

The evolution of the Ly α forest effective optical depth following He II reionisation

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ABSTRACT

Three independent observational studies have now detected a narrow ($\Delta z \simeq 0.5$) dip centred at $z = 3.2$ in the otherwise smooth redshift evolution of the Ly α forest effective optical depth. This feature has previously been interpreted as an indirect signature of rapid photo-heating in the IGM during the epoch of He II reionisation. We examine this interpretation using a semi-analytic model of inhomogeneous He II reionisation and high resolution hydrodynamical simulations of the Ly α forest. We instead find that a rapid ($\Delta z \simeq 0.2$) boost to the IGM temperature ($\Delta T \simeq 10^4$ K) beginning at $z = 3.4$ produces a well understood and generic evolution in the Ly α effective optical depth, where a sudden reduction in the opacity is followed by a gradual, monotonic recovery driven largely by adiabatic cooling in the low density IGM. This behaviour is inconsistent with the narrow feature in the observational data. If photo-heating during He II reionisation is instead extended over several redshift units, as recent theoretical studies suggest, then the Ly α opacity will evolve smoothly with redshift. We conclude that the sharp dip observed in the Ly α forest effective optical depth is instead most likely due to a narrow peak in the hydrogen photo-ionisation rate around $z = 3.2$, and suggest that it may arise from the modulation of either reprocessed radiation during He II reionisation, or the opacity of Lyman limit systems.

Key words: methods: numerical - intergalactic medium - quasars: absorption lines - diffuse radiation - cosmology:theory.

1 INTRODUCTION

Observations of the He II Gunn & Peterson (1965) trough in the spectra of intermediate redshift quasars, coupled with large fluctuations in the mean transmission of the He II Ly α forest, provide the most direct evidence for the tail-end of He II reionisation occurring around $z = 3$ (Jakobsen et al. 1994; Davidsen et al. 1996; Heap et al. 2000; Zheng et al. 2004; Shull et al. 2004; Fechner et al. 2006). However, to date only a handful of ultraviolet (UV) quasar spectra suitable for detailed analyses of the He II Ly α forest have been obtained (although see Syphers et al. 2009). Moreover, the Ly α transmission is sensitive to small He II fractions only, and is thus unsuitable for probing the earliest stages of He II reionisation.

Fortunately, the impact of He II reionisation on the intergalactic medium (IGM) can still be probed, albeit indirectly, over a wider redshift baseline with the existing wealth of high quality, high signal-to-noise H I Ly α forest

data. Firstly, the increased temperature of the IGM expected following He II reionisation (Miralda-Escudé & Rees 1994; Abel & Haehnelt 1999; Paschos et al. 2007) will thermally broaden absorption lines in the Ly α forest (Haehnelt & Steinmetz 1998). There is some evidence that observed line widths are consistent with a sudden increase in the IGM temperature around $z = 3.3$ (Schaye et al. 2000; Ricotti et al. 2000), although the error bars on the measurements are large and not all studies agree on this result (McDonald et al. 2001). Secondly, the residual H I fraction in the IGM, which is in photo-ionisation equilibrium with the metagalactic UV background, will be lowered through the temperature dependence of the H II recombination coefficient ($n_{\text{HI}} \propto T^{-0.7}$). If the temperature of the IGM rises suddenly following He II reionisation, a similarly sudden decrease in the H I fraction, and hence the observed Ly α forest opacity, will result (Theuns et al. 2002; Bernardi et al. 2003; Faucher-Giguère et al. 2008b).

Bernardi et al. (2003), in a study using 1061 moderate resolution quasar spectra obtained from the Sloan Digital Sky Survey (SDSS), statistically measured such a departure from the otherwise smooth, power law evolution of the Ly α forest opacity. A sudden decrease of around 10 per cent was observed in the effective optical depth at $z = 3.2$, followed by a recovery to its former power law evolution by $z = 2.9$. Faucher-Giguère et al. (2008b) have subsequently reconfirmed the Bernardi et al. (2003) result by using 86 high resolution, high signal-to-noise spectra obtained with Keck/HIRES/ESI and Magellan/MIKE to directly measure the Ly α opacity. Furthermore, Dall’Aglio et al. (2008) have also recently detected a narrow dip in the Ly α forest effective optical depth at $z = 3.2$, albeit at a low level of statistical significance (2.6σ), using another set of 40 high resolution spectra obtained with VLT/UVES.

Theuns et al. (2002) (hereafter T02) compared detailed hydrodynamical simulations of the Ly α forest to the Bernardi et al. (2003) data and interpreted this feature as indirect evidence for photo-heating during He II reionisation. However, recent theoretical studies indicate it is difficult to explain this narrow feature in the Ly α opacity evolution by invoking a rapid temperature boost in the IGM alone. Bolton et al. (2009) used analytical and numerical arguments to demonstrate that there are too few hard photons available to heat the entire IGM by the required amount ($\sim 10^4$ K) over the timescale on which the observed opacity decreases ($\Delta z = 0.1 - 0.2$). State-of-the-art radiative transfer simulations also indicate that He II reionisation and reheating will be an extended rather than sudden process, with the volume averaged IGM temperature rising gradually from higher redshifts (McQuinn et al. 2009). A rapid injection of energy into the IGM during He II reionisation is thus unlikely to explain the sharp initiation of the observed dip in the Ly α forest opacity.

Somewhat separate to this argument, however, is the explanation for the subsequent *recovery* of the observed Ly α opacity to its former power law evolution by $z = 2.9$. Let us assume that the narrow dip in the Ly α forest effective optical depth is indeed initiated by a large, rapid temperature boost in the IGM following He II reionisation (but see Bolton et al. 2009; McQuinn et al. 2009). Although this will not occur over the whole IGM, such rapid heating is not necessarily excluded in localised patches close to quasars with very hard spectra, where the requisite hard photons are abundant. Following reionisation, the thermal evolution of the IGM at low density is dominated by the balance between photo-heating and adiabatic cooling due to Hubble expansion (Hui & Gnedin 1997; Theuns et al. 1998). However, if the Ly α opacity evolution were driven by the thermal state of the IGM alone through the H II recombination rate, any subsequent recovery in the opacity due to adiabatic cooling would occur over a Hubble time – too long to explain the rapid recovery in the Bernardi et al. (2003) feature. The study of T02 instead demonstrated that this interpretation is too simplistic. They found that the hydrodynamical response of the IGM following a sudden reheating, coupled with the impact of the extra electrons liberated during He II reionisation, are enough to drive the Ly α opacity back to its former power law evolution by $z = 2.9$ as observed by Bernardi et al. (2003) – much earlier than one would expect using the simple argument above.

To date, the only dedicated study of the impact of He II reionisation on the Ly α forest effective optical depth using high resolution hydrodynamical simulations of the IGM has been performed by T02. McQuinn et al. (2009) recently presented a detailed analysis of the evolution of the Ly α forest opacity using radiative transfer simulations of He II reionisation. However, their study necessarily used low resolution, post-processed dark matter simulations, and they were thus unable to model the impact of the gas hydrodynamics on the recovery of the Ly α forest opacity. In light of these recent observational and theoretical results we therefore re-examine the impact of He II reionisation on the evolution of the Ly α forest opacity using semi-analytic modelling and high resolution hydrodynamical simulations of the IGM. In particular, we shall focus on the explanation for the *recovery* of the dip observed in the Ly α forest opacity. This paper is therefore closely related to the work presented in Bolton et al. (2009), where the issues surrounding the *initiation* of the feature first detected by Bernardi et al. (2003) (*i.e.* the plausibility, or lack of, for the rapid photo-heating of the IGM) were examined in detail.

The structure of this paper is as follows. We begin in §2 with a brief review of the relationship between the Ly α forest opacity and the underlying physical properties of the IGM. In §3 we use a semi-analytic model to examine the evolution of the Ly α forest opacity following photo-heating in the IGM during inhomogeneous He II reionisation. Motivated by these results, we then proceed to model the Ly α forest opacity in more detail using hydrodynamical simulations of the IGM. The simulations are described in §4, and the evolution of the Ly α forest opacity in the simulations is presented in §5. In §6 we investigate the impact of sudden reheating on the IGM gas distribution and peculiar velocity field; we find that our simulations are unable to reproduce the narrow feature observed in the Ly α forest opacity evolution. Finally, we consider alternative explanations for the observational data in §7 before concluding in §8.

2 THE FLUCTUATING GUNN-PETERSON APPROXIMATION

The forest of Ly α absorption lines observed in the spectra of high redshift quasars originates from the neutral hydrogen remaining in the intervening, low density IGM following H I reionisation (Bi et al. 1992; Zhang et al. 1995; Hernquist et al. 1996; Miralda-Escudé et al. 1996; Theuns et al. 1998). The mean normalised flux of the Ly α forest, $\langle F \rangle = \langle I_{\text{observed}}/I_{\text{emitted}} \rangle$, is the simplest observable quantity, and it is often expressed as an effective optical depth

$$\tau_{\text{eff}} = -\ln\langle F \rangle \equiv -\ln\langle e^{-\tau} \rangle, \quad (1)$$

where τ is the underlying Ly α optical depth in each pixel of the spectrum or spectra from which $\langle F \rangle$ is measured. Assuming the IGM is highly ionised and in photo-ionisation equilibrium with the metagalactic UV background, and the low density IGM ($\Delta = \rho/\langle \rho \rangle \leq 10$) follows a power-law temperature density relation, $T = T_0 \Delta^{\gamma-1}$ (Hui & Gnedin 1997; Valageas et al. 2002), the Ly α optical depth at $z \gtrsim 2$ may be written as (*e.g.* Weinberg et al. 1999; McDonald & Miralda-Escudé 2001)

$$\tau \simeq 1.0 \frac{(1 + \chi_{\text{He}})}{\Gamma_{-12}} \left(\frac{T_0}{10^4 \text{ K}} \right)^{-0.7} \left(\frac{\Omega_b h^2}{0.024} \right)^2 \left(\frac{\Omega_m h^2}{0.135} \right)^{-1/2} \times \left(\frac{1+z}{4} \right)^{9/2} \Delta^{2-0.7(\gamma-1)}, \quad (2)$$

where Ω_b and Ω_m are the present day baryon and matter densities as a fraction of the critical density, $h = H_0/100 \text{ km s}^{-1} \text{ Mpc}^{-1}$ for the present day Hubble constant H_0 , $\Delta = \rho/\langle\rho\rangle$ is the normalised gas density, T_0 is the gas temperature at mean density, γ is the slope of the temperature density relation and $\Gamma_{-12} = \Gamma_{\text{HI}}/10^{-12} \text{ s}^{-1}$ is the hydrogen photo-ionisation rate. The power-law temperature dependence is due to the case-A H II recombination coefficient, and χ_{He} accounts for the extra electrons liberated during He II reionisation; $\chi_{\text{He}} = 1.08$ prior to He II reionisation and $\chi_{\text{He}} = 1.16$ afterwards for a helium fraction by mass of $Y = 0.24$ (Olive & Skillman 2004). The effective optical depth can then be estimated by integrating over all possible IGM densities Δ .

