Could the cosmological recombination spectrum help us understand annihilating dark matter?

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ABSTRACT

In this paper, we explore the potential effects of dark matter (DM) annihilations on the cosmological recombination spectrum. With this example, we want to demonstrate that the cosmological recombination spectrum in principle is sensitive to details related to possible extra energy release during recombination. We restrict ourselves to DM models which produce a negligible primordial distortion of the cosmic microwave background (CMB) energy spectrum (usually characterized as $\mu$- and $\gamma$-type distortions). However, since during the epoch of cosmological recombination ($z \sim 1000$) a large fraction of the deposited energy can directly go into ionizations and excitations of neutral atoms, both the cosmological recombination spectrum and ionization history can still be affected significantly. We compute the modifications to the cosmological recombination spectrum using our multilevel H I and He I recombination code, showing that additional photons are created due to uncompensated loops of transitions which are induced by DM annihilations. As we illustrate here, the results depend on the detailed branching of the deposited energy into heating, ionizations and excitations. This dependence in principle should allow us to shed light on the nature of the underlying annihilating DM model (or more generally speaking, the mechanism leading to energy injection) when measuring the cosmological recombination spectrum. However, for current upper limits on the potential DM annihilation rate during recombination the cosmological recombination spectrum is only affected at the level of a few per cent. Nevertheless, we argue here that the cosmological recombination spectrum would provide another independent and very direct way of checking for the presence of sources of extra ionizing or exciting photons at high redshifts. This would open a new window to possible (non-standard) processes occurring before, during and between the three epochs of recombination.

Key words: atomic processes – radiative transfer – cosmic microwave background – dark matter – cosmology: theory.

1 INTRODUCTION

The anomalies in the cosmic-ray spectra of electrons and positrons at 1–1000 GeV seen using PAMELA (Adriani et al. 2009), FERMI (Abdo et al. 2009), HESS (Aharonian et al. 2009) and ATIC (Chang et al. 2008) could be interpreted as signatures from annihilating dark matter (DM; e.g. see Hooper, Taylor & Silk 2004; Cholis et al. 2008; Arkani-Hamed et al. 2009, and references therein). This very intriguing possibility has recently motivated several independent groups to reconsider the effects of such DM models on the ionization history of our Universe and the cosmic microwave background (CMB) temperature and polarization anisotropies (Cirelli, Iocco & Panci 2009; Huetsi, Hektor & Raidal 2009; Kanzaki, Kawasaki & Nakayama 2009; Slatyer, Padmanabhan & Finkbeiner 2009; Galli et al. 2009a) using modified versions of RECFAST (Seager, Sasselov & Scott 1999, 2000; Wong, Moss & Scott 2008).

It is clear that the ionization history is sensitive to the number of extra ionizations and excitations of neutral hydrogen or helium atoms at $z \sim 1000$ (e.g. Peebles, Seager & Hu 2000). These extra ionizations or excitations can be mediated by particles (e.g. electrons or positrons) or photons, which, for example, could both be produced as a consequence of DM annihilations (e.g. Padmanabhan & Finkbeiner 2005) or decaying particles (e.g. Chen & Kamionkowski 2004). Therefore, both DM annihilations and decays of long-lived unstable particles in principle are able to delay recombination, introducing changes to the Thomson visibility function (Sunyaev & Zeldovich 1970a) at $z \sim 1100$, which then affect the CMB temperature and polarization anisotropies, changing the position of the acoustic peaks and their (relative) heights. This allows one to place
interesting constraints on possible DM annihilations during recombination using current and future CMB data (e.g. Zhang et al. 2006; Slatyer et al. 2009; Galli et al. 2009a). Similarly, energy release by decaying particles (Zhang et al. 2007), or more general sources of additional ionizations or excitations of neutral atoms during recombination (Bean, Melchiorri & Silk 2003, 2007; Galli et al. 2008) can be a constraint. Since the presence of any such sources could compromise our ability to measure the spectral index of the primordial power spectrum and its running, it is very important to consider these possibilities carefully, along with other physical corrections to the modelling of cosmological recombination (see Fendt et al. 2009; Sunyaev & Chluba 2009, for detailed overview on recently considered processes).

