be radiated (even without considering the effects of feedback, see below).

### 2.3.2 Cooling & Feedback

It is well known that feedback of energy and momentum associated with formation of stars plays a fundamental role in regulating the dynamics of the IGM, although a clear and coherent understanding of how is still lacking. For the purposes of the present investigation, the effects of feedback can be taken into account by computing the increase in the cooling time, \( \tau_{\text{cool}} \), caused by the injection of additional energy. The average energy deposited in each halo of mass \( M \) can be computed through two quantities: the fraction of baryons \( f_s \) that is converted into stars and, the amount of energy, \( \zeta_{SN} \), released by supernova explosions per unit of stellar mass formed. When normalized to the halo volume-integrated enthalpy, \( W = (5/2)n k_B TV = (5/2)f_s M k_B T/\mu m_p \), where \( V \) is the halo volume and \( \mu m_p \) is the gas mean molecular weight, the deposited energy is

\[
\frac{\Delta E_{FB}}{W} = f_s f_c M \zeta_{SN} \approx 2.25 \left( \frac{f_s}{0.15} \right) \left( \frac{\zeta_{SN}}{10^{49} \text{erg} M_\odot^{-1}} \right) \left[ \frac{T(M)}{10^6 K} \right]^{-1}.
\]

(7)

The value assumed for \( \zeta_{SN} \) varies in the literature. If we take a ratio of supernova per total mass converted into stars \( \sim 4 \times 10^{-3} \), as suggested by a standard Salpeter initial mass function, then \( \zeta_{SN} \sim 4 \times 10^{48} \text{ erg} M_\odot^{-1} \), although based on metallicity arguments \( ? \) suggests a value 4 times larger \((?, \text{see also})\) [bookbinderetal80, ceo92]. We notice, incidentally, that the deposition of the assumed amount of energy corresponds to \( \mu m_p f_s \zeta_{SN} \approx 1 \text{ keV/part.} \) as observed in the core of small groups of galaxies \((?)\). Obviously, for the assumed parameters, the energy injection will only substantially affect gas within halos with temperatures \( \lesssim 10^5 \text{ K} \). This result is in agreement with the more sophisticated feedback prescriptions presented in \( ? \).

With the additional energy expressed in eq. (7) the radiative lifetime of a halo, as in eq. (6), will be extended by a factor \((1 + \eta \Delta E_{FB}/W)\). Here we have introduced yet another parameter, \( \eta(M) \), the fraction of energy deposited through feedback processes that is available for conversion into thermal radiation. We do so in order to account for various effects produced by feedback. For example, when \( \Delta E_{FB}/W \geq 1 \) gas is more likely to be blown out of its host halo, thus strongly inhibiting cooling \((?)\). This would imply \( \eta \leq 0 \). In addition, the evolution of a cooling core will be altered and possibly disrupted by the occurrence of merger events \((? \text{, cf. }?)\), making \( \eta < 1 \). However, as we shall show in the following, most of the radiation is

\[\text{The code assumes the emitting plasma in collisional equilibrium and neglects the effects of an ambient radiation field. It includes the following atomic processes: collisional ionization, collisional excitation followed by auto-ionization, radiative recombination, and dielectric recombination. Collisional ionization and excitation are assumed to be produced by free electrons with a Maxwellian distribution of energies (}).\]

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contributed by halos within a narrow mass range. Thus, in practice, we only explore the case \( \eta \simeq 1/2 \), which we consider appropriate for the mass range of interest here (cf. ?).

### 2.3.3 Halo Emissivity Distribution

In Fig. 2 we show which halos are primarily responsible for the production of thermal far UV emission. For this purpose we define arbitrarily a UV luminosity \( L_{\nu} \), as the integral between 13.6 and 100 eV of \( j_T(\nu, z) \) given in eq. (3). We then plot for different redshifts and as a function of the temperature of the emitting halo population, histograms of \( dL_{\nu}/d\log T \), that is the total UV luminosity of halos with virial temperature in a logarithmic interval centered on \( T \). According to our model, most of the emission is produced by halos with temperatures between 10\(^6\) K and a few \( \times 10^7 \) K, corresponding to masses \( 10^{11-13} M_\odot \). The contribution from halos below this range is reduced because of their short cooling time. This may be seen by comparing with the highest (magenta) curve which corresponds to the case in which the corrective term that accounts for the effects of cooling is omitted. Notice that the “bumpy” shape of these curves for \( T \leq 10^7 \) K is a reflection of the structured temperature dependence of the cooling function. On the other hand, the bend at the high mass end is solely due to the paucity of massive halos above a certain threshold which depends on both redshift and \( \sigma(M) \) for \( \sigma(M) \propto M^{\beta} \), \( dL_{\nu}/d\log T \propto dL_{\nu}/d\log M \propto M^{\beta/2-1/3} \). The blue line corresponds to the feedback case described in §2.3.2. As expected, the extra injected energy primarily affects small structures which are now able to produce a larger amount of radiation (by reradiating the injected energy). Compared to the no-feedback case (cyan line) the spectra in Fig. 4 are increased by a factor of a few in the low energy UV range but are basically unchanged in soft X-rays. This is a reflection of the fact that feedback mostly affects smaller (colder) halos. In any case the objects which dominate the thermal emission are only weakly affected by the specific feedback prescription we adopt.