Eq. (2) is the Fluctuating Gunn-Peterson Approximation (FGPA), and although it ignores the effect of redshift space distortions on the Ly α forest opacity, it clearly elucidates the relationship between the opacity and the underlying physical properties of the IGM. A sudden decrease in the Ly α effective optical depth (*e.g.* Bernardi et al. 2003) can thus be attributed to an increase in the IGM temperature, either by raising T_0 or changing¹ γ , an increase in the photo-ionisation rate, or a combination of both. On the other hand, an increase in the free electron fraction will raise the opacity by reducing the recombination timescale. Both an increase in the IGM temperature and an 8 per cent change in the free electron fraction will be associated with He II reionisation. Keeping these points in mind, we now proceed to describe a semi-analytical model for inhomogeneous He II reionisation which utilises this useful approximation.

3 THE Ly α FOREST OPACITY DURING INHOMOGENEOUS He II REIONISATION

3.1 A semi-analytic model

Before analysing our detailed hydrodynamical simulations, we first consider a simplified semi-analytic model which illustrates the difficulty of reproducing the localised feature observed by Bernardi et al. (2003) in the Ly α forest. Although this model is unable to compute many of the detailed properties of the Ly α forest (such as line broadening and the

¹ Lowering (raising) the value of γ will increase the temperature in underdense (overdense) regions of the IGM while decreasing the temperature in overdense (underdense) regions. Whether raising or lowering γ subsequently decreases the effective optical depth of the Ly α forest thus depends on redshift. At $z \geq 3$ the Ly α forest opacity is mainly sensitive to underdense regions in the IGM. Lowering γ therefore decreases the Ly α effective optical depth at $z \geq 3$ by producing hotter, more highly ionised voids (*e.g.* Bolton et al. 2005; Faucher-Giguère et al. 2008a). Typical values adopted for γ in Ly α forest models lie in the range $1 \leq \gamma \leq 1.6$, although there is some evidence to suggest a more complex relationship between temperature and density, perhaps with $\gamma < 1$, may be required to reproduce the observed Ly α forest flux distribution (Becker et al. 2007; Bolton et al. 2008).

impact of peculiar velocities), it has the virtue of including inhomogeneous reionisation (Furlanetto & Oh 2008b), which our hydrodynamical simulations do not. The model is based on the calculations presented in Furlanetto & Oh (2008a), who examined how He II reionisation can affect the temperature-density relation of the IGM, and we refer the reader there for more details.

In brief, the model has three parts. The first determines the reionisation history of gas elements of a given density in the IGM (Furlanetto & Oh 2008a). For most of our calculations, we assume that the probability that an element is reionised at any particular time is independent of its density, simply tracing the mean ionisation history, $\bar{x}_i(z)$. We also compare to a “density-driven” model in which high-density regions are ionised first (because they sit near the biased regions that host quasars). However, recent numerical simulations of He II reionisation suggest that the ionised regions are large and rare enough that this correlation is weak (McQuinn et al. 2009), so we usually use the density-independent model. We will assume that He II reionisation ends at $z_{\text{He}} = 3.2$, consistent with the calculations below, and that hydrogen reionisation (which only affects the temperature of gas for which helium is still singly-ionised) occurs at $z_{\text{H}} = 8$; the latter has little effect on our calculation.

To make this comparison as straightforward as possible, we assume that $\bar{x}_i(z)$ is simply proportional to the mass in galaxies with $m \gtrsim 5 \times 10^{11} M_{\odot}$ (in other words, these massive galaxies contain supermassive black holes that have gone through bright quasar phases). This provides a reasonable approximation to the quasar emissivity (compared to, *e.g.*, the luminosity function of Hopkins et al. 2007) and also leads to fast He II reionisation, with over 70% of the ionisation occurring after $z = 4.2$, if He II reionisation ends at $z = 3.2$ (Furlanetto & Oh 2008b).

The second ingredient is to follow the thermal evolution of each gas element after reionisation. We again use the method presented in Furlanetto & Oh (2008a), which is in turn based on Hui & Gnedin (1997). We include all of the relevant atomic cooling, heating, recombination, and photo-ionisation processes, as well as adiabatic expansion (and collapse for overdense regions). After reionisation, the dominant processes are photo-heating and adiabatic cooling, which together determine the asymptotic, nearly power law, temperature-density relation (Hui & Gnedin 1997; Hui & Haiman 2003).

The chief uncertainty in this model is the initial temperature after reionisation, T_i , which depends in a non-trivial manner on the spectrum of the photons ionising the gas parcel (Abel & Haehnelt 1999; Tittley & Meiksin 2007; Paschos et al. 2007; Bolton et al. 2009; McQuinn et al. 2009). This may itself be inhomogeneous in the IGM, because low-energy photons will be absorbed near their hosts, leaving many of the ionisations to be done by hard photons (Abel & Haehnelt 1999). We therefore consider a range of possibilities for T_i . Our fiducial model takes $T_i = 4 \times 10^4 \text{ K}$, larger than is typically expected in order to exaggerate the effect of heating (*e.g.*, Bolton et al. 2009; McQuinn et al. 2009).

Finally, in order to compute the average optical depth we need to assume an IGM density distribution, $P_V(\Delta)$ (averaged by volume), as well as the distribution of temperatures and ionisation rates at each density. We will use the

density distribution recommended by Miralda-Escudé et al. (2000):

$$P_V(\Delta) d\Delta = A_0 \Delta^{-\beta} \exp \left[-\frac{(\Delta^{-2/3} - C_0)^2}{2(\delta_0/3)^2} \right] d\Delta. \quad (3)$$

This form fits cosmological simulations at $z = 2-4$ quite well. Note that the underlying simulation had somewhat different cosmological parameters than the currently preferred values; it is, however, accurate enough for the qualitative calculations that follow.

3.2 The Ly α effective optical depth

To compute the Ly α effective optical depth, we assign each volume element of known density an optical depth according to Eq. (2)². Our model makes no predictions about the amplitude of the ionising background, and for simplicity we set $\Gamma_{-12} = 1$ over the entire redshift range $z = 2-5$. This is consistent with constraints derived by comparing observational data to simulations of the Ly α forest at $2 \leq z \leq 4$ (Bolton et al. 2005), and yields a mean transmission reasonably close to the observed values.

We choose the temperatures by following gas elements of the appropriate density after their most recent reionisation event using the thermal evolution code described above, with the reionisation redshifts distributed according to the overall ionisation history. We then calculate

$$e^{-\tau_{\text{eff}}} = \int_0^\infty d\Delta P_V(\Delta) \int dT P(T|\Delta) e^{-\tau(T,\Delta)}, \quad (4)$$

where $P(T|\Delta)$ is the probability distribution of temperatures for elements at density Δ and τ_{eff} is the effective optical depth. Again, we emphasise that this ignores peculiar velocities of the gas, line blending (and indeed the wings of every line), and the clustering of the absorbers, but it provides a qualitative description of the evolving transmission.

Fig. 1 shows some example histories computed in this manner. As a basis for comparison, the dotted curve ignores He II reionisation and assumes a constant ionising background and IGM temperature ($T = 20,000$ K). This does not quantitatively match the observed evolution; it merely serves to show that, without He II reionisation, our model produces smooth, nearly featureless evolution.

The other curves assume that helium reionisation ends at $z_{\text{He}} = 3.2$. The short-dashed, solid, and dot-dashed curves assume $T_i = (3, 4, 6) \times 10^4$ K, respectively (all relatively large in order to exaggerate the effect; see Bolton et al. 2009; McQuinn et al. 2009 for self-consistent estimates of the plausible temperature range). The long-dashed curve takes $T_i = 4 \times 10^4$ K and assumes that reionisation is density-dependent (Furlanetto & Oh 2008a); this aspect of reionisation has only a small effect on the mean transmission. This is because the Ly α forest is sensitive to only a narrow range of densities near the mean, where the density modulation is modest anyway.

According to these models, He II reionisation can certainly induce a feature in τ_{eff} , so long as the temperature

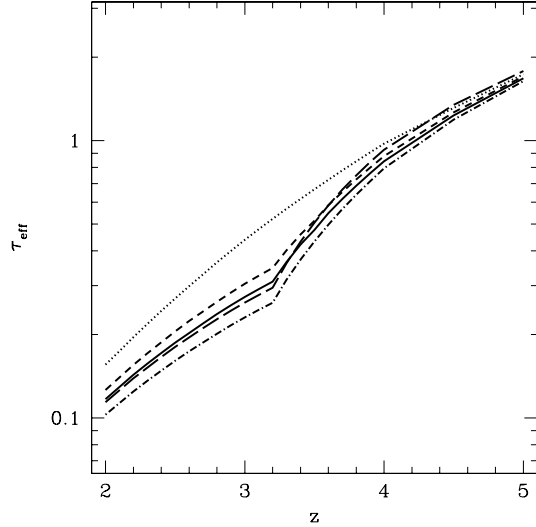


Figure 1. Mean transmission histories in our semi-analytic models. The dotted curve shows $\tau_{\text{eff}}(z)$ if we ignore He II reionisation and set $T = 20,000$ K throughout. The other curves include He II reionisation at $z_{\text{He}} = 3.2$, as described in the text. The short-dashed, solid, and dot-dashed curves assume $T_i = (3, 4, 6) \times 10^4$ K, respectively. The long-dashed curve assumes $T_i = 4 \times 10^4$ K but that reionisation proceeds from high to low densities.

increase is large enough. In all cases, τ_{eff} falls relatively steeply until z_{He} before levelling off and returning closer to the expected evolution without He II reionisation. The shape is generic within these reionisation models, although note that the downward turn in τ_{eff} is still significantly shallower than that observed in the observational data (Bernardi et al. 2003; Faucher-Giguère et al. 2008b). The pre-reionisation phase can be steepened by making reionisation occur faster, but the model assumed here is not far from empirical estimates of the evolution of the quasar emissivity. Moreover, compressing He II reionisation into a short time interval limits the temperature jump that it can induce, because the photo-ionisation timescale associated with high-energy photons (which provide the most efficient heating) is quite long (Bolton et al. 2009).

However, the post-reionisation evolution is completely generic and easy to understand. In this regime, the thermal evolution is dominated by the competition between photo-heating and cooling by adiabatic expansion. The timescale for the cooling is therefore the expansion, or Hubble, time – a substantial fraction of the age of the Universe. There is no way to avoid this behaviour for the temperature history, and to the extent that the Ly α forest depends only on these temperatures, He II reionisation cannot induce a *narrow* feature in the transmission which recovers quickly to its pre-reionisation evolution.