However, the CMB temperature and polarization anisotropies are not the only (direct) signals that can tell us about (non-standard) processes occurring during cosmological recombination. It is well known that the recombination of hydrogen and helium in the Universe leads to the emission of several photons per baryon, modifying the CMB energy spectrum (Peebles 1968; Zeldovich, Kurt & Syunyaev 1968; Dubrovich 1975; Dubrovich & Stolyarov 1997). Recently, detailed computations of the cosmological recombination spectrum were carried out (e.g. Rubinó-Martín, Chluba & Sunyaev 2006; Chluba & Sunyaev 2006a; Rubinó-Martín, Chluba & Sunyaev 2008), showing that the recombinations of hydrogen and helium lead to relatively narrow spectral features in the CMB energy spectrum. These features were created at redshifts \( z \sim 1300–1400 \), \( z \sim 2100–2400 \) and \( z \sim 6000 \), corresponding to the times of H\(_{\alpha}\), He\(_{\alpha}\) and He\(_{\alpha}\) recombinations, respectively, and, due to redshifting, today they should still be visible at mm, cm and dm wavelengths. Observing these signatures from cosmological recombination may offer an independent way to determine some of the key cosmological parameters, such as the primordial helium abundance, the number density of baryons and the CMB monopole temperature at recombination (e.g. Chluba & Sunyaev 2008a). Furthermore, it will allow us to directly check our understanding of the recombination process and possible non-standard aspects (e.g. see Sunyaev & Chluba 2009, for an overview), for example, in connection with early energy release (Chluba & Sunyaev 2009b).

In this work, we will demonstrate that the cosmological recombination spectrum is also sensitive to the branching of energy released due to DM annihilations into ionizations and excitations. As we show here, it is not only important how much energy is deposited in total, but also when. Depending on the underlying model for the annihilating DM, these efficiencies will differ so that observing the cosmological recombination spectrum may offer another very direct way for constraining such models.

In earlier considerations of the possible effects in connection with energy release during recombination, this branching was either parametrized with single numbers (Peebles et al. 2000; Bean et al. 2003, 2007) or simple approximations were used (Chen & Kamionkowski 2004; Padmanabhan & Finkbeiner 2005; Mapelli, Ferrara & Pierpaoli 2006). A detailed account for all possible aspects of the problem related to a computation of these efficiencies is beyond the scope of this paper. Based on earlier investigations, we will therefore restrict ourselves to some simple examples (see Section 2), which are mainly intended to show the principle dependencies of the recombination spectrum on DM annihilations. However, one can carry out similar computations in connection with decaying particles or other, more speculative sources of extra ionizations (e.g. due to superconducting strings (Ostriker & Thompson 1987) or evaporating primordial black holes (Carr 1976)) during recombination. Since in those cases the cosmological recombination should also exhibit signatures of these non-standard processes while at the same time the CMB anisotropies may be unaffected (see Section 4), we hope that this work will provide further motivation towards refined studies in connection with the cosmological recombination spectrum and potential future experiments measuring it in detail.

But how do additional ionizations or excitations during recombination actually affect the cosmological recombination spectrum? Every extra ionization of hydrogen liberates an electron and proton. At \( z \sim 1000 \), this process is reversed by a recombination of the proton with another free electron after a rather short time. Since the H\(_{\alpha}\) Lyman continuum is completely blocked (Chluba & Sunyaev 2007), the electron is captured into some excited state \( \ell \geq 2 \), emitting at least two photons in the subsequent cascade towards the ground state. This increases the total emission of photons by hydrogen during recombination, since in every additional uncompensated loop of transitions several quanta can be produced. Similarly, extra excitations allow additional electrons to reach high levels, so that in total more recombination photons will be released, by both hydrogen and at \( z \sim 2200 \) also by helium.

The physics of this problem is very similar to the effect of helium photons on hydrogen (Chluba & Sunyaev 2009d), leading to additional feedback-induced emission during recombination. Here the source of extra ionizations and excitations is related to photons emitted by helium at \( z \sim 2200 \) and \( \sim 6000 \) in the normal recombination process. Another example is connected with the changes introduced to the cosmological recombination spectrum as a consequence of early energy release, which produces a primordial y-type distortion (Zeldovich & Sunyaev 1969) of the CMB in the pre-recombinational epoch, leading to uncompensated loops of atomic transitions that attempt to restore the CMB blackbody spectrum (Chluba & Sunyaev 2009b). In this case, the extra ionizations and excitations are caused by the excess of photons in the Wien tail of the distorted CMB.

For the cases considered here the latter process is not important, since the required energy that goes into additional ionizations or excitations during recombination can be tiny in comparison with the energy density of the CMB. In such cases the CMB spectrum is not affected significantly by the additional energy release throughout the entire pre-recombinational epoch (see Appendix A for some estimates), while the dynamics of cosmological recombination can still be strongly modified. This is because the number of hydrogen and helium nuclei is a factor of \( \sim 2 \times 10^{10} \) smaller than the number of CMB photons. This makes it rather easy to perturb the recombination process, while at the same time the CMB energy spectrum itself remains practically unaltered.