### 2.4 Comparison with Observational Data

We will now attempt to test our model against observable quantities that are related to the occurrence of radiative cooling. For this purpose, we have used SDSS data presented in ? for a representative sample of more than 10\(^5\) galaxies, to construct a function, \( g_*(M) \), describing how the total stellar mass is distributed over halo mass in the present day universe. The dataset includes the following quantities of interest to us: colors in the \( g, r, i \) bands, stellar masses and concentration index.

For each galaxy, the halo mass is inferred from its luminosity through either the Tully-Fisher or Faber-Jackson relations depending on whether the galaxy is disk or bulge dominated, respectively. We first divide the sample into spirals and ellipticals on the basis of the concentration index defined as the ratio of the Petrosian half-light radius to the Petrosian 90\% light radius, \( C = r_{90}/r_{50} \). Thus, a galaxy is classified as spiral or elliptical/spheroidal depending on whether or not \( C > 3 \) (?). The SDSS color indexes are converted to Johnson’s \( UBV RCIC \) system according to the relations given by ?). Next, for spirals we use Verheijen’s (2001) relation between the \( I \)-band magnitude, \( M_I \), and the rotational velocity \( V_{\text{flat}} \), measured in the outer parts of the galaxy disk where the rotation curve flattens out. For ellipticals we employ the relation between \( B \)-band magnitude, \( M_B \), and dispersion velocity as determined by ?). Temperature and velocity are then related according to \( k_B T = (1/2)\mu m_p v^2 \) and temperature and mass as in eq. (5). The results for \( g_\nu \) are shown in Fig. 3, where the histogram represents the fraction of stellar mass in the local universe as a function of the host halo virial temperature and mass.

The formation of a mass \( M_* \) of stars implies that the thermal energy/enthalpy associated with the equivalent gas mass, \( \alpha M_* k T \), must have been radiated away. Here \( T(M) \) is the virial temperature of the host halo and \( \alpha = 5/2 \) as in eq. (6). In fact, this is just a lower limit. More precisely, the amount of radiated thermal energy is connected with the rate at which cold gas forms inside a halo, \( \dot{M}_{\text{cold}} \), rather than the rate at which stars form, \( M_* \). Obviously, averaged over the halo lifetime \( \dot{M}_{\text{cold}} \gg M_* \); however, there is no straightforward way of inferring the former from the latter. A number of empirical facts and theoretical arguments suggest a proportionality between the cosmic-volume averages of the two quantities,

\[
\langle \dot{M}_{\text{cold}} \rangle = \lambda(M,z)\langle M_* \rangle (M,z) \tag{8}
\]

with the proportionality factor \( \lambda(M,z) \) taking values of order of several at high redshifts (in appendix A we provide an estimate of \( \lambda(M,z) \) based on Press-Schechter formalism). Since we only have poor statistical information about how the time history of star formation depends on the halo mass, we assume that, on average, stars in halos of a given mass \( M \) formed at the measured cosmic star formation rate (see, e.g. B01) multiplied by our stellar mass distribution function, \( g_\nu(T) \). That is

\[
\langle \dot{M}_* \rangle (M,z) = g_\nu(M) \text{ SFR}(z). \tag{9}
\]

This is the simplest assumption consistent with the results in Fig. 3. Thus, the differential amount of energy radiated as a consequence of star formation by halos of mass \( M \) is

\[
\frac{d\dot{E}}{dM}(z, M) = \alpha \frac{k_B T(M)}{\mu m_p} \langle \dot{M}_{\text{cold}} \rangle = \alpha \frac{k_B T(M)}{\mu m_p} g_\nu(M) \lambda(M,z) \text{ SFR}(z). \tag{10}
\]

which amounts to a differential volume-averaged thermal emissivity

\[
\frac{d\dot{E}}{dM}(z, M) \frac{\epsilon[T(M), \nu]}{\Lambda[T(M), Z]} \tag{11}
\]
Figure 2. Histograms of the integrated thermal UV emission between 13.6 and 100 eV as a function of the virial temperature of the emitting halo population for four different redshifts. The various lines correspond to cases with (cyan) and without (magenta) accounting for the effects of radiative cooling, and with cooling+feedback (blue). The red thick line was obtained based on the observed cosmic star formation history and the distribution of stellar mass as a function of halo virial temperature as reconstructed from SDSS data (see Sec. 2.4 for details).