Of course, we have emphasised that the forest is indeed more complex than this model, because of peculiar velocities, geometric effects, and line blending. In principle, these can induce a narrower feature, and T02 appealed to just such an effect to explain the Bernardi et al. (2003) feature. In their simulations, once the reheated gas in the IGM became overpressurised with respect to its surroundings it started to expand, resulting in a sudden change in the peculiar veloc-

² We drop the assumption of a power law temperature-density relation in this instance; see Furlanetto & Oh (2008a) for some example temperature distributions.

ity gradients in the IGM. This extended the Ly α absorption lines in redshift space by shifting absorption from the saturated line cores to the wings, increasing the line equivalent widths and hence the mean Ly α opacity. In the remainder of this paper, we will examine all of these effects in more detail using new hydrodynamical simulations that do include a detailed reconstruction of the Ly α forest.

4 HYDRODYNAMICAL SIMULATIONS OF THE Ly α FOREST

4.1 Initial conditions

The hydrodynamical simulations in this study are performed using an upgraded version of the publicly available parallel Tree-SPH code GADGET-2 (Springel 2005). All simulations have a box size $15h^{-1}$ comoving Mpc and contain 2×400^3 gas and dark matter particles. The mass of each gas particle is $9.4 \times 10^5 M_\odot$ and the gravitational softening length is $1/30^{\text{th}}$ of the mean linear interparticle spacing. This adequately resolves the Ly α forest at $2 \leq z \leq 4$ (Theuns et al. 1998; Bolton et al. 2008) and provides just over twice the mass resolution of the simulations used in the T02 study.³ Star formation is included using a simplified prescription which converts all gas particles with overdensity $\Delta > 10^3$ and temperature $T < 10^5$ K into collisionless stars, significantly speeding up the simulations. Outputs are saved every $\Delta z = 0.05$ in the redshift range $2 \leq z \leq 4$, enabling a very fine sampling of the simulation data with redshift.

The ionisation state of the gas particles is computed in the optically thin limit using a non-equilibrium ionisation algorithm which follows the abundances of six species: H I, H II, He I, He II, He III and e^- (Bolton & Haehnelt 2007a). The ultraviolet background (UVB) model of Haardt & Madau (2001) (hereafter HM01) including emission from quasars and galaxies is used to compute the H I, He I and He II photo-ionisation and heating rates. This model is in good agreement with observational constraints on the metagalactic H I photo-ionisation rate (Tytler et al. 2004; Bolton et al. 2005), although we set the HM01 photo-ionisation and heating rates for all species to zero at $z > 6$ in order to match the simulations of T02.

The UVB is assumed to be spatially uniform in all our simulations. This will be a reasonable approximation at the H I photo-ionisation edge at $z < 4$, when the mean free path for hydrogen ionising photons is much larger than the average separation between ionising sources (Bolton & Haehnelt 2007b; Faucher-Giguère et al. 2008a; Furlanetto 2008). However, a full radiative transfer implementation is required to model the impact of inhomogeneous He II reionisation on the Ly α forest opacity (e.g. Maselli & Ferrara 2005; Tittley & Meiksin 2007; Paschos et al. 2007; Bolton et al. 2009; McQuinn et al. 2009). Although He II reionisation does not *directly* impact Ly α forest effective optical depth, our simulations do not capture the effect of the resulting fluctuations in the IGM temperature, which do impact on the effective optical depth through the temperature dependence of

the H II recombination coefficient. However, our simulation volume is $15h^{-1}$ comoving Mpc, which is small on comparison to the ~ 50 comoving Mpc scales on which large temperature fluctuations are expected (McQuinn et al. 2009). Although not ideal, our assumption of instantaneous reheating on this smaller scale is therefore a reasonable approximation. Furthermore, note that HM01 implicitly assumes that He II is reionised at all redshifts, *i.e.* He II only exists in the dense Ly α absorbers which they model. Radiative transfer effects obviously modify the form of the UV background prior to He II reionisation (e.g. Madau & Haardt 2009), both in He II ionising radiation above 4 Ry and in reprocessed radiation from He II Lyman series and two photon emission. Thus, the UV background could be quite different before reionisation is complete. We comment on this more in §7.2.

The simulations were all started at $z = 99$, with initial conditions generated using the transfer function of Eisenstein & Hu (1999). The cosmological parameters adopted are $\Omega_m = 0.26$, $\Omega_\Lambda = 0.74$, $\Omega_b h^2 = 0.024$, $h = 0.72$, $\sigma_8 = 0.85$ and $n_s = 0.95$. These are consistent with the fifth year *Wilkinson Microwave Anisotropy Probe* (WMAP) data (Dunkley et al. 2009), aside from a slightly larger value for the matter power spectrum normalisation. The gas is assumed to be of primordial composition with a helium mass fraction of $Y = 0.24$ (e.g. Olive & Skillman 2004).

4.2 Thermal and ionisation histories

Five hydrodynamical simulations, listed in Table 1, were performed for this study. The simulations use different models for the thermal and ionisation evolution of the IGM, but aside from the S3 model, which we shall discuss later, are identical in all other respects. The different thermal histories are constructed by increasing the HM01 He II photo-heating rate in the simulations (e.g. T02; Bolton et al. 2005; Jena et al. 2005). This mimics spectral hardening due to radiative transfer effects during He II reionisation by boosting the mean excess energy per He II photo-ionisation (Abel & Haehnelt 1999; Bolton et al. 2004). Note that we only model the jump in heating rates, which affects the thermal evolution; apart from model S2 which we discuss below, we do not model the jump in the He II fraction itself.

The IGM temperatures at mean density, T_0 , in all five simulations are displayed in the left panel of Fig. 2 as a function of redshift. These data are compared to observational constraints obtained by Schaye et al. (2000) from an analysis of absorption line widths in the Ly α forest. Note that these thermal histories are not fits to the observational data; they are instead merely chosen to be representative of IGM thermal histories presented in the literature. The S1 model (solid curve) closely resembles the T_0 evolution in the simulations of T02, with a sudden increase ($\Delta z = 0.1$) in temperature, $\Delta T \sim 10^4$ K, beginning at $z = 3.4$. Note that this temperature boost occurs over a much shorter timescale than that in our semi-analytical model of He II reionisation (Furlanetto & Oh 2008a). The E1 model (dot-dashed curve) is qualitatively similar to recent results from detailed three dimensional radiative transfer simulations of He II reionisation (McQuinn et al. 2009). The temperature boost develops over a longer timescale as quasars gradually photo-heat the IGM. The third simula-

³ T02 use a modified version of Hydra (Couchman et al. 1995) to run hydrodynamical simulations in a $12h^{-1}$ comoving Mpc box with a gas particle mass of $2.0 \times 10^6 M_\odot$.

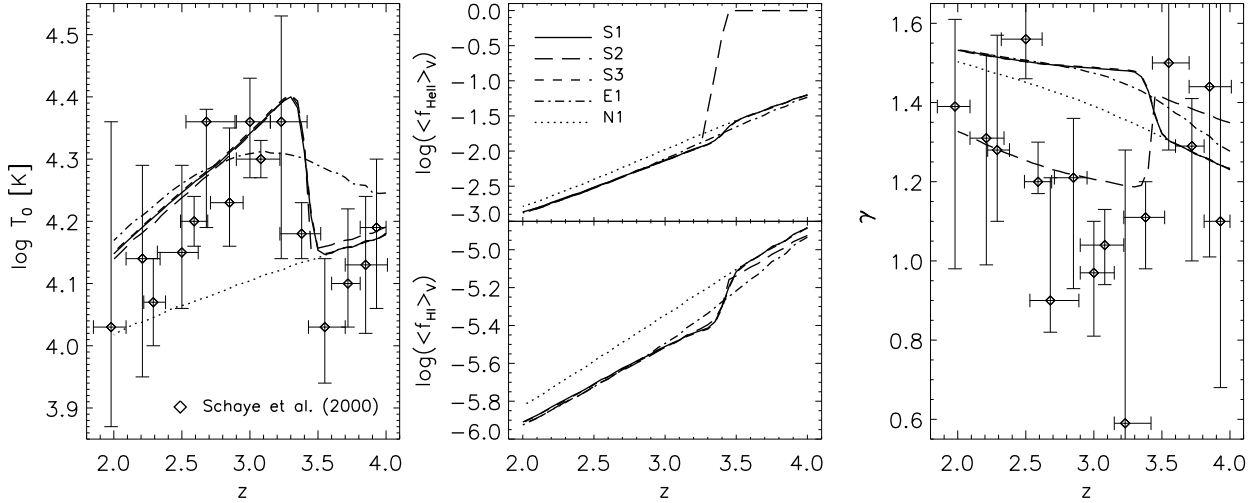


Figure 2. *Left:* The temperature of the IGM at mean density, T_0 , in the five simulations used in this study: S1 (solid curve), S2 (long-dashed curve), S3 (short-dashed curve), E1 (dot-dashed curve) and N1 (dotted curve). The temperatures are given at intervals of $\Delta z = 0.05$ over the redshift range $2 \leq z \leq 4$. These data are compared to observational constraints obtained by Schaye et al. (2000) from an analysis of absorption line widths in the Ly α forest. *Centre:* The corresponding volume weighted He II (upper panel) and H I (lower panel) fractions. *Right:* The slope of the power law temperature-density relation, $T = T_0 \Delta^{\gamma-1}$. The values are obtained by a least squares power law fit to the volume weighted temperature-density plane in the simulations at $0.1 \leq \Delta \leq 1$ and $T \leq 10^5$ K. The data are again compared to observational constraints from Schaye et al. (2000). Note that the S1 and S3 models are almost indistinguishable in all three panels.

Table 1. Hydrodynamical simulations used in this study. All the simulations have a box size of $15h^{-1}$ comoving Mpc and contain 2×400^3 gas and dark matter particles.

Model	Thermal history description
S1	Sharp temperature boost
E1	Extended temperature boost
N1	No temperature boost; control model
S2	Similar to S1, but with a rapid change in n_e
S3	Identical to S1, but with a stricter timestep limit

tion, N1 (dotted curve), is included as a control model, and has no temperature boost.

The corresponding volume weighted He II and H I fractions in the simulations are displayed in the central panel of Fig. 2. In the S1, E1 and N1 simulations, both H I and He II reionisation commence at $z = 6$. This choice is deliberate; these three models are designed to exclude the 8 per cent increase in the free electron fraction following He II reionisation (see discussion in §2). Instead, they will be used to explore the impact of differences in the IGM thermal state *alone* on the τ_{eff} evolution. Sudden He II reionisation at $z = 3.4$ is instead included in a fourth model, S2 (long-dashed curve), which is tailored to have a similar temperature at mean density to the S1 model. The S2 simulation is therefore the most similar to the model used by T02, who also assumed H I reionisation at $z = 6$ and He II reionisation at $z = 3.4$.