2 MODELLING OF THE DIFFERENT PROCESSES

In this section, we give a brief summary of the required equations for our multilevel recombination code to take the effect of DM annihilations into account. We closely follow the approach outlined by Chen & Kamionkowski (2004) in connection with energy injection from decaying particles, including recent modifications and updates related to DM annihilations (Padmanabhan & Finkbeiner 2005; Huetsi et al. 2009; Slatyer et al. 2009; Galli et al. 2009a). However, some of the details here are slightly different, and we also introduce the basis for further improvements of our multilevel recombination code in connection with this problem.
2.1 Overall energy injection rate

Envisioning some self-annihilating DM particle $\chi$ and its antiparticle $\bar{\chi}$, the total rate of energy release per unit volume is given by

$$\frac{dE}{dr} \bigg|_{x^H} = 2 M \chi e^2 (\sigma v) N_x N_{\chi}$$

$$\approx 2.9 \times 10^{-31} \left[ 1 + z \right]^k \text{eV s}^{-1} \text{cm}^{-3}$$

$$\times \left[ \frac{M \chi e^2}{100 \text{GeV}} \right]^{-1} \left[ \frac{\Omega h^2}{0.13} \right] \left[ \frac{\left( \sigma v \right)}{3 \times 10^{-26} \text{cm}^3 \text{s}^{-1}} \right].$$

Here $M \equiv M_{\chi}$ is the mass of the DM particle and its antiparticle, $\langle \sigma v \rangle$ is the thermally averaged product of the cross-section and relative velocity of the annihilating DM particles and $N_x \equiv N_{\chi} = N_{0,\chi} \approx 1.4 \times 10^{-8} \text{cm}^{-3}$.

Depending on the DM model, $\langle \sigma v \rangle$ in general is a function of redshift. In particular, it could also include the effect of Sommerfeld enhancement during the epoch of cosmological recombination (e.g. Slatyer et al. 2009; Galli et al. 2009a), which can be important since at that time the relative velocities of the DM particles become small. One can incorporate these possibilities by replacing $\langle \sigma v \rangle = S(z) \langle \sigma v \rangle_0$, where $\langle \sigma v \rangle_0 = \text{constant}$; however below we will restrict ourselves to cases with $\langle \sigma v \rangle = S(z) \langle \sigma v \rangle_0 = \text{constant}$.

Another aspect of the problem is connected with the clustering of DM (Huetsi et al. 2009). This effect only becomes important at low redshift ($z \lesssim 100$) and leads to an enhancement $\langle N_x N_{\chi} \rangle = B(z) / N_{0,\chi}^2$ of the average squared DM particle number density, with clustering boost factor $B(z) > 1$. This effect can become very pronounced, depending on the assumed halo concentration model and the lower mass cut-off for the halo mass function. However, here we are mainly interested in the CMB spectral distortions generated at $z \gtrsim 200$, where the clustering of matter is negligible. We therefore neglect this aspect of the problem here, using $B(z) = 1$ throughout, so that equations (1) and (2) remain unaltered.

However, it is very important that depending on the involved annihilation channels (e.g. photons, leptons, hadrons, neutrinos) only a fraction, $f_d$, of the released energy will be deposited into the intergalactic medium (IGM), going into heating, and ionizations or excitations of atoms (i.e. hydrogen and helium). For example, energy released in the form of neutrinos (at redshifts of interest to us here) will be carried away so that one can usually expect $f_d < 1$. Furthermore, because the transparency of the Universe to photons and the energy deposition efficiency of different particles (e.g. electrons and positrons) depend on the redshift of injection and the cosmological model (e.g. densities and expansion rate), $f_d$ is a function of time and cosmology.

To include this aspect of the problem into the computations, we therefore write

$$\frac{dE}{dr} \bigg|_{x^H} = f_d(z) \frac{dE}{dr} \bigg|_{x^\chi} = f_d(z) \epsilon_0 N_{\chi} \left[ 1 + z \right]^k \text{eV s}^{-1},$$

with the dimensionless parameter

$$\epsilon_0 = 1.5 \times 10^{-24} \left[ \frac{M \chi e^2}{100 \text{GeV}} \right]^{-1} \left[ \frac{\Omega h^2}{0.13} \right] \left[ \frac{\left( \sigma v \right)}{3 \times 10^{-26} \text{cm}^3 \text{s}^{-1}} \right].$$

Here we also used $N_{\chi} \approx 1.9 \times 10^{-7} \text{cm}^{-3} \left[ 1 + z \right]^3$ for the number density of hydrogen nuclei in the Universe.

A detailed computation for $f_d$ as a function of time and cosmology is beyond the scope of this paper. However, recently Slatyer et al. (2009) computed the function $f_d$ for different models of annihilating DM for the concordance model. Their models give typical values for $f_d$ ranging from $f_d \sim 0.1$ to $f_d \sim 1$, where the largest values are reached for annihilation channels $\chi \bar{\chi} \rightarrow e^+ e^-$ at high ($z \gtrsim 10^3$) redshifts. They also provided some simple fitting formulae for specific DM models, which we will use below. For the purpose of this paper, this should be sufficient.