The last term in the above equation describes the spectral distribution of the radiated energy. Using the above ‘empirical’ thermal emissivity, and the cosmic star formation rate $SFR(z)$ summarized in B01, we have recomputed the quantity $dL_{\text{uv}}/d\log T$ presented in Fig. 2 in order to compare it with our model predictions. The results are illustrated by the thick (red) line plotted in Fig. 2 for each redshift (the adopted value of $\lambda$ was computed as detailed in appendix A).

The comparison between model and ‘empirical’ predictions is only meaningful for halos with $\tau_{\text{cool}} < \tau_h$ (no star formation can occur otherwise), i.e. with virial temperatures somewhat below the overlap point of black (cooling accounted for) and magenta (no cooling) curves in Fig. 2. The actual dividing line should occur a few times below this point of overlap, typically above several $\times10^6$ K. This is because the cooling suppression factor in the former curve is based on the assumption that the halo age equals the Hubble time and is thus a bit overestimated. With this clarification, we conclude that the ‘empirical’
Figure 3. Histograms representing \( g_* \), i.e. the fraction of stars in the local universe hosted in halos of virial temperature \( T \), and mass \( M \). \( M \) and \( T \) are related through eq. (5) where we set \( z = 0 \). The histogram was built on the basis of the SDSS data (7).

curve agrees quite well with the model predictions for star forming halos with virial temperatures above a few \( \times 10^5 \) K, where the cooling effects are strong. Below \( \times 10^5 \) K the agreement worsens considerably. This is somewhat expected and is typically attributed to the suppression of star formation and/or gas blow-away due to feedback effects. Nevertheless, this is not too worrisome since the contribution of these objects to the UV background is not crucial, even in the ‘feedback-case’.  

Finally, we note that the model and empirical curves have been derived using completely different methods and assumptions, with values for the few free parameters involved (e.g. \( \theta \)) taken from the studies in which these were introduced, rather than in order to improve our results. Nevertheless, the two approaches give very similar estimates of the UV thermal emission; this is encouraging as, in principle, the agreement could have been much worse.
3 RESULTS

3.1 Ionizing Flux

Fig. 4 shows the mean ionizing flux due to thermal emission with \((\eta = 0.5)\) and without feedback effects, that due to QSOs, and that due to galaxies \((f_{esc} = 1\%)\) at four different redshifts. The imprint of the H i and He ii IGM absorption, with a significant reduction of the flux close to the Lyman limits of the two species, is clearly visible in each spectrum. We find that thermal emission provides an important contribution to the average UV radiation flux. Apart from this, the three sources of radiation have very different spectra and redshift dependence. At \(z \sim 3\) and near the H i ionization threshold, thermal emission with feedback effects corresponds to \(\sim 10\%\) of the QSO contribution and is comparable to the stellar component. When feedback effects are not included the thermal flux in this spectral region is reduced by a factor of a few. At higher redshifts the thermal flux becomes progressively more important relative to the QSO component, whereas its relation to the

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stellar component is virtually unchanged. This latter feature makes sense because the production of stellar photons is linked to the formation of stars and, therefore, of cold gas, which in turn produces the thermal photons. Since the two components were modeled independently, this provides a nice consistency check. At photon energies near the He II ionization threshold the trends are similar, except that now the thermal flux plays a more dominant role because of its relatively hard spectrum. Thus, in this energy range, and independently of feedback effects, thermal emission dominates the stellar contribution (even for $f_{esc} \sim 10\%$). It is comparable to the QSOs flux at $z \sim 3$, and becomes the dominant source at redshifts $\geq 4$.