The right hand panel in Fig. 2 displays the evolution of the slope of the temperature-density relation, $T = T_0 \Delta^{\gamma-1}$ (Hui & Gnedin 1997; Valageas et al. 2002), in each of the simulations. The power-law index γ gradually increases towards lower redshifts in the E1 and S1 models, asymptot-

ically approaching the upper limit of $\gamma \sim 1.6$ achieved by the balance between photo-heating and adiabatic cooling in the low density IGM. Note however, the S1 and S2 models, although having very similar T_0 values, exhibit different behaviour for γ . This is due to the different ionisation histories adopted in the two models. During the reionisation of He II at $z = 3.4$ in the S2 model, photo-ionisation equilibrium is no longer a good assumption and the He II photo-heating rate is *independent* of density (Bolton et al. 2009). This is because the He II fraction $-f_{\text{HeII}} = n_{\text{HeII}}/n_{\text{He}} \sim 1$ immediately prior to reionisation – is independent of density. This flattens the temperature-density relation and lowers the value of γ . However, in the S1 model, where the He II is already in ionisation equilibrium with the UV background, the He II photo-heating rate is instead *proportional* to the IGM density (Theuns 2005). The He II fraction is proportional to density (due to higher recombination rates in denser regions) and the extra energy injected at $z = 3.4$ instead increases γ . As we shall see, these differences will also play a small role (relative to T_0) in the simulated τ_{eff} evolution. These values are compared to measurements of γ from Schaye et al. (2000). Although the error bars are large, somewhat lower values of γ are preferred at $z < 3.5$.

4.3 Time integration in GADGET-2 and Hydra

We run the fifth and final model, S3, to quantify the effect of the GADGET-2 time integration scheme on our results. The study of T02 used hydrodynamical simulations performed with a modified version of the P³M – SPH code Hydra (Couchman et al. 1995). Hydra employs a single step time integration scheme where all particle positions are advanced on the minimum timestep required throughout the simulation volume. In contrast, GADGET-2 uses individual

timesteps for each particle. This means that particles in high density regions, where dynamical timescales are short, have timesteps which are orders of magnitude smaller than particle timesteps in the lowest density regions (Springel 2005). This enables efficient use of computational resources and substantially improves code performance. However, if a sudden boost to the IGM temperature occurs in the middle of a particle timestep, the energy injected will be smeared over the timescale corresponding the particle timestep. In regions where the dynamical timescale is long, such as the low density IGM which dominates the transmission in the Ly α forest at $z = 3$, this can delay the impact photo-heating on the gas and could potentially affect our numerical results.

Unfortunately, forcing a single minimum timestep in GADGET-2 to test this possibility would be prohibitively expensive. We instead impose a stricter upper limit on the maximum timestep size for *all* particles in the S3 model. In our first four simulations, the maximum allowed timestep for all particles is $\Delta \ln[1/(1+z)] = 0.05$, although in practice this value will vary and can be up to a factor of two smaller at any given redshift. This is equivalent to specifying the maximum allowed timestep as a fraction of the current Hubble time in the simulations. In the S3 model we instead impose a maximum timestep of $\Delta \ln[1/(1+z)] = 0.001$ at $z < 4.2$. This choice should provide a useful test of the impact of timestepping on the gas hydrodynamics during a rapid change in the IGM temperature (V. Springel, private communication). The thermal and ionisation history of the S3 model corresponds to the short-dashed curves in Fig. 2. Note the differences between the thermal and ionisation histories of the S1 and S3 models are minimal, and that the curves are almost indistinguishable in all three panels.

5 THE Ly α EFFECTIVE OPTICAL DEPTH FROM HYDRODYNAMICAL SIMULATIONS

5.1 Construction of synthetic Ly α spectra

We now turn to analysing the Ly α forest effective optical depth in each of our simulations. Synthetic Ly α forest spectra are constructed from each simulation by extracting randomly selected sight-lines parallel to the simulation box boundaries at redshift intervals of $\Delta z = 0.05$ over the range $2 \leq z \leq 4$. At each redshift a total of 1024 sight-lines are extracted, each with 1024 pixels. Every pixel in each sight-line has a neutral hydrogen number density n_{HI} , temperature T , peculiar velocity v_{pec} and Hubble velocity v_{H} associated with it. A standard SPH interpolation procedure (*e.g.* Theuns et al. 1998) is used to extract the first three of these quantities from the simulation data. In each line of sight with N pixels, the Ly α optical depth in pixel i is then given by

$$\tau(i) = \frac{c\sigma_{\alpha}\delta R}{\pi^{1/2}} \sum_{j=1}^N \frac{n_{\text{HI}}(j)}{b_{\text{HI}}(j)} H(a, x). \quad (5)$$

Here $b_{\text{HI}} = (2k_{\text{B}}T/m_{\text{H}})^{1/2}$ is the Doppler parameter, $\sigma_{\alpha} = 4.48 \times 10^{-18} \text{ cm}^2$ is the Ly α cross-section, δR is the pixel width and $H(a, x)$ is the Voigt-Hjerting function (Hjerting 1938)

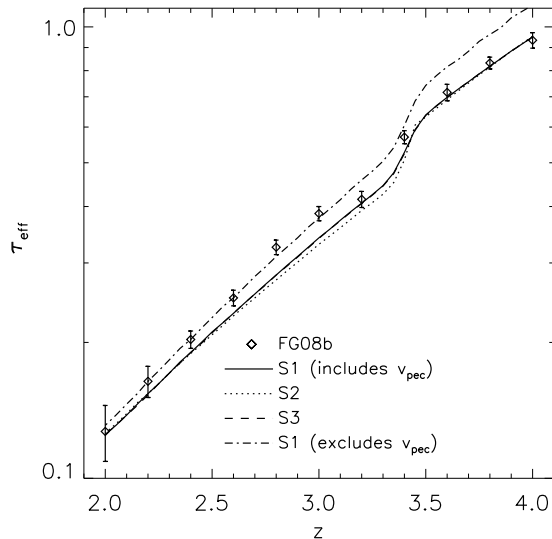


Figure 3. Comparison of τ_{eff} measured from our synthetic Ly α forest spectra to the observational data of FG08b (open diamonds). The error bars correspond to the statistical uncertainties only. The synthetic data correspond to the models with a sharp temperature boost at $z = 3.4$ (S1, solid curve), an 8 per cent increase in the free electron fraction in addition to the sharp temperature boost (S2, dotted curve) and a stricter upper limit on the particle timesteps (S3, dashed curve). The latter model is indistinguishable from the S1 data. A fourth curve (dot-dashed) corresponds to the effective optical depth measured from the S1 model ignoring the effect of peculiar velocities on the Ly α forest. As noted by T02, neglecting peculiar velocities raises the absolute value of τ_{eff} by deblending absorption lines.

$$H(a, x) = \frac{a}{\pi} \int_{-\infty}^{\infty} \frac{e^{-y^2}}{a^2 + (x-y)^2} dy, \quad (6)$$

where $x = [v_{\text{H}}(i) - u(j)]/b_{\text{HI}}(j)$, $u(j) = v_{\text{H}}(j) + v_{\text{pec}}(j)$, $a = \Lambda_{\alpha}\lambda_{\alpha}/4\pi b_{\text{HI}}(j)$, $\Lambda_{\alpha} = 6.265 \times 10^8 \text{ s}^{-1}$ is the damping constant and $\lambda_{\alpha} = 1215.67 \text{ \AA}$ is the wavelength of the Ly α transition. We use the analytic approximation for $H(a, x)$ provided by Tepper García (2006).

5.2 Comparison to the τ_{eff} evolution observed by FG08b

We shall compare our simulation data to the recent direct measurement presented by Faucher-Giguère et al. (2008b) (hereafter FG08b), who found a feature in the Ly α forest opacity evolution consistent with the one detected statistically by Bernardi et al. (2003). FG08b measured τ_{eff} over the redshift range $2 \leq z \leq 4.2$ using a sample of 86 high-resolution, high signal-to-noise quasar spectra obtained using three different instruments; the ESI and HIRES spectrographs on Keck and the MIKE spectrograph on Magellan.

We firstly renormalise the optical depths of the synthetic spectra by the same constant at every redshift, $A = 0.87$, to approximately match the normalisation of the Ly α effective optical depth measured by FG08b. For a set of spectra with N pixels

$$e^{-\tau_{\text{eff}}} = \langle F \rangle = \frac{1}{N} \sum_{i=1}^N e^{-A\tau_i}, \quad (7)$$

where τ_i is the optical depth in each pixel of the synthetic spectra and $\langle F \rangle$ is the mean observed flux. From Eq. (2), this renormalisation corresponds to a straightforward rescaling of Γ_{-12} upwards by around 15 per cent if all other parameters in the simulation remain fixed. We have verified this renormalisation has no impact on the shape of the τ_{eff} evolution.

Fig. 3 shows the comparison between the τ_{eff} evolution measured from our synthetic spectra and the observational data of FG08b. Note we have yet not processed the synthetic spectra to resemble the FG08b data (aside from the overall renormalisation); the synthetic spectra are noiseless and have a substantially longer path length than the observational data, but nevertheless provide a useful illustration of the τ_{eff} evolution in the simulations. The solid, dotted and dashed curves display the S1, S2 and S3 models, respectively. These data all display the same generic features observed in the semi-analytic model; a downturn in τ_{eff} due to the reheating of the IGM, followed by a slow recovery driven mainly by adiabatic cooling. In all instances it is clear the simulated data do not exhibit the narrow feature observed in the FG08b data.

Firstly, note the S3 and S1 models are again almost indistinguishable, giving us confidence that the adaptive time integration scheme used in GADGET-2 will not significantly alter our results. Secondly, a small difference between the τ_{eff} evolution observed in the S1 and S2 models is apparent at $z < 3.4$ – recall that the S2 simulation includes the 8 per cent increase in the free electron abundance due to sudden He II reionisation. As noted in §2, the extra electrons will reduce the H II recombination timescale and thus increase the Ly α opacity. In a highly ionised IGM the equilibration timescale is $t_{\text{eq}} \simeq f_{\text{HI}} t_{\text{rec}}$, where $f_{\text{HI}} = n_{\text{HI}}/n_{\text{H}}$ is the IGM H I fraction and the recombination timescale is

$$t_{\text{rec}} \simeq \frac{9.4 \times 10^9 \text{ yrs}}{(1 + \chi_{\text{He}})\Delta} \left(\frac{T}{2 \times 10^4 \text{ K}} \right)^{0.7} \left(\frac{1+z}{4} \right)^{-3}. \quad (8)$$

Therefore, for $f_{\text{HI}} \simeq 10^{-5.5}$, the equilibration timescale is already very short at mean density, $t_{\text{eq}} \sim 3 \times 10^4$ yrs (*cf.* 2.9×10^8 yrs between $z = 3.4$ and $z = 3$), and the increase in f_{HI} (and hence τ_{eff}) following the additional 8 per cent increase in the electron number density should be very rapid. However, we instead find τ_{eff} in the S2 model is 3-4 per cent *lower* than the S1 model at $3 \leq z \leq 3.4$. In this instance the effect of the extra electrons on the Ly α opacity is countered by the associated flattening of the temperature-density relation following sudden, homogeneous He II reionisation. Recall that most of the Ly α forest at $z > 3$ is dominated by transmission from underdense regions in the IGM, which become hotter in the S2 model as the temperature-density relation flattens. Consequently, including a sudden change in the ionisation state of He II in addition to a large temperature boost *increases* rather than decreases the magnitude of the dip in τ_{eff} at $z < 3.4$. The extra electrons therefore do not help in reproducing the τ_{eff} data in our model.