2.2 Heating of the medium

As mentioned above, only part of the energy that is deposited into the IGM will go into heating of the medium. The rest will lead to ionizations or excitations of hydrogen and helium atoms. If we denote the fraction of the deposited energy that goes into heating by $g_{\text{th}}(z)$, then we can write the additional term in the evolution equation of the temperature of the medium which is related to DM annihilations as

$$\frac{dT_M}{dt} \bigg|_{\text{th}} = \frac{2}{3k} \frac{g_{\text{th}}(z)}{N_{\text{H}} [1 + f_{\text{He}} + X_e]} \frac{dE}{dr} \bigg|_{x^\chi}.$$

Here $N_{\text{H}}$ is the total number of hydrogen nuclei. $f_{\text{He}} \sim 8$ per cent is the number of helium nuclei relative to the number of hydrogen nuclei and $X_e = N_e / N_{\text{H}}$ is the usual free electron fraction.

It is clear that at high redshifts, well before the epoch of recombination, practically all the deposited energy goes into heating of the medium, so that $g_{\text{th}}(z) \sim 1$. Due to energy conservation, one also has $g_{\text{th}}(z) = g_d(z) - g_{\text{ion}}(z) - g_{eX}(z)$, where $g_{\text{ion}}(z)$ is the fraction of the deposited energy that goes into ionizations of atoms and $g_{eX}(z)$ the fraction that goes into excitations. Note that every ionization event also leads to partial heating of the medium, since the liberated electron (and nucleus) will usually have some (large) excess energy which will then be dissipated, e.g. in the form of secondary particles, which can again lead to the heating and ionization of the medium. In detailed computations of the efficiencies $g_{\text{ion}}$ and $g_{eX}$, this has to be accounted for. Below, we will specify the approximations for these functions that we will use in this work.

2.3 Ionizations and excitations of atoms

Knowing that a fraction $g_{\text{ion}}(z)$ of the energy deposited by DM annihilations is going into ionizations of both hydrogen and helium, one can write $g_{\text{ion}}(z) = g_{H}^{\text{ion}}(z) + g_{\text{He}}^{\text{ion}}(z)$, where $g_{H}^{\text{ion}}(z)$ and $g_{\text{He}}^{\text{ion}}(z)$ are the partial contributions of hydrogen and helium, respectively. Similarly, one has $g_{eX}(z) = g_{eX}^{H}(z) + g_{eX}^{\text{He}}(z)$ for the energy that goes into excitations. Since the chemical mixture of the medium can vary, it is useful to introduce the specific contributions $g_{\text{ion}}^{H}(z) = g_{\text{ion}}^{H}(z) N_{\text{H}} / N_{\text{H}}$, with $i = H$ or $i = \text{He}$. Knowing that $a = \text{ion}$ or $i = eX$. For hydrogen one has $N_{\text{H}} / N_{\text{H}} = (1 + f_{\text{He}})$ and for helium $N_{\text{He}} / N_{\text{H}} = (1 + f_{\text{He}}) / f_{\text{He}}$.

For the net ionization rate from the ground states of neutral hydrogen and helium related to DM annihilations, one then finds

$$\frac{dN_{\text{H}}^{\text{ion}}}{dr} \bigg|_{\text{H}} = \frac{1}{1 + f_{\text{He}}} \frac{g_{\text{ion}}^{\text{H}}(z)}{E_{\text{ion}}^{\text{H}}} \frac{dE}{dr} \bigg|_{x^\chi},$$

$$\frac{dN_{\text{He}}^{\text{ion}}}{dr} \bigg|_{\text{He}} = \frac{f_{\text{He}}}{1 + f_{\text{He}}} \frac{g_{\text{ion}}^{\text{He}}(z)}{E_{\text{ion}}^{\text{He}}} \frac{dE}{dr} \bigg|_{x^\chi},$$

where $E_{\text{ion}}^{\text{H}} = 13.6 \text{eV}$ and $E_{\text{ion}}^{\text{He}} = 24.6 \text{eV}$ are the ionization potentials of hydrogen and helium, respectively. For equations (6) and (7), it was assumed that He $\text{He}$ is not important. Furthermore, it is
assumed that the energy which is consumed in each ionization is equal to the ionization energy so that \( \frac{d\tilde{\epsilon}_{\text{ion}}}{d\tau} \) gives the rate of ionization events per unit volume. By definition of \( \tilde{\epsilon}_{\text{ion}} \), this should be possible.