### 3.2 Softness parameter

Both for the individual spectral components and for the total spectrum, we have computed the redshift evolution of the spectral softness parameter,
Figure 6. Left: Photo-ionization rates defined in eq. 13 as a function of redshift for ionization of H I due to emission from: QSO (black dash), stellar for $f_{\text{esc}} = 1\%$ (red long dash) and shocked IGM with feedback effects (blue solid) and without (cyan dot). The thick solid line is the total. The data points are from ?) and corrected for our cosmological model. Right: same as for left panel but now for coefficients relative to He II (right). The stellar component (red long dash) has now been multiplied by a factor 10 for visualization purposes.

\begin{equation}
S_L \equiv \frac{J_{\text{HI}}}{J_{\text{HeII}}},
\end{equation}

where $J_{\text{HI}}$ and $J_{\text{HeII}}$ are the UVB intensities at the H I and He II Lyman limits respectively. These are shown in Fig. 5. As already anticipated, the thermal emission is characterized by small values of $S_L$, ranging from a few at low redshift to 100, at $z \approx 6$, with minor differences introduced by the adopted feedback prescription. A similar evolution is seen for the QSO spectra, although with higher $S_L$ values in the range 80-300; the stellar component instead shows less pronounced evolution, increasing by only a factor of $\sim 3$, but maximal $S_L$ values that exceed 1000 at $z \gtrsim 3$. The composite spectra (feedback case) $S_L$ evolution resembles very closely that of QSOs, even at high redshift: this is somewhat fortuitous as at $z \gtrsim 4$ the UVB is dominated by the sum of stellar and thermal contribution. The above values can be compared with the available data. ?), from an analysis of the quasar Q0302-003 ($z = 3.286$), find that around $z = 3.2$ the ratio of the H I to He II photoionization rates, is larger than 400. As $S_L = r^{-1} \Gamma_{\text{HI}}/\Gamma_{\text{HeII}}$, where $r$ depends on the spectrum shape and is about 4 for our composite spectrum, from Fig. 5 we find that this value corresponds to our estimate for $z = 3.4$, a very close match. In our model the results of Heap et al. would not be necessarily caused by a jump in the parameter $S$ at $z \approx 3$, but by the evolution of the composite spectrum, particularly the decrease in the ratio of thermal to the QSO flux. In the final Section we will reconsider this point.

3.3 Photo-ionization Rates

Fig. 6 shows the photo-ionization rates in units of $10^{-12} \text{s}^{-1}$ defined as

\begin{equation}
\Gamma_{12}(z) \equiv \frac{\Gamma(z)}{10^{-12} \text{s}^{-1}} = 4\pi \int_{\nu_s}^{\infty} \sigma_s(\nu) \frac{J(\nu_z)}{h\nu} d\nu
\end{equation}

for the various UV radiation components discussed above and as a function of redshift with the same notation as in Fig. 4. In addition, we also plot the photo-ionization rates inferred by ?) from the observed Gunn-Peterson effect in high-z QSO spectra using semianalytic representations of Lyα absorption.

The left panel of Fig. 6 is relative to ionization of H I and contains a number of important features. First, we notice that although the emission from QSOs is able to produce the ionizing flux observed at $z \sim 2-3$, it falls short at higher redshifts, a well known fact. In our formulation, it results from the assumed rapid decline of the QSO number density for $z > 3$, as derived from the SDSS (?§2.2)FanAJ2001. We note here, however, that ?) in a recent paper questioned this view: by tweaking the bright end of the SDSS QSO luminosity function and allowing for objects fainter than the survey limit, QSOs alone can in principle account for the UV background at $z > 4$. 

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As a second result, the comparison between the $\Gamma$ values predicted by \cite{7} and the stellar ionization rates assuming $f_{esc} = 1\%$ (red dash curve) imply, according to our model, that $f_{esc}$ has to be smaller than a few %. In fact, had we assumed, e.g., $f_{esc} \approx 10\%$ the observed $\Gamma$ would have been overpredicted by a factor 5 for $z \geq 4$ and by a factor 2 at $z \sim 2 - 3$. In B01 the value $f_{esc} = 10\%$ was preferred, based on comparisons of the predicted ionizing flux with estimates from the proximity effect at $2 \leq z \leq 4$. Estimates of the UV background from the proximity effect are known to be larger than those obtained via theoretical models of the IGM opacity (i.e. the work of \cite{7} we are using here), probably because of a bias of the QSO distribution towards the denser environments the (\cite{7}) or because of systematic errors due to line blending (\cite{7}). Furthermore, proximity-based UV backgrounds have been traditionally derived in a flat Einstein-De Sitter cosmology (\cite{7}); this may cause a further overestimation by 40% if the true cosmology is the one we are using here (\cite{7}).

Our conclusions on $f_{esc}$ are not affected by the high values of the measured photo-ionization rates for $5 \leq z \leq 6$. As already pointed out in \S 2.1 (see Fig. 1 there) that bump is due to a deviation of the IGM optical depth from the assumed smooth power-law evolution. Had we modeled in more detail the behavior of the IGM optical depth we would have also recovered higher values for $\Gamma$ in the same redshift range. Therefore our statement above is valid throughout the whole redshift range.