The fourth, dot-dashed curve corresponds to spectra extracted from the S1 simulation, but this time excluding the effect of redshift space distortions induced by peculiar velocities in the IGM. As noted by T02, excluding the peculiar

Table 2. Instrumental resolution (FWHM), pixel sizes and signal-to-noise properties adopted for the synthetic Ly α forest spectra. The values are based on those reported by FG08b.

Instrument	FWHM	Pixel size	S/N per pixel
HIRES	6 km s ⁻¹	2 km s ⁻¹	30
ESI	33 km s ⁻¹	11 km s ⁻¹	30
MIKE	11 km s ⁻¹	2 km s ⁻¹	20

velocities tends to deblend Ly α absorption lines, increasing τ_{eff} . It is also apparent the peculiar velocity field flattens the τ_{eff} evolution somewhat; the gradient of the dot-dashed curve is steeper in comparison to the S1 data including peculiar velocities, especially at $z < 3$. These differences decrease towards lower redshift, indicating line blending becomes less widespread as the line number density and Ly α opacity both fall. However, in contrast to T02 (see their fig. 2), we find these redshift space distortions are insufficient to produce the rapid upturn in τ_{eff} by $z = 2.9$ seen in the observational data. The redshift space distortions induced by peculiar velocities provide no change in the generic shape of the τ_{eff} evolution with redshift. We have verified that this also holds for the S3 model with the stricter timestep limit. This result will be considered in more detail in §6.

However, to make a fair comparison to the FG08b data, our synthetic spectra must be processed to resemble their observational data set as closely as possible. Our renormalised synthetic spectra are therefore also convolved with a Gaussian instrument profile and resampled onto pixels of the required size. Gaussian distributed random noise is then added to each pixel. The parameters used for this procedure are summarised in Table 2, corresponding to the three instruments used in the FG08b data set. A random subsample of synthetic HIRES, ESI and MIKE spectra with a total path length corresponding to the values displayed in fig. 2 of FG08b are then drawn from our simulated data set in intervals of $\Delta z = 0.2$ between $2 \leq z \leq 4$.

The results are displayed as the open diamonds in Fig. 4 for the S1 (left panel), E1 (central panel) and N1 (right panel) models. Following FG08b, the error bars (statistical only) correspond to the standard error of the mean, computed by subdividing the synthetic data into chunks 3 proper Mpc in length. The solid curves in each panel correspond to the underlying τ_{eff} evolution measured from the unprocessed synthetic Ly α spectra. Our results are compared to the best fitting curve to the FG08b data, displayed by the dashed line in each panel.

As expected, the synthetic data in Fig. 4 now exhibit some additional scatter due to the variation in cosmic structure probed by the random sight-lines, but in all instances are within 1σ of the underlying τ_{eff} evolution. The S1 data in the left-hand panel again exhibit the generic behaviour seen in the semi-analytical model, and are inconsistent with the narrow feature observed by FG08b. For comparison, the central panel displays the data obtained from the E1 model with an extended period of heating (*e.g.* McQuinn et al. 2009), while the right hand panel displays the N1 model with no additional heating from He II reionisation. In both instances τ_{eff} smoothly evolves with redshift. This result is in agreement with the predictions for τ_{eff} from the radiative

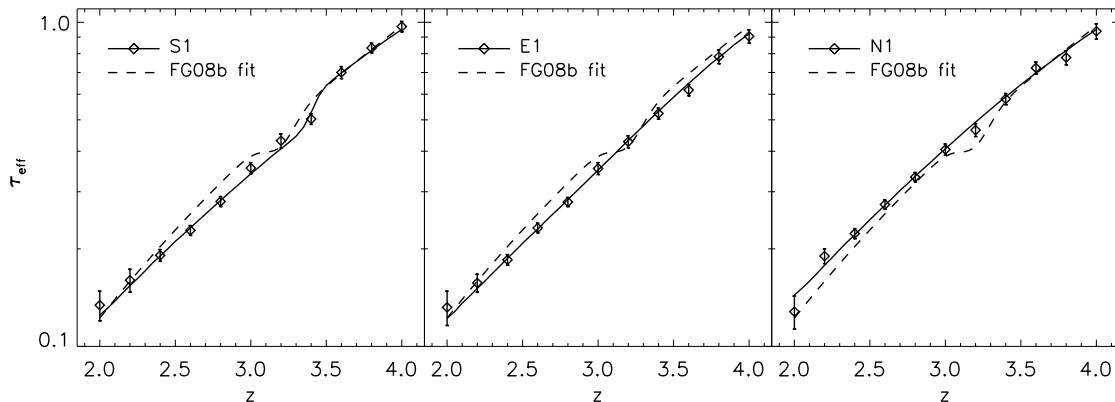


Figure 4. A comparison of the τ_{eff} evolution measured from synthetic Ly α spectra drawn from three of our hydrodynamical simulations to the best fitting curve from FG08b (dashed curves). The solid curves in each panel display τ_{eff} measured from the unprocessed synthetic Ly α spectra, while the open diamonds with error bars correspond to τ_{eff} measured in bins of width $\Delta z = 0.2$ after the spectra are processed to closely resemble the FG08b data set. The error bars correspond to the standard error of the mean. *Left:* The S1 model. *Centre:* The E1 model. *Right:* The N1 model.

transfer simulations of McQuinn et al. (2009); an extended period of reheating during He II reionisation will induce no sharp features in the τ_{eff} evolution due to changes in the IGM temperature alone.

Our detailed simulations therefore indicate that hydrodynamical effects in the IGM following rapid reheating (but see Bolton et al. 2009; McQuinn et al. 2009) will not aid in reproducing the sharp dip in the Ly α opacity observed by FG08b. Instead, our results suggest any recovery in τ_{eff} following sudden reheating will be driven primarily by adiabatic cooling in the low density IGM, consistent with the behaviour predicted in our semi-analytical model.

6 HYDRODYNAMICS AND REDSHIFT SPACE DISTORTIONS FOLLOWING SUDDEN REHEATING

In the last section we demonstrated our simulations were unable to reproduce the narrow feature observed by FG08b. We now consider this result in more detail by examining the impact of hydrodynamical effects and redshift space distortions on the Ly α effective optical depth evolution.

6.1 The gas density distribution

We begin by briefly discussing the effect of sudden reheating on the IGM density distribution. Following an injection of energy into the IGM during reionisation, the subsequent increase in gas pressure will smooth the IGM density distribution over scales corresponding to the local Jeans length (Schaye 2001; Pawlik et al. 2009). However, as noted by Gnedin & Hui (1998), this Jeans smoothing will not occur instantaneously. The low density gas responsible for the Ly α forest absorption will expand on a sound-crossing time, which can be considerable – comparable to the Hubble time – in low density systems. Indeed, hydrostatic equilibrium is only restored once $t_{\text{sc}} \sim t_{\text{dyn}}$, where t_{sc} and t_{dyn} are the sound crossing and dynamical timescales respectively (Schaye 2001). The dynamical time t_{dyn} in the low density

IGM is around the Hubble time, as is evident from the Friedmann equation, $H^2 = 8\pi G\rho$.

This can be appreciated in the comparison between sight-line data drawn from the S1 (solid curves) and N1 (dotted curves) simulations displayed in Fig. 5. The left panel displays a subsection of a single sight-line through the simulations at $z = 3$, while the right panel shows the same sight-line at $z = 2$. From top to bottom, the normalised gas densities, the H I fractions, the gas temperatures, the peculiar velocities and the resulting Ly α forest spectra, computed in three different ways, are displayed. The spectra in the third row from bottom are computed using Eq. (5) and include the effect of redshift space distortions due to the peculiar velocity field. In the next row down, however, the N1 spectrum has been recomputed by fixing the gas temperature in the N1 model to be equal to the S1 values and then rescaling the N1 H I fraction, such that $n'_{\text{HI}} = n_{\text{HI}}(T_{\text{S1}}/T_{\text{N1}})^{-0.7}$. This removes most of differences in the spectra attributable to Doppler broadening and the H I fraction, although note the $T^{-0.7}$ scaling for the recombination coefficient is not exact and small differences will remain.⁴ Finally, in the bottom row the S1 and temperature adjusted N1 spectra both exclude the effect of the peculiar velocity field.

It is clear that differences in the Ly α spectra due to the gas density alone are very small soon after reheating. The density distribution at $z = 3$ has only been slightly smoothed around the peak of the overdensity located at $6.9h^{-1}$ Mpc. In contrast, pressure smoothing of the gas distribution is clearly apparent by $z = 2$, as there has been sufficient time for the gas distribution to be altered by the increased pressure in the high density regions. It is precisely because of the finite timescale required to achieve hydrostatic equilibrium that changes in the gas density distribution due to Jeans smoothing will have a negligible impact

⁴ In the left hand panel of Fig. 5 prior to this rescaling, the larger difference between the S1 and N1 line profiles at $5.8h^{-1}$ Mpc on comparison to the broader lines at $6.7h^{-1}$ Mpc is because these absorption features lie on the linear and logarithmic parts of the curve of growth, respectively.