With the same arguments, for excitations of hydrogen and helium from the ground state one can write

\[
\frac{d\tilde{\epsilon}_{\text{ex}}(z)}{dr}_{\text{He}} = - \frac{1}{1 + \tilde{\epsilon}_{\text{He}}} \frac{dE_{\text{ex}}}{dr}_{\text{He}}
\]

(8)

\[
\frac{d\tilde{\epsilon}_{\text{ex}}(z)}{dr}_{\text{H}} = - \frac{\tilde{\epsilon}_{\text{He}}}{1 + \tilde{\epsilon}_{\text{He}}} \frac{dE_{\text{ex}}}{dr}_{\text{He}},
\]

(9)

where now \( E_{\text{He}} = 10.2 \text{ eV} \) and \( E_{\text{He}}^\text{ex} \approx 21.0 \text{ eV} \) are the transition energies to the second shell of hydrogen and helium, respectively. Here, in particular, it is assumed that states with principle quantum numbers \( n > 2 \) are not directly excited. For the purpose of this paper, this approximation will do; however, this may lead to an underestimation of the total number of additional secondary low-frequency photons produced in our computations. This is because excitations to highly excited levels \( (n \gg 2) \) will directly allow some electrons to make transitions among excited states, which leads to emission at low frequencies. On the other hand, for excitations to the second shell the electron will more likely stay within lower levels, and hence produce emission at high frequencies.

2.3.1 Expressions for \( \tilde{\epsilon}_{\text{ion}}(z) \) and \( \tilde{\epsilon}_{\text{ex}}(z) \)

Detailed computations of the specific ionization and excitation fractions \( \tilde{\epsilon}_{\text{ion}}(z) \) and \( \tilde{\epsilon}_{\text{ex}}(z) \) require to follow the evolution of primary and secondary, non-thermal electrons and photons produced by the DM annihilation process (see Kanzaki et al. 2009, for a recent study in this connection), respectively. Such computations have to include several cooling (e.g. Compton and Coulomb cooling for electrons) and particle creation processes (e.g. pair production by photons) and various aspects of radiative transfer (e.g. photon feedback, escape of photons from the main resonances). This is far beyond the scope of this paper; however, based on calulations by Shull & van Steenberg (1985), Chen & Kamionkowski (2004) proposed

\[
\tilde{\epsilon}_{\text{ion}}(z) = \tilde{\epsilon}_{\text{ion}}^\text{H}(z) \approx \frac{1 - X_p}{3},
\]

(10)

where \( X_p = N_p/N_H \) is the free proton fraction. In the context of cosmological recombination with DM annihilations, this approximation has already been used by several authors (Padmanabhan & Finkbeiner 2005; Mapelli et al. 2006; Zhang et al. 2006; Huetsi et al. 2009; Slatyer et al. 2009; Galli et al. 2009a). Furthermore, Padmanabhan & Finkbeiner (2005) also applied a similar expression for the specific ionization and excitation fraction of helium, i.e. \( \tilde{\epsilon}_{\text{ion}}^\text{He}(z) = \tilde{\epsilon}_{\text{ion}}^\text{H}(z) \approx \frac{1 - Z_{\text{He}}}{3}, \) where \( Z_{\text{He}} = N_{\text{He}}/N_H \) is the fraction of singly ionized helium atoms relative to the total number of helium nuclei. In this approximation \( g_n = [1 + 2X_p + f_{\text{He}}(1 + 2Z_{\text{He}})]/3[1 + f_{\text{He}}], \) where both ionizations and excitations were included.

These are very rough approximations, since many details of the computations are not represented or recoverable in this way. For simple order of magnitude computations this approach is certainly acceptable, but for more detailed calculations that aim at including model dependencies, refinements become necessary. For example, without further details it is not easy to say what fraction of excitations is due to photons and what fraction is due to electrons or collisions in more general. In the former case, one should introduce modifications to \( \tilde{\epsilon}_{\text{ex}}(z) \) due to photon escape, which will also strongly depend on the actual density of hydrogen atoms and the expansion rate of the medium, and not only on the ionization fraction. For conditions in our Universe, in that case it will be possible to neglect excitations (e.g. like in Padmanabhan & Finkbeiner 2005), since the escape probabilities, \( P_{\text{esc}} \), in the main resonances of hydrogen and helium are extremely small, so that \( \tilde{\epsilon}_{\text{ex}}^\text{H}(z) \approx P_{\text{esc}} \tilde{\epsilon}_{\text{ex}}^\text{H}(z) \approx 0. \) On the other hand, if excitations mainly occur due to interactions with primary or secondary electrons some such modification is necessary. At low energies \( (E < 100 \text{ eV}) \), this seems to be the case (Shull 1979). As we will see below, the results do depend on this assumption, and a more detailed computation will be necessary.