A third interesting point concerns the relative contribution of stellar and thermal emission to the metagalactic UV flux at high redshifts where the QSO emission drops rapidly. Although one can in principle reproduce the observed photo-ionization rates (thick solid line) with an escape fraction $f_{esc} \sim 2\%$, it is also possible that at least half of those rates are due to thermal emission. This is the case when a fraction $\eta \approx 0.5$ of the feedback energy is re-radiated through thermal processes (blue solid line). However, thermal emission is characterized by a much harder spectrum than the stellar one. This might serve not only as a way to discriminate observationally between the two components, but it might have implications for the temperature evolution of the IGM, studied in the next Section.

Finally, the right panel of Fig. 6 shows the He $\Pi$ ionization rates. According to the plot, above 4 Ry thermal emission is comparable to QSOs at $z \sim 3$ but completely dominates the radiation flux at higher redshifts. Stellar emission is negligible at these energies. These conclusions are independent of the assumed $f_{esc}$ and depend only weakly on feedback. This result is very important in terms of the IGM evolution and has not previously been noticed.

3.4 He $\Pi$ Reionization

The previous results hint at the intriguing possibility that He $\Pi$ reionization could have been powered by UV light from cosmic structure formation. A necessary condition for reionization is that the number of ionizing photons emitted per He $\Pi$ atom (we assume all He to be in this form prior to its reionization) in one recombination time is equal to, or larger than, unity or, equivalently, that the photoionization rate equals the recombination rate. The volume-averaged recombination time, $t_{rec}$, is given by

$$t_{rec} = \frac{1.16 n_p \alpha(T) C^{-1}}{\langle n_p^2 \rangle / \langle n_p \rangle^2},$$

where $\alpha$ is the radiative recombination coefficient, and the clumping factor $C \equiv \langle n_p^2 \rangle / \langle n_p \rangle^2 > 1$ is meant to include the effects of density inhomogeneities inside the ionized region; we have used a helium to hydrogen number ratio $y = 0.08$ and assumed a temperature of the reionized gas $T \approx 40000$ K. Then, from eq. 14 we find

$$t_{rec} = 0.17 C^{-1} \left( \frac{1 + z}{10} \right)^{-3} \text{Gyr}.$$  \hfill (15)

Note that the expansion time scale is $H^{-1}(z) = 0.87 \text{Gyr} \left( (1 + z)/10 \right)^{-3/2}$, and hence below redshift $z = 10$ in the diffuse ($C \approx 1$) IGM the recombination time scale is always comparable or longer than the Hubble time. From Fig. 6 we find that at $z \sim 6$ thermal emission dominates the photoionization rates and alone provides $5.5 \times 10^{-17}$ (10$^{-16}$) He $\Pi$ photoionizations/s in the no-feedback (feedback) case. According to eq. (15), the He $\Pi$ recombination rate at the same redshift is $6.4 \times 10^{-17} C$ s$^{-1}$, i.e. a comparable value. Hence, it appears that structure formation can produce He $\Pi$ reionization around $z = 6$, without the contribution from any other process and essentially independently of the feedback prescription adopted.

An additional check of the above results can be performed. At $z = 6$ the fraction of mass in halos with virial temperature (calculated from Press-Schechter, eq. 4) larger than 4Ry/$k_B$ is equal to $f_4 = 3\%$ (\cite{7}). In order to provide at least one He $\Pi$ ionizing photon per He atom (we have assumed a helium abundance $y = 0.08$ and that He is all singly ionized), the required mean photon energy has to be equal to 4 Ry/$f_4$ = 145 eV. By averaging again over the mass distribution we find that this mean energy is 1.7 keV. Hence this simple argument confirms that He $\Pi$ reionization can be caused by structure formation.

Whether or not this possibility is fully compatible with other observational results is not clear. Nevertheless, in \S 4 we show that this scenario predicts a temperature evolution of the IGM that maybe consistent with existing observations.