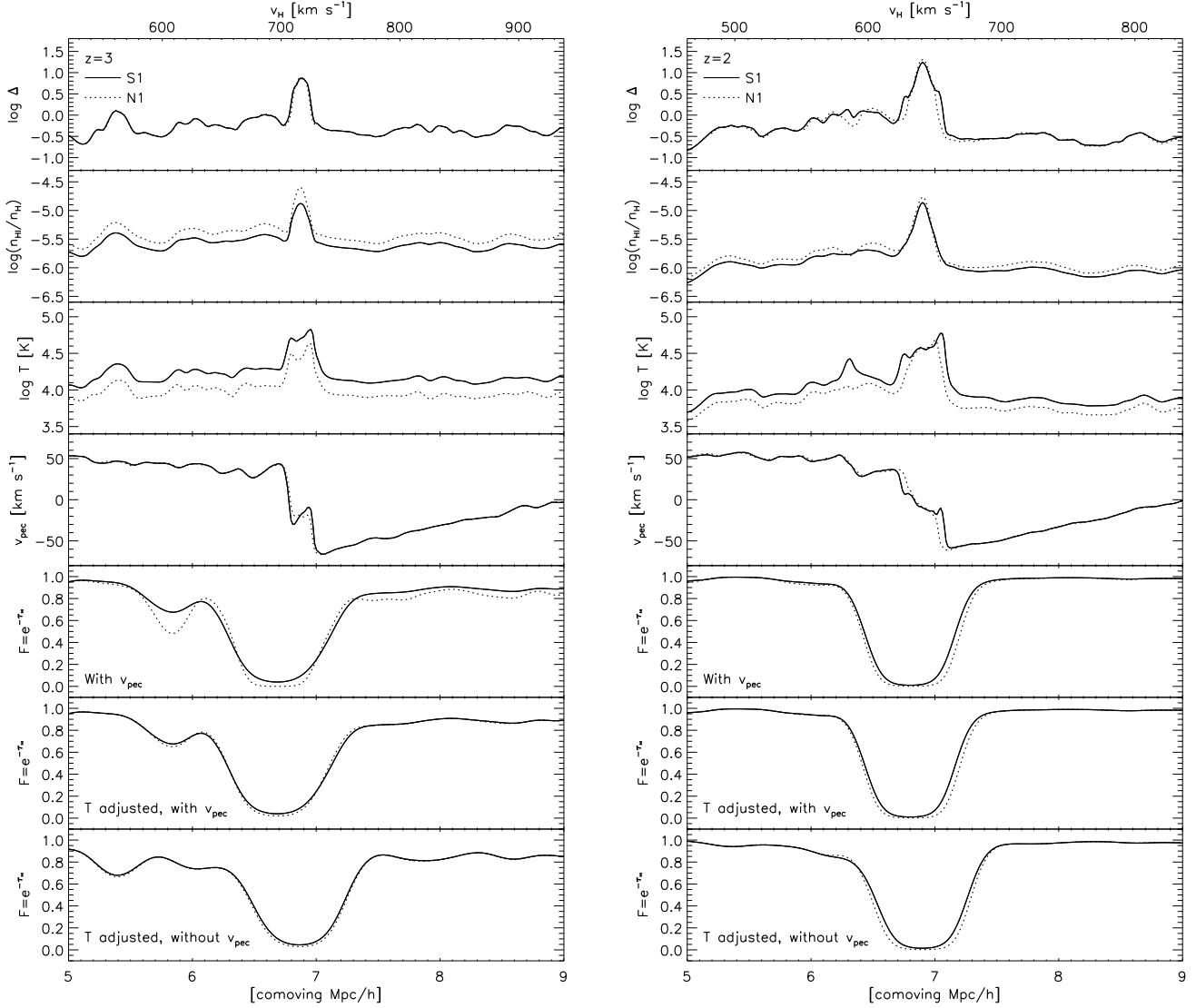


Figure 5. Comparison between various IGM properties along a subsection of a sight-line extracted from identical locations in the S1 (solid curves) and N1 simulations (dotted curves) at $z = 3$ (left panel) and $z = 2$ (right panel). From top to bottom, the normalised gas densities, the H I fractions, the gas temperatures, the peculiar velocities and the resulting Ly α forest spectra, computed in three different ways, are displayed. The spectra in the third row from bottom are computed using Eq. (5) and include the effect of the peculiar velocity field. In the next row, however, the N1 spectrum has been recomputed by fixing the gas temperature in the N1 model to be equal to the S1 values and then rescaling the N1 H I fraction, such that $n'_{\text{HI}} = n_{\text{HI}}(T_{\text{S1}}/T_{\text{N1}})^{-0.7}$. This removes most of differences in the spectra attributable to Doppler broadening and the H I fraction. The S1 and temperature adjusted N1 spectra both exclude the effect of the peculiar velocity field in the bottom row.

on the τ_{eff} evolution by $z = 3$ if the IGM is suddenly heated at $z = 3.4$.

6.2 Peculiar velocity gradients

However, the finite time required to change the gas distribution does not exclude rapid changes in the peculiar velocity *gradients* responsible for the Jeans smoothing. Redshift space distortions associated with these gradients may then induce changes in the Ly α line profiles (Bryan et al. 1999; Theuns et al. 2000). T02 appealed to these peculiar velocity gradients to explain the recovery in the τ_{eff} feature observed by Bernardi et al. (2003). On examination of the peculiar velocity field in the left panel of Fig. 5, the S1 model does

indeed exhibit a positive peculiar velocity gradient in the centre of the overdensity at $6.9h^{-1}$ Mpc, as it must if the gas is to expand. The increased gas pressure has reversed the collapse of the overdensity which is occurring in the colder N1 model.

This effect of gas temperature on the peculiar velocity gradients is displayed in more detail in the left hand panel Fig. 6, where the probability distribution of the peculiar velocity derivative with respect to the Hubble velocity, $\frac{dv_{\text{pec}}}{dv_{\text{H}}}$, in all 1024 synthetic sight-lines is shown for the S1 (solid curve), E1 (dotted curve) and N1 (dashed curve) simulations. The distribution is shown for overdense gas with $\Delta > 3$ only (*cf.* fig. 3 in Theuns et al. 2000). The shift in the probability distribution towards more positive

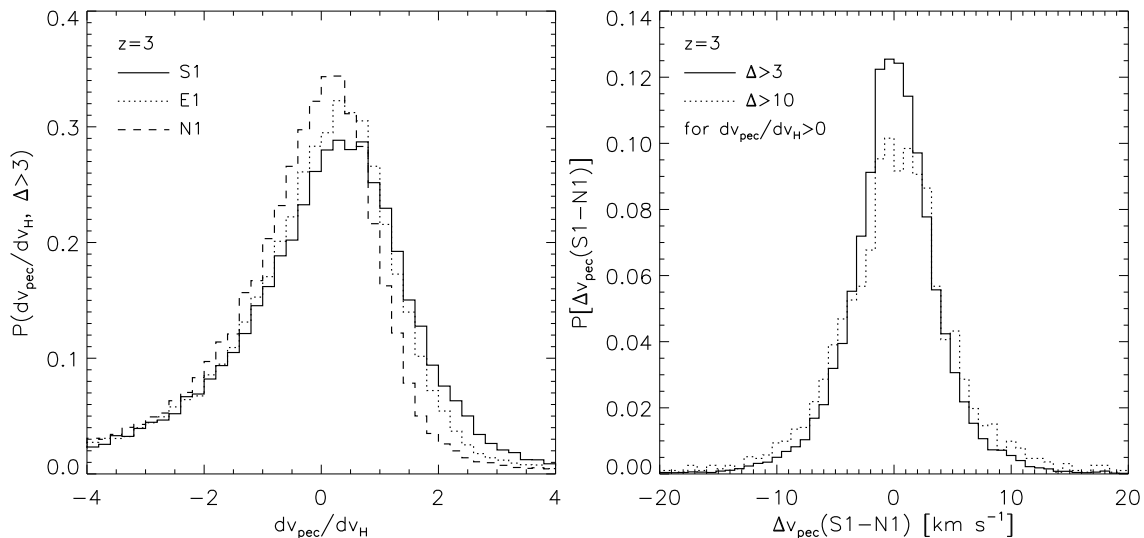


Figure 6. *Left:* Probability distribution of the peculiar velocity derivative with respect to the Hubble velocity, v_H , in the S1 (solid curve), E1 (dotted curve) and N1 (dashed curve) simulations at $z = 3$. The probability distribution is for gas with $\Delta > 3$ only. The shift in the probability distribution towards more positive gradients for progressively hotter models is associated with expansion in the higher density regions in the simulations. *Right:* The probability distribution function of the *difference* between the peculiar velocity fields in the S1 and N1 models in all regions where there is a positive velocity gradient in the S1 model (*i.e.* regions which are expanding). The solid and dotted curves display the distributions for all gas with $\Delta > 3$ and $\Delta > 10$, respectively. The largest values of Δv_{pec} correspond to the densest regions in the simulations.

gradients for progressively hotter models is due to expansion in the hotter, high density regions which have become overpressurised with respect to their cooler surroundings (Bryan et al. 1999; Theuns et al. 2000). The fact that the probability distributions look qualitatively similar to the data presented in fig. 3 of Theuns et al. (2000) is encouraging, and indicates the higher temperatures in our GADGET-2 simulations have a similar impact on the peculiar velocity field. Theuns et al. (2000) used simulations performed with Hydra in a $2.5h^{-1}$ comoving Mpc box with a gas particle masses of $1.14 \times 10^6 M_\odot$ (similar to our mass resolution). Note the larger velocity gradients present in our simulations are most likely due to the additional large scale power present in our significantly bigger simulation boxes ($15h^{-1}$ comoving Mpc).

In the example sight-line in Fig. 5, the maximum difference between the peculiar velocities in the S1 and N1 models is $\sim 10 \text{ km s}^{-1}$ in the centre of the overdensity. Although this difference indeed has some impact on the line profile, the change in the peculiar velocity field is small on comparison to the line width, and it is not enough to significantly alter the broad Ly α line shown in the lower panels. Once differences attributable to the different temperatures of the simulations have been scaled out of the spectra, the S1 and N1 models produce very similar absorption line profiles.

This anecdotal evidence is displayed more quantitatively in the right hand panel of Fig. 6, where we plot the probability distribution of the *difference* between the peculiar velocities in the S1 and N1 simulations for all pixels with positive peculiar velocity gradients and $\Delta > 3$ (solid curve) or $\Delta > 10$ (dotted curve) in the S1 data. In the majority of these regions there is in fact only a small change in the peculiar velocity field, with the largest differences

($\sim 10 \text{ km s}^{-1}$) associated with the rarer, high density peaks which produce broad lines like the one shown in Fig. 5. Since an absolute change in the peculiar velocity of a few km s^{-1} is small in comparison to typical line widths of 20 km s^{-1} (*e.g.* Kim et al. 2002), this explains why the impact on τ_{eff} is correspondingly small. Instead, it is the *instantaneous* temperature of the IGM which primarily influences the τ_{eff} evolution in our simulations.

These findings differ from those of T02, who found changes in the peculiar velocity gradients were partially responsible for inducing the recovery in τ_{eff} observed by Bernardi et al. (2003). The discrepancy between these results may be due to differences between the numerical methods used. We have tested the impact of timestepping on simulations, and we find our results are robust in this respect. However, we cannot be absolutely certain that other numerical effects do not play a role, and ultimately we can only speculate on the origin of this difference. Ideally, an independent numerical study is required to verify or refute our claims. In agreement with T02, however, we do indeed find that line blending due to the peculiar velocity field lowers τ_{eff} , and that higher temperatures steepen the peculiar velocity gradients in overdense regions, giving us confidence that our simulations are at least broadly consistent with T02. We find these effects are nevertheless insufficient to reproduce the narrow feature observed in the τ_{eff} evolution, suggesting that a sudden increase in the IGM temperature at $z = 3.4$ following He II reionisation is unable to adequately explain the τ_{eff} data.

7 ON THE POSSIBLE ORIGIN OF THE OBSERVED τ_{eff} FEATURE

Thus far we have established that even in the presence of a sudden increase in the IGM temperature at $z = 3.4$ (but see Bolton et al. 2009; McQuinn et al. 2009), the narrow dip observed in the Ly α effective optical depth by FG08b cannot be reproduced in our hydrodynamical simulations of the Ly α forest. However, before proceeding it is worth briefly emphasising that the purported “narrowness” of the τ_{eff} feature is a somewhat model-dependent statement. For instance, if we consider the $\tau_{\text{eff}} - z$ plane, FG08b characterised the width of the feature by fitting a power-law and a Gaussian “bump”. This significance of the bump then depends on how good an approximation a power law is to the underlying evolution of the IGM opacity excluding the effects associated with He II reionisation. In our simulations without sudden heating, τ_{eff} does indeed evolve smoothly with redshift, indicating this should be a reasonable approximation if the redshift evolution of the ionising background is gradual. In contrast, the FG08b feature stands out visually, and more objective measures we have applied to the data (*e.g.*, a regularised derivative of the data as used in edge-finding algorithms) tend to confirm that impression. The fact that it is seen in at least three independent data sets further hints that it is not a data artifact.