In addition, Shull & van Steenberg (1985) assumed that the ionization fractions of hydrogen and helium are always equal. This assumption is not valid in the cosmological recombination problem, where helium is completely recombined at \( z \approx 1700 - 1800 \) while hydrogen is still fully ionized. It is not clear that this extreme case can be obtained from their results, since details in the radiative transfer would be very different, likely leading to non-linear scalings. The computations of Shull & van Steenberg (1985) suggest that the scaling of \( \tilde{\epsilon}_{\text{ex}}(z) \) with the ionization fraction from neutral to fully ionized media already appears to be faster than linear. In the case of hydrogen, they provide a fit that is close to

\[
\tilde{\epsilon}_{\text{ion}}^{\text{H, Shull}}(z) \approx \frac{2}{5} \left[ 1 - X_p^{25} \right]^{7/4}.
\]

(11)

Also, their results suggest that for nearly neutral media one has \( \tilde{\epsilon}_{\text{ion}}^\text{H} \approx 1/2, \tilde{\epsilon}_{\text{ex}}^\text{H} \approx 1/2, \tilde{\epsilon}_{\text{He}}^\text{H} \approx 3/5 \) and \( \tilde{\epsilon}_{\text{He}}^\text{He} \approx 3/10 \) (see their table 2 with \( f_{\text{He}} = 10 \text{ per cent} \)) instead of \( \tilde{\epsilon}_{\text{ion}}^\text{He} \approx 1/2, \tilde{\epsilon}_{\text{ex}}^\text{He} \approx 3/5 \). This also means that for more accurate computations within the cosmological context, relatively large differences to the above approximations for \( \tilde{\epsilon}_{\text{ex}}(z) \) can still be expected. A more detailed computation, where we also plan to include doubly ionized helium, excitations of levels with \( n > 2 \) and detailed radiative transfer, will be left for some future work. However, below we will demonstrate that the cosmological recombination spectrum is sensitive to the form of \( \tilde{\epsilon}_{\text{ex}}(z) \), while for cases in connections with DM annihilations the dependence on \( f_d \) is much weaker.

2.4 Modifications to the multilevel recombination code

For the computations presented in this paper, we use our multi-level hydrogen and helium recombination code (Chluba & Sunyaev 2009d). In principle, it allows us to include the effect of feedback of photons and several other recently considered physical processes that modify the recombination history of the Universe at the per cent level (for overview, see Slatyer et al. 2009; Sunyaev & Chluba 2009). However, here we will not include most of these corrections, since the ambiguities introduced due to DM annihilations are much larger in any case. We only want to demonstrate the principal aspects of the problem and show that the cosmological recombination spectrum is sensitive to DM annihilations.

As a first step, one should add the term given by equation (5) to the normal evolution equation for the temperature of the medium. Due to the tight coupling of the photon and electron temperature by Compton scattering, the additional heating will not affect the results of our multilevel recombination code until low redshifts \( (z \lesssim 200) \), where the energy exchange between electrons and photons becomes inefficient. However, the continuous heating of the medium at high redshift in principle will lead to some (small) primordial \( \mu \)- or \( \gamma \)-type spectral distortion of the CMB well before the epoch of recombination (e.g. Zeldovich & Sunyaev 1969; Sunyaev & Zeldovich 1970b; Illarionov & Sunyaev 1975a,b).
We do not take this modification of the background radiation into account, but according to our estimates (see Appendix A), for the DM models under discussion here, their annihilation should never lead to any important primordial CMB distortion (e.g. the $\gamma$ parameter was always smaller than $\sim 10^{-10}$–$10^{-8}$). Therefore, the effects discussed in our previous work on pre-recombinational energy release (Chluba & Sunyaev 2009b) are negligible here. Note that with respect to those earlier computations, the main difference is that we now include direct ionizations and excitations by DM-induced particles, which, as mentioned in Section 1, can be very efficient.

In the multilevel recombination code, one should also add equations (6)–(7) and (8)–(9) to the ground-state rate equations of hydrogen and helium, accordingly. The additional ionizations introduced by DM annihilations will liberate an electron (and nucleus) and, depending on the epoch at which this ionization occurs, the ionization will be directly compensated by the recombination of another electron. In the pre-recombinational epochs one should also allow for direct recombinations to the ground state (Chluba & Sunyaev 2009d), while during the recombination epochs electrons will be captured to excited states (with principle quantum numbers $n > 2$), potentially liberating several (low frequency) photons in the cascade towards the ground state. Also, equation (8) should be subtracted from the rate equation for the H I 2p state in order to ensure conservation of the electron number. Furthermore, we subtract equation (9) from the rate equation of the $2^1P_1$ level, assuming that with $n = 2$ only this level is reached. Since the transition rate to the $2^3P_1$ level is $\sim 10^9$ times smaller, this should be significant. Similarly, we subtract equations (6) and (7) from the rate equation for the free electrons.