4 DISCUSSION

Fig. 7 illustrates the thermal history of the IGM. Data points are from \cite{7}, who measured the cut-off in the $b - N_{HI}$ relation in a sample of nine quasar spectra in the redshift range 2.5-4.0. Despite the large uncertainties, there is an indication of a
temperature peak at redshift $z \approx 3$. On the other hand, the plotted curves represents different evolutionary scenarios and are (numerical) solutions to the equation

$$\frac{dT}{T} = (\gamma - 1) \left[ \frac{(\gamma - 1)dt}{y_in_Hk_BT} + \frac{dn_H}{n_H} + \frac{dy_i}{y_i} - d \frac{1}{\gamma - 1} \right] + \gamma \frac{d\mu}{\mu}$$

(16)

where $y_i = \sum y_i$, $y_i = n_i/n_H$ is the concentration of the $i$-th species, $\mu = y_i^{-1} \sum y_i \mu_i$ is the mean molecular weight and the IGM mean density evolves as $n_H(z) = n_H(0)(1 + z)^3$. The calculation accounts for photoheating from a time dependent ionizing background, Compton cooling on Cosmic Microwave Background photons and non–equilibrium evolution of hydrogen and helium ionic species in a cosmological context (see further details; we do not include metal cooling which is negligible for typical IGM metallicities). The temperature of the medium at $z = 10$, where we start our integration, has been computed self-consistently.

We emphasize that the various curves presented in Fig. 7 are meant to illustrate a variety of possibilities and not to depict a specific scenario. In this perspective, we first discuss the dotted line, representing the case for a UVB generated by QSO+stellar contributions given by the photoionization rates presented in Fig. 13. In this case $\text{H} \, \text{I}$ reionizes at $z \sim 6$ and $\text{He} \, \text{II}$ around $z \sim 3$ (therefore their heating and ionization rates are set to zero prior to their respective epoch of ionization). At hydrogen reionization the IGM temperature jumps to $\log T = 4.13$ and decreases later on due to adiabatic cooling. A sharp rise is seen at $\text{He} \, \text{II}$ reionization at $z_{\text{He} \, \text{II}} = 3$, although not quite enough to fit all the data points. When the contribution from thermal emission is added (lower solid curve), the temperature is slightly larger at both reionization events, and also thereafter. As expected, the effect is small since the previous photoionization rates were already sufficient to keep the gas fully ionized and, therefore, their increase changes little the energy input to the IGM.

As a second example the long-dash-dot curve illustrates a similar case but with the photoheating rates increased twofold to mimic radiative transfer effects as discussed in 7). Again, $\text{H} \, \text{I}$ reionizes at $z \sim 6$ and $\text{He} \, \text{II}$ around $z \sim 3$. This scenario provides a somewhat better fit. Obviously it is possible to improve even further the agreement with the data points by fiddling with the available parameters so as to produce the desired behavior. At least in principle the data points can be reproduced (see also 7).

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Motivated by our findings in 3.4 we also explore a less conservative scenario in which He $^\text{II}$ reionization occurs as early as $z_{\text{HeII}} = 6$ and is mainly driven by thermal emission. By construction, the temperature peak around $z \sim 3$ does not develop in this case. Nevertheless, most of the data are fit at the 1 $\sigma$ level and all except one point at 2$\sigma$ level.

The possibility that cosmic structure formation is responsible for He $^\text{II}$ reionization at an early epoch is intriguing and challenges the common wisdom that such process is due to QSOs and occurs around $z = 3$. It is worthwhile then to revisit the arguments in support of the standard conjecture. These are essentially three: (i) an apparent boost of the IGM temperature around $z = 3$ which could be due to heating associated with He $^\text{II}$ reionization; (ii) an abrupt change of the $CIV$/$SiIV$ ratio at $z \approx 3$ indicating a significant change in the ionizing spectrum; (iii) the detection of a patchy He $^\text{II}$ Ly$\alpha$ absorption at about the same redshift, suggestive of the final (overlapping) stages of reionization process. Taken together these facts seem to justify the conjecture above.

However, as already pointed out by $\cite{7}$, these three observations probe different structures: metal line ratios probe high density regions; the effective He $^\text{II}$ Ly$\alpha$ absorption is mainly produced by neutral gas in the voids; the IGM temperature measurements concern the ionized gas with density around the cosmic mean. Radiative transfer calculations show that the ionization fronts propagate first through the low density regions and only subsequently penetrate the denser clumps such as filaments and halos $\cite{7}$.

We also emphasize the following difficulties. The existence of the peak in the IGM temperature evolution needs further data to be fully confirmed. Although a wavelet analysis seems to confirm it $\cite{7}$, the data are also marginally consistent with no feature at $z = 3$, in which case early reionization would provide a better fit to the temperature data (Fig. 7). The conclusion about the jump in the $CIV$/$SiIV$ ratio at $z = 3$ is still debated $\cite{7}$, e.g. $\cite{bosara98,Boksenberg1998}$. A more recent VLT/UVES study $\cite{7}$ using 7 QSOs finds no strong discontinuity in the ratio around $z = 3$ and suggests that $N_{CIV}/N_{SiIV}$ might not be a good indicator of He $^\text{II}$ reionization. On the other hand, we have shown (Sec. 3.2) that our smoothly increasing $S_L$ is quantitatively consistent with the $\cite{7}$ observations. Finally, (?) points out that a considerable number of hard photons are required at $z > 3$ to reproduce the observed abundances of O $\text{VI}$. At that epoch a sizable fraction of cosmic volume has to be transparent at photon energies above the He $^\text{II}$ ionizing threshold.