Additionally, a decomposition into the required evolution in the ionising background (as in the next section) will also depend on the assumed IGM thermal evolution. The Ly α opacity scales as $\tau \propto T^{-0.7}/\Gamma_{\text{HI}}$; the effect of gas temperature and the photoionisation rate on τ_{eff} are impossible to disentangle without independent estimates for one or the other. In this work we have demonstrated that sudden heating on its own cannot reproduce the τ_{eff} feature. However, a sudden global heating event followed by a gradual downturn in the photo-ionisation rate at $z < 3.2$ could still reproduce the data. On the other hand, there is good reason to suppose that such a sudden, global boost to the IGM temperature is unlikely in the first place. Our previous work has demonstrated that He II reionisation likely produces too little heating to produce a substantial opacity decrease over the short timescale required (Bolton et al. 2009). In the next section we therefore assume the temperature boost during He II reionisation occurs over an extended period of time (our E1 model). This model also resembles the results from recent radiative transfer simulations performed by McQuinn et al. (2009). We therefore dispense with rapid changes in the temperature entirely, and now turn to discuss the remaining possibility for the origin of the τ_{eff} dip: a narrow peak in the metagalactic hydrogen photo-ionisation rate.

7.1 The hydrogen photo-ionisation rate required by the FG08b τ_{eff} evolution

Many studies have used simulations of the Ly α forest, combined with measurements of τ_{eff} , to place constraints on the metagalactic hydrogen photo-ionisation rate, $\Gamma_{-12} = \Gamma_{\text{HI}}/10^{-12} \text{ s}^{-1}$ (*e.g.* Rauch et al. 1997; McDonald & Miralda-Escudé 2001; Cen & McDonald 2002; Schaye et al. 2003; Meiksin & White 2004; Tytler et al. 2004; Bolton et al. 2005; Jena et al. 2005). We use the same

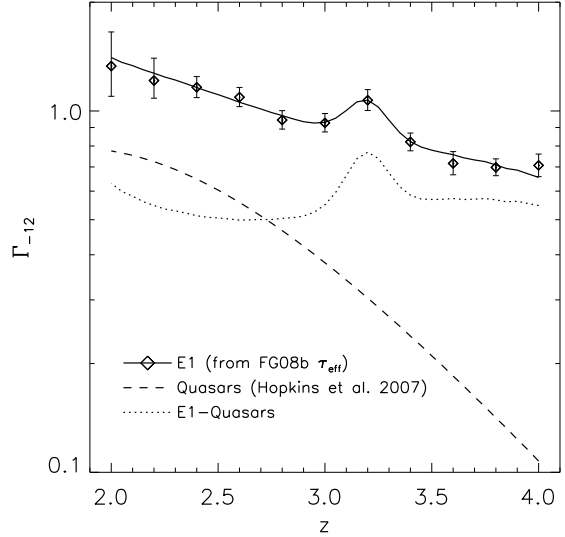


Figure 7. The solid curve corresponds to the hydrogen photo-ionisation rate, $\Gamma_{-12} = \Gamma_{\text{HI}}/10^{-12} \text{ s}^{-1}$, required to reproduce the fit to τ_{eff} evolution presented by FG08b. The narrow peak in Γ_{-12} at $z = 3.2$ corresponds to the dip in the FG08b τ_{eff} evolution. The dashed curve displays the expected contribution of quasars to the total Γ_{-12} , based on the Hopkins et al. (2007) model for the quasar luminosity function, and the dotted line corresponds to the difference between the total Γ_{-12} required by τ_{eff} and the quasar only contribution. To give an indication of the statistical uncertainty in this peak, the open diamonds with error bars are derived from the FG08b τ_{eff} data points including their 1σ statistical errors.

procedure in this work to estimate Γ_{-12} from our E1 simulation, which was constructed to have a similar thermal history at mean density to the recent radiative transfer simulations of McQuinn et al. (2009). We use Eq. (7) and the best fit⁵ to τ_{eff} obtained by FG08b for this purpose. Following from Eq. (2), the photo-ionisation rate which reproduces the τ_{eff} fit is given by $\Gamma_{-12}(z) = \Gamma_{\text{HM}}(z)/A(z)$, where $\Gamma_{\text{HM}}(z)$ is the photo-ionisation rate from the HM01 UVB model divided by 10^{-12} s^{-1} .

The result of this procedure is displayed as the solid curve in Fig. 7. A gradual increase in Γ_{-12} towards lower redshift is required, and the prominent bump centred at $z = 3.2$ is responsible for the narrow τ_{eff} feature. The open diamonds with error bars are derived from the FG08b τ_{eff} data points including their 1σ statistical errors; these give an indication of the statistical uncertainty in this peak.

The dashed curve shows the expected contribution from quasars to Γ_{-12} . We compute this using the recent Madau & Haardt (2009) parameterisation for the Hopkins et al. (2007) comoving quasar emissivity at the H I Lyman limit, ν_{L} , assuming a power law spectrum with $\epsilon_{\nu}^{\text{Q}} = \epsilon_{\text{L}}(\nu/\nu_{\text{L}})^{-1.6}$ (*e.g.* Telfer et al. 2002). The expected contribution to the photo-ionisation rate from quasars is then

⁵ Specifically, we use the best fit FG08b present for their τ_{eff} measurements in redshift bins of width $\Delta z = 0.2$ when using the Schaye et al. (2003) metal correction.

$$\Gamma_{-12}^Q = \frac{(1+z)^3}{10^{-12} \text{ s}^{-1}} \int_{\nu_L}^{\infty} \frac{\epsilon_{\nu}^Q \lambda_{\nu} \sigma_{\nu}}{h_{\text{p}} \nu} d\nu, \quad (9)$$

where the photo-ionisation cross-section $\sigma_{\nu} = 6.3 \times 10^{-18} \text{ cm}^2 (\nu/\nu_L)^{-3}$ and the mean free path $\lambda_{\nu} = \lambda_L (\nu/\nu_L)^{1.5}$; the latter assumes the H I column density distribution $f(N_{\text{HI}}, z) \propto N_{\text{HI}}^{-1.5}$ (e.g. Petitjean et al. 1993; Miralda-Escudé 2003). We take $\lambda_L = 50[(1+z)/4]^{-4}$ proper Mpc, which is towards the lower end of the range recently advocated by Faucher-Giguère et al. (2008a). Note, however, that the expected contribution to the photo-ionisation rate from quasars depends on the uncertain quasar luminosity function (particularly the faint end slope), mean free path and ionising spectrum, as well as radiative transfer effects (for instance, reprocessed radiation can contribute significantly; Haardt & Madau 1996).

The dotted curve displays the difference between the total photo-ionisation rate and the contribution from quasars alone, $\Gamma'_{-12} = \Gamma_{-12} - \Gamma_{-12}^Q$. This curve remains almost flat for $2 \leq z \leq 4$, with $\Gamma'_{-12} \simeq 0.55$, aside from the narrow bump ($\Delta z \simeq 0.4$) which peaks with $\Gamma'_{-12} \simeq 0.75$ at $z = 3.2$. Thus, in order to reproduce the observed τ_{eff} , we require a boost of around 35–40 per cent in Γ'_{-12} over a short redshift interval.

Finally, note that although the FG08b statistical errors are relatively small, the normalisation of Γ_{-12} is still somewhat uncertain. In particular, although the Γ_{-12} values we derive from our E1 simulation are consistent with other estimates from hydrodynamical simulations (Tytler et al. 2004; Bolton et al. 2005; Jena et al. 2005), they are systematically higher by up to a factor of two in comparison to the recent estimates presented by Faucher-Giguère et al. (2008a) using the same τ_{eff} data. This is partially because Faucher-Giguère et al. (2008a) do not model the Ly α forest in detail, instead obtaining analytical estimates for Γ_{-12} using the FGPA combined with the IGM density distribution derived by Miralda-Escudé et al. (2000) (see Eqs. 2 and 3). Systematic uncertainties on Γ_{-12} due to other parameters, such as the IGM temperature which is still poorly constrained, are also large (Bolton et al. 2005). Thus, although the overall normalisation remains somewhat uncertain, the *shape* of these curves should be fairly robust, and the requirement for a peak in Γ_{-12} remains so long as the volume weighted IGM temperature evolves slowly (McQuinn et al. 2009).

7.2 Possible causes of a peak in the hydrogen photo-ionisation rate

We have found that, even if the IGM temperature increases suddenly, we cannot reproduce a narrow feature in the optical depth.⁶ Thus a narrow peak at $z = 3.2$ in the otherwise approximately constant value of $\Gamma'_{-12} = \Gamma_{-12} - \Gamma_{-12}^Q$ from $2 \leq z \leq 4$ is instead required to reproduce the τ_{eff} evolution measured by FG08b. This could be associated with a sharp modulation in the ionising emissivity or the mean free path

⁶ The exception to this is if the photo-ionisation rate subsequently drops at $z < 3.2$ following a sudden temperature increase. However, as discussed earlier, recent studies indicate that a rapid global temperature boost is unlikely in the first place (Bolton et al. 2009; McQuinn et al. 2009).

of ionising photons, or indeed a combination of both effects (see eq. 9). However, while appealing to the stellar contribution to the photo-ionisation rate for an increase in the emissivity is in principle acceptable given the uncertainties involved in deriving the stellar ionising rate from observations at $z > 3$ (e.g. Madau et al. 1999; Miralda-Escudé 2003; Bolton & Haehnelt 2007b), there is no apparent reason why such a narrow peak should occur at $z = 3.2$.

Instead, we suggest it is more likely that any modulation of the hydrogen photo-ionisation rate is associated with dense He II and H I Lyman limit systems. Unlike the low density IGM, such systems have dynamical, cooling and recombination times which are comparable or shorter than the $\sim 10^8$ years associated with the width of the τ_{eff} feature. They can modulate the hydrogen ionising background, either by changing the hydrogen ionising emissivity (by the reprocessing of He II ionising photons into He II Lyman series, Balmer, or two-photon emission, all of which can ionise hydrogen), or the mean free path of ionising photons, since their opacity will be altered by changes in their size or temperature⁷. Furthermore, a sharp change in the emissivity or opacity will then be imprinted on the ionising background on a timescale comparable to the mean free time of an ionising photon, or $\Delta z \sim (dN/dz)^{-1} [\Gamma(\beta - 1)]^{-1} \sim 0.24$ where $dN/dz = 3.3[(1+z)/5]^{1.5}$ is the abundance of H I Lyman limit systems (Storrie-Lombardi et al. 1994), and we assume $dN/dN_{\text{HI}} \propto N_{\text{HI}}^{-\beta}$ with $\beta = 1.5$. This is certainly well within the range required to explain the feature. We now proceed to give some brief examples of these effects.