3 HOW DOES THE COSMOLOGICAL RECOMBINATION SPECTRUM CHANGE DUE TO DARK MATTER ANNIHILATIONS?

In this section, we illustrate the dependencies of the cosmological recombination spectrum and ionization history on the energy injection by annihilating DM. We start the discussion with cases that assume $f_d(z)\epsilon_\theta = $ constant and then go to more complicated models in the following. We want to demonstrate that the cosmological recombination spectrum is sensitive to differences in the DM annihilation model, in particular related to the branching of the deposited energy into heating, ionizations and excitations.

3.1 Effect of DM annihilation on $N_e$: constant $f_d(z)\epsilon_\theta$

Fig. 1 shows the effect of DM annihilation on the free electron fraction. We included 10 shells for hydrogen and helium into our computations and assumed that all $g_n^f(z)$ are given by equation (10). One can clearly see that DM annihilations have the strongest effect at low redshifts, leading to a delay of recombination (at $z \sim 1000$) and an increase in the residual free electron fraction at very low redshift ($z \sim 200$). The recombination of neutral helium (at $z \sim 2000$) is hardly changing, even in the most extreme cases considered here, implying that the net recombination rate for helium is not affected as strongly by ionizations due to DM annihilation.

One reason for this behaviour is that for a given energy deposition rate, $dE_d/dt$, due to the difference in the ionization potentials, there are approximately two times fewer ionizing photons per helium atom available than for hydrogen (see equations 6 and 7 for confirmation). Similarly, the effective DM-induced excitation rate is approximately two times smaller. Another reason is that any small relative difference $\Delta N_{e}^{He}/N_{e}^{He}$ in the number of free electrons from helium, due to its small abundance ($\sim$ 8 per cent in comparison

to hydrogen), will have an $\sim 13$ times smaller effect on the total ionization history $N_e = N_e^H + N_e^{He}$, which includes the electrons from hydrogen. In the early stages of helium recombination, one therefore expects that the ionization history can only be affected by a comparable amount as during hydrogen recombination when increasing the DM annihilation rate $\sim 20$–$30$ times. Looking at Fig. 1, and comparing the curves for $f_d(z)\epsilon_\theta = 2 \times 10^{-22}$ and $f_d(z)\epsilon_\theta = 5 \times 10^{-22}$ at $z \sim 1100$ and $z \sim 2000$, seems to confirm this statement.

However, towards the end of helium recombination the main reason for the rather small effect of DM annihilations on the free electron fraction is connected with the acceleration of helium recombination caused by the absorption of resonant He I photons in the Lyman continuum of hydrogen (Kholupenko, Ivanchik & Varshalovich 2007; Rubin-Martin, Chluba & Sunyaev 2008; Switzer & Hirata 2008). This shifts the end of helium recombination from $z \sim 1600$ to $z \sim 1750$, because the effective recombination rate of helium is increased many times by this process. It is extremely hard to delay helium recombination with DM annihilations once this process has started working well ($z < 1900$). Here it is also important that the He II ions interact with a bath of free electrons from hydrogen. For each He II ion there are about 13 electrons available for recombinations, while for each proton during hydrogen recombination there is only one. This number of electrons per He II ion remains practically constant until the recombination of hydrogen begins. Therefore, without DM annihilations practically all helium atoms recombine, leaving basically no free He II ions at low redshifts (see Fig. 2).

Only at very late stages, when recombinations of helium are already slow, one again expects some modifications in the helium ionization history. This is because there DM annihilations can (partially) reionize helium atoms, without this process being (significantly) reversed by recombinations. In Fig. 2 one can see that at low redshifts, the number of residual He II ions indeed increases strongly when accounting for DM annihilations. Nevertheless, in all

1 This statement is also true when neglecting the acceleration of He I recombination by the H I continuum absorption.
cases considered no more than $\sim 0.3$ per cent of helium atoms are reionized by DM annihilation at low redshifts, and the total contribution of free electrons from helium to the residual free electron fraction at $z \sim 200$ does not exceed a few per cent.

### 3.2 Effect of DM annihilation on the cosmological recombination spectrum: constant $f_A(z)e_0$

As we have seen in the previous section, the largest modifications in the free electron fraction appear at the end of hydrogen recombination. For the Thomson visibility function and the CMB power spectra the modifications around $z \sim 1100$ are most important, while the huge relative changes in the residual electron fraction at $z \lesssim 500$ actually do not matter that much.

It is known that the recombination lines from hydrogen mainly appear at $z \sim 1300–1400$ (e.g. see Rubiño-Martín et al. 2006; Chluba & Sunyaev 2006b; Chluba, Rubiño-Martín & Sunyaev 2007), where about $\sim 20$ per cent of the hydrogen atoms recombined, while at maximum visibility ($z \sim 1100$) already $\sim 86$ per cent of all H$^+$ were formed. From the differences in the free electron fraction, it is therefore already clear that the H$\alpha$ recombination lines will mainly be modified on the blue sides of the recombination features, with an increase of the emission due to additional ionizations and subsequent recombinations. With increasing DM annihilation efficiency, the changes will become more strong and should eventually also affect the maxima of the recombination features. This will lead to shifts in their positions towards higher frequencies and an increase in the overall amplitude and width of the recombination features.