In conclusion, there is no conclusive argument against an early ($z \approx 6$) start of He $^\text{II}$ reionization which may result from the hard UV field emitted during the virialization of large galaxies/small groups. Clearly, this conclusion needs to be confirmed by numerical simulations, although this task might be far from easy, as the correct treatment of the strongly radiating cooling inhomogeneities in the gas poses challenging numerical problems. Resorting to improved observations of the proximity effect might then provide a check of the scenario proposed here.

5 SUMMARY AND CONCLUSIONS

We have shown that UVB ionizing photons can be copiously produced by thermal emission from shock-heated gas in collapsing cosmic structures. Our calculations are based on an implementation of the extended Press-Schechter theory. However, the estimated amount of thermal radiation is consistent with that inferred from an independent analysis based on observed, high redshift star formation rates (7, B01); $\cite{lanzettaetal02}$, and the distribution of stellar mass as a function of halo virial temperature as reconstructed from recent SDSS data (7).

Thermal radiation is characterized by a hard spectrum extending up to photon energies of order $h\nu \sim k_BT$. This is well above the H $\text{i}$ and He $^\text{II}$ ionization thresholds for virial temperatures above $10^6$ K. The bulk of the emission is produced by halos with temperatures between $10^6$ K and a few $\times 10^7$ K, corresponding to masses $10^{11}-13 M_\odot$.

We use simplified radiative transfer to compute the transmitted flux due to QSO, stellar and thermal emissions. Importantly, the resulting associated photoionization rates, when compared to measurements of the Lyman series Gunn-Peterson effect in the spectra of SDSS high redshift quasars (77), imply an escape fraction of UV ionizing photons from galaxies, $f_{\text{esc}}$, below a few %. This result is in agreement with very recent and independent determinations of $f_{\text{esc}}$ carried out by $\cite{7}$, who set a $3\sigma$ (statistical) upper limit $f_{\text{esc}} \lesssim 4\%$ for a sample of spectroscopically identified galaxies of redshift $1.9 < z < 3.5$ in the Hubble Deep Field.

It turns out that near the H $\text{i}$ ionization threshold, thermal emission is comparable to the stellar component and amounts to about 5-10 $\%$, 15-30 $\%$ and 20-50 $\%$ of the total at redshifts of 3, 4.5 and higher respectively. The quoted range depends on the fraction of feedback energy allowed to be re-radiated through thermal processes. Near the ionization threshold for He $^\text{II}$, the thermal contribution is much stronger. It is comparable to the QSO input already at $z \sim 3$, and it dominates for $z > 4$. Thus, this contribution, with a typical softness parameter $S_L = 10-100$, is expected to play a major role in He $^\text{II}$ reionization. In principle structure formation alone provides enough photons to produce and sustain He $^\text{II}$ reionization at $z \sim 6$. These conclusions are independent of our feedback prescriptions, which only affect low virial temperature systems. In our scenario, He $^\text{II}$ ionizing photons are produced primarily by relatively large collapsing structures (with $T_{vir} \gtrsim 10^6$ K).

The thermal spectrum $J_{\nu} \propto \nu^{\alpha}$ is very hard, with a slope in the range $\alpha \approx 1-2$ depending on the importance of the supernova feedback we include. The latter process primarily affects the smallest systems, increasing their emissivity through reradiation of SN input. As these smaller objects dominate the 1-4 Ry band, the spectrum in this energy range becomes flatter as feedback is increased. Measuring the evolution of the thermal component of the UVB provides a powerful method to evaluate the importance of SNe input into the intergalactic medium.

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ACKNOWLEDGMENTS

We are indebted to X. Fan for providing us with his measurements of the photo-ionization rates and to G. Kauffmann for making available to us her SDSS data sample. In addition, we wish to thank B. Ciardi, G. De Lucia, F. van den Bosch and V. D’Odorico for useful discussions. This work was partially supported by the Research and Training Network ‘The Physics of the Intergalactic Medium’, EU contract HPRN-CT2000-00126 RG29185.

APPENDIX A: PRESS-SCHECHTER ESTIMATE OF $\lambda(M, Z)$.