• Helium recombination radiation

A large fraction of the hydrogen ionising background may be radiation from higher frequencies that is reprocessed by dense systems. For instance, Haardt & Madau (1996) find that ~ 40 per cent of Γ_{HI} comes from reprocessed radiation at $z \sim 3$; similarly (Fardal et al. 1998) find that about ~ 20 per cent of Γ_{HI} comes from reprocessed radiation. This has two components: recombinations of hydrogen to the ground state, and reprocessing of He II ionising photons into He II Lyman series, Balmer, or two-photon emission, all of which can ionise hydrogen. The latter obviously undergoes rapid evolution during the process of He II reionisation, particularly toward the tail end of reionisation when the mean ionising background can rise rapidly (although there could be large fluctuations in the background throughout the reionisation process; Furlanetto 2008). On the other hand, recent calculations indicate the contribution of recombination radiation to the UV background could be smaller than previous estimates (Faucher-Giguère et al. 2009).

Another potentially important effect comes from the reprocessing of He II Lyman series photons in an optically thick IGM (Madau & Haardt 2009). Lyman series photons between Ly β (at $E = 3.56$ Ry) and the Lyman limit (4 Ry) are degraded to He II Ly α photons, Balmer or lower frequency radiation. The magnitude of this effect depends strongly on the opacity of the IGM, which of course evolves rapidly during He II reionisation. While this may not appear to be particularly significant (since such photons can already ionise hydrogen), the degradation to

⁷ The heating and expansion of Lyman limit systems was also briefly discussed by McQuinn et al. (2009).

lower energies implies that the hydrogen photo-ionisation rate can be significantly modulated. In particular, selection rules forbid Ly β photons from being converted into Ly α photons, and the reduction of two-photon emission from Ly β reprocessing at the end of He II reionisation could reduce the hydrogen ionisation rate, helping to produce the downturn in the photo-ionisation rate.

- *The size of Lyman limit systems*

In this work, we have thus far argued that hydrodynamic effects associated with the heating of the IGM during He II reionisation have little effect on τ_{eff} . While this is true of the low-density gas responsible for most of the absorption in the Ly α forest, it is not true of higher density systems which have smaller dynamical timescales, $t_{\text{dyn}} \propto \Delta^{-1/2}$. Since H I Lyman limit systems (LLSs) dominate the opacity of the post-reionisation IGM, a sudden change in the sizes of these systems can then modulate the mean free path of H I ionising photons.

We may consider this as follows. At the end of He II reionisation, although most of the IGM has been reheated, self-shielded He II LLSs, in which the helium is still predominantly in the form of He II, will still remain. Self-shielding to He II ionising photons occurs for He II column densities $N_{\text{HeII}} > 6.7 \times 10^{17} \text{ cm}^{-2}$. Assuming the size of an absorber with overdensity Δ is the local Jeans length (Schaye 2001), this column density corresponds to a characteristic overdensity

$$\Delta \simeq 56 \left(\frac{T}{10^4 \text{ K}} \right)^{2/15} \left(\frac{1+z}{4} \right)^{-3} \Gamma_{-14}^{2/3}, \quad (10)$$

where $\Gamma_{-14} = \Gamma_{\text{HeII}}/10^{-14} \text{ s}^{-1}$ is the He II photo-ionisation rate. The H I column density of an absorber with overdensity Δ is (Schaye 2001)

$$N_{\text{HI}} \simeq 3.8 \times 10^{16} \text{ cm}^{-2} \frac{\Delta^{3/2}}{\Gamma_{-12}} \left(\frac{T}{10^4 \text{ K}} \right)^{-1/5} \left(\frac{1+z}{4} \right)^{9/2}. \quad (11)$$

Self-shielding to He II ionising photons therefore happens at an H I column density of $N_{\text{HI}} \sim 1.6 \times 10^{16} \text{ cm}^{-2} (\Gamma_{-14}/\Gamma_{-12})$, well short of the column density of $N_{\text{HI}} \sim 1.6 \times 10^{17} \text{ cm}^{-2}$ associated with H I LLSs. Thus, at face value it may seem that He II reionisation cannot substantially affect H I LLSs, and hence the mean free path of H I ionising photons. However, He II LLSs are also likely to be associated with the less dense outer regions of H I LLSs. A strong increase in external pressure due to the reheating of the IGM during He II reionisation then means that the system will no longer be in hydrostatic equilibrium. The pressure gradient compresses the gas,⁸ until pressure equilibrium is once again restored (a similar mechanism has been invoked for globular cluster formation, *e.g.* Cen 2001). As a result, H I LLSs could decrease in size, and increase in density. Alternatively, if hard photons can penetrate and heat the gas in the high density regions (*e.g.* Bolton et al. 2009), the H I LLS could instead expand.

For simplicity, let us for the moment assume that any change in density occurs isothermally (see below for more

⁸ Note that in our optically thin hydrodynamical simulations the self-shielding of dense systems is not modelled. This effect is therefore absent in our simulation data.

discussion). If the absorber expands or contracts to a new overdensity $\Delta \rightarrow f\Delta$, then (since $n_{\text{HI}} \propto \Delta^2$ and $R \propto \Delta^{-1/3}$), $N_{\text{HI}} \propto f^{5/3}$, while the cross-section of the absorber $\sigma \propto f^{-2/3}$. Assuming that f is independent of Δ (the opacity is dominated by H I LLSs, so the result will be most affected by f for such systems), the mean free path of ionising photons will be altered by a factor

$$\frac{\lambda}{\lambda_{\text{L}}} = \frac{\int_{0.1}^{\infty} \tau^{-1.5} [1 - \exp(-\tau)] d\tau}{\int_{0.1}^{\infty} f^{-2/3} \tau^{-1.5} [1 - \exp(-f^{5/3}\tau)] d\tau}, \quad (12)$$

where the lower limit $\tau \sim 0.1$ is the H I optical depth at which systems self-shield from He II ionising radiation. This is a non-monotonic function with a minimum at $f \sim 0.3$; however, it does not exceed unity until $f > 1$. Thus, if systems compress, the hydrogen photo-ionisation rate falls; if they expand, it increases. However, more quantitative exploration of this possibility requires careful simulation of the LLS with coupled hydrodynamics and radiative transfer, particularly since the equilibrium temperature also strongly affects opacity. In addition, a caveat to this argument is that although the dynamical timescale of an individual Lyman limit system is comparable to that required for the modulation of Γ_{HI} , a globally averaged change in the mean free path still requires the size of all the Lyman limit systems to change over a short time interval.

- *The temperature of Lyman limit systems*

Yet another possibility is that the temperature of H I LLSs itself changes during He II reionisation, altering their opacity. A rapid increase in the injection of photons at 3 Ry from reprocessed radiation implies that the ionising background hardens significantly, which could result in a significant change in the equilibrium temperature. This could be further modulated by the hydrodynamic effects mentioned above. Note that the equilibrium temperature of LLSs is a non-trivial function of density in photo-ionisation equilibrium, particularly in the presence of metal line cooling (Wiersma et al. 2009). Generally, as a system becomes denser, the equilibrium temperature falls, since the efficiency of cooling increases as the ionisation parameter falls (Efstathiou 1992; Wiersma et al. 2009). This exacerbates the increase in opacity due to the density increase. Non-monotonic evolution of the temperature of the LLSs (which have short cooling times since they cool radiatively, rather than adiabatically) could thus modulate their opacity, as well as the mean free path and the H I photo-ionisation rate.

8 CONCLUSIONS

We have used a semi-analytic model of inhomogeneous He II reionisation and high resolution hydrodynamical simulations of the Ly α forest to investigate the impact of sudden reheating on the evolution of the Ly α forest effective optical depth. Our semi-analytic model indicates that any injection of energy into the IGM during inhomogeneous He II reionisation will produce a well understood and generic evolution in τ_{eff} , where a reduction in the opacity from $z = 4$ to $z = 3$ is followed by a gradual, monotonic recovery driven largely by adiabatic cooling in the low density IGM. This behaviour is inconsistent with the narrow

dip ($\Delta z = 0.4$) of around 10 per cent in τ_{eff} at $z = 3.2$ which has now been detected by three independent observational studies (Bernardi et al. 2003; Dall’Aglio et al. 2008; Faucher-Giguère et al. 2008b).

However, our semi-analytic model does not include a detailed reconstruction of the Ly α forest. We therefore also analyse five high resolution hydrodynamical simulations of the IGM to investigate the effect of various thermal histories on the Ly α effective optical depth. We find that sudden reheating at $z = 3.4$ results in a sharp decrease in τ_{eff} , although we note that such a large, sudden increase in the IGM temperature is nevertheless unlikely to occur over the entire IGM at once (Bolton et al. 2009; McQuinn et al. 2009). This assumption may nevertheless be appropriate in localised regions around quasars with hard spectra, and in small volumes similar to the box size of our hydrodynamical simulations.

Although the assumption of sudden, homogeneous reheating does indeed successfully reproduce the initiation of the τ_{eff} dip observed by FG08b, our simulations are still unable to account for the rapid *recovery* of the narrow dip in the τ_{eff} evolution by $z = 2.9$. The effect of the 8 per cent increase in the free extra fraction following He II reionisation on τ_{eff} is small, and any resulting increase in τ_{eff} is instead countered by the simultaneous flattening of the power-law temperature-density relation during instantaneous, homogeneous He II reionisation at $z = 3.4$. Redshift space distortions in the Ly α forest attributable to the response of the gas to the extra energy injected into the IGM are also unable to account for the observed recovery of τ_{eff} . We find that sudden reheating does indeed produce larger, more positive peculiar velocity gradients corresponding to regions of newly expanding gas, but the absolute changes in the velocity field are generally small in comparison to typical line widths, except in the most overdense regions in the simulation. Such regions are, however, rare and thus contribute little to the average opacity. We have tested the robustness of this result with respect to the time integration scheme employed in GADGET-2, and we find this has little impact on our numerical results. In contrast, if He II reionisation is an extended process, then in agreement with the recent study by McQuinn et al. (2009) we find the Ly α effective optical depth will evolve smoothly with redshift.

As a consequence, we must instead appeal to a narrow peak in the metagalactic hydrogen photo-ionisation rate at $z = 3.2$ to reproduce the τ_{eff} feature in our simulations. This could potentially be modulated by Lyman limit systems, which have recombination, cooling and dynamical times comparable to or less than the timescale associated with the width of the feature. In particular, we suggest that radiative transfer effects from He II reionisation itself could be responsible, either by altering the emissivity of reprocessed He II recombination photons, or by changing the opacity of H I Lyman limit systems and hence the mean free path of ionising photons. However, further detailed investigation is still required to establish the origin of this intriguing feature in the redshift evolution of the Ly α forest opacity.

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