The lifetime of a halo is

$$\tau_h(z, z_f, M) \equiv \tau_H(z) - \tau_H(z_f(M)), \quad (A1)$$

with the formation redshift, $z_f(M)$, of a halo of mass $M$ defined as the formation time of a progenitor with mass $M/2$ (??). Its relation to the running redshift $z$ (assumed $\gg 1$) is statistical and $M$-dependent, and is described by

$$\frac{1 + z_f}{1 + z} = 1 + \frac{1}{1 + z} \left[ \tilde{\omega}(0) \right] \left[ \sigma^2(M/2) - \sigma^2(M) \right]^{1/2}. \quad (A2)$$

where $\tilde{\omega}$ takes values (effectively between a few $\times 10^{-3}$ and a few) according to a probability function defined within the extended Press-Schechter formalism (??, see also approximation formulae in ??) and $\delta_c$ and $\sigma(M)$ have been defined in §2.3.1. Thus, at a given redshift $z$ the formation rate of cold gas within a halo of mass $M$ that formed at $z_f$, averaged over the halo lifetime $\tau_h(z, z_f, M)$, is

$$\langle \dot{M}_{\text{cold}} \rangle_{\text{time}}(z, z_f, M) = \frac{M_{\text{cold}}}{\tau_h(z, z_f, M)} + (1 + R) \langle \dot{M}_* \rangle_{\text{time}} \quad (A3)$$

where $R$ ($\approx 1$) is the mass return fraction by stars and the subscript explicitly indicates time, as opposed to cosmic-volume. After introducing the characteristic star formation timescale: $\tau_* \equiv M_{\text{cold}}/\dot{M}_*$ (??, e.g.) we rewrite eq. (A3) as

$$\langle \dot{M}_{\text{cold}} \rangle_{\text{time}} = \left( 1 + R + \frac{\tau_*}{\tau_h} \right) \langle \dot{M}_* \rangle_{\text{time}} \quad (A4)$$

It follows that at early epochs when $\tau_h \ll \tau_*$ radiatively cooled gas is produced at a much faster rate than stars. However, the two quantities converge at late times, when $\tau_h \gg \tau_*$. This is consistent with, e.g., ?) semianalytic model of a Milky Way-like galaxy, in which star formation occurs at a rate proportional to the amount of available cold gas. Based on the diagram of Fig. 2 of his paper we readily infer that $\lambda \simeq 5.5$, at $z \simeq 9$ shortly after the galaxy starts evolving, $\lambda \sim 2.0$ by $z \simeq 3$ and $\lambda \sim 1.2$ by $z \simeq 0$.

Values of $\tau_* \sim \text{Gyr}$ or so have been inferred from observations of nearby galaxies by, e.g., ??), although the author suggests higher values during the early stages of the evolution of galactic disks. Alternatively, $\tau_*$ can be taken from semianalytic models which successfully reproduce a number of observed galactic properties. In the model of ??), the star formation rate is prescribed according to $\dot{M}_* = 0 \dot{M}_{\text{cold}}/\tau_{\text{dyn}}$, where $\theta = 0.1 - 0.3$ is a control parameter and $\tau_{\text{dyn}} \equiv \tau_{\text{gal}}/v_{\text{gal}} = 0.08 \sqrt{\Omega_m/0.3} \tau_H(z)$. In this formulation

$$\tau_* \equiv \frac{\tau_{\text{dyn}}}{\theta} = 0.08 \left( \frac{\Omega_m}{0.3} \right)^{1/2} \frac{\tau_H(z)}{\theta}. \quad (A5)$$

It implies lower values than those found by ??) and, therefore, is more conservative for our present calculations. After using this expression in eq. (A4) and volume averaging over halos of a given mass $M$, a comparison to eq. (8) leads to

$$\lambda(M, z) = 1 + R + \frac{1}{\theta} \left( \frac{\tau_{\text{dyn}}(z_f)}{\tau_h(z_f)} \right) \simeq 1 + R + 0.1 \frac{\Omega_m}{\theta} \left[ \left( \frac{1 + z_f}{1 + z} \right)^{3/2} - 1 \right]^{-1} \quad (A6)$$

where the last equality holds for $z \gg 1$. The averaging operation is effectively over the halo formation time which we achieve through the relation in eq. (A2) and the probability functions for $\tilde{\omega}$, as in ??). In fact, for $\theta \sim 0.1 - 0.3$, values of $\lambda$ computed through eq. (A6) are of order of several. This is consistent with the idea that the bulk of the cold gas in galactic systems formed at $z > 2 - 3$ and, on average, had not yet been depleted by star formation occurring at roughly constant rate before that time (??, see plot 2 in B01; see also) who used $\theta = 0.2$ and $R = 1$, although results are not too sensitive to $R$. 

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