In Fig. 3, we illustrate this behaviour of the different components in the H$\alpha$ recombination spectrum when taking the effect of DM annihilations into account. We included 10 shells for hydrogen and helium into our computations and assumed that all $g^2(z)$ are given by equation (10). For all components, the DM annihilation increases the overall amplitude of the emission and hence the total number of photons released during recombination. This is simply related to the fact that every extra ionization caused by DM annihilation will liberate an electron which then can recombine to some excited state of hydrogen. From there, it will cascade towards lower levels emitting several photons on its way. These photons are released in addition to those from the normal recombination epoch. Because at $z \lesssim 2000$ the H$\alpha$ Lyman continuum is completely blocked (e.g. Chluba & Sunyaev 2007), every ionization caused by DM annihilations will lead to at least two photons (i.e. Balmer continuum and Lyman $\alpha$) at lower frequency. However, since electrons can also be captured into some highly excited state ($n > 2$), the effective number of emitted photons per ionization can be larger than 2. In Fig. 3 one can see that indeed the emission in transitions among highly excited levels, appearing at low frequencies in the cosmological recombination spectrum, also increases when including the effect of DM annihilations. This shows that a significant number of electrons are captured to states with $n > 2$.

Furthermore, from Fig. 3 it is clear that in particular for the H$\alpha$ bound–bound dipole emission lines the positions and widths of the recombination features are affected by DM annihilations. Both the free–bound and $2s–1s$ two-photon continuum emissions are less sensitive in this respect since they are initially very broad. In this context, especially the spectral features due to Paschen $\alpha$ (visible at $v \sim 120$ GHz) and Balmer $\alpha$ (visible at $v \sim 350$ GHz) are interesting, as they are both very prominent and not overlapping so
much with other spectral features. Nevertheless, the distortions from the recombination epoch are affected in basically all spectral bands. From an observational point of view, it will be important to look at the CMB distortions in many frequency channels and to determine the positions and width of several features simultaneously.

Here we would also like to mention that the very high frequency spectral feature visible at $v \sim 4$ THz is created by the H$_1$ Lyman $\alpha$ resonance at $z \sim 600$. Although from an observational point of view this distortion is not very interesting (the cosmic infrared background is far too strong in this spectral band), for computations of the low redshift chemistry (e.g. see Schleicher et al. 2008, for recent computations), also including the effect of non-equilibrium background radiation (Switzer & Hirata 2005; Vonlanthen et al. 2009), such a feature may be relevant. However, in this case processes directly related to secondary particles from DM annihilation may still be more important (for example, as found in the case of cosmic rays by Jasche, Ciardi & Enßlin 2007).

### 3.2.2 Effect on the helium recombination spectrum

In Fig. 5, we also present the changes in the CMB spectral distortions introduced by He$^i$. Here the effect of DM annihilations is significantly smaller than for hydrogen: for $\Delta f_{\text{dm}} = 5 \times 10^{-22}$ the differences are of the order of $\sim 10$–20 per cent in some bands, while for hydrogen changes of up to $\sim 60$–70 per cent were found (cf. the upper panel in Fig. 3). Still this is much larger than the changes seen in the helium ionization history, where for $\Delta f_{\text{dm}} = 5 \times 10^{-22}$ the corrections were $\Delta N_{\text{he}}^s / N_{\text{he}}^s \sim 0.1 \times 13 \sim 1$ per cent (cf. Fig. 1). This implies that the recombination spectrum is more sensitive to energy deposition during helium recombination than the ionization history itself. Also, it is clear that the small changes in $N_e$ during helium recombination will not propagate very much to the CMB power spectra so that one cannot expect to see any signature of DM annihilation during helium recombination in the CMB. However, directly observing the helium recombination spectrum in principle could shed light on processes occurring during this epoch.

How does this work? In the early stages of helium recombination and in its pre-recombinational epoch ($z \gtrsim 2600$), the number of extra He$^i$ ionizations caused by DM annihilations is very small ($\Delta N_{\text{he}}^s(z) \propto [1 - Z_{\text{He,II}}] \ll 1$). According to our parametrization, most of the deposited energy is going into heating of the medium at that
time (Section 2.3). The same is true for the pre-recombinational epoch of hydrogen, explaining why practically no extra emission is produced, even though DM is continuously annihilating (Section 3.2.1).

Our discussion in Section 3.1 has already shown that the ionization history of helium is not affected as much by DM annihilations; this suggests also that very few extra photons are produced by He I. There, it was most important that the small fraction of neutral hydrogen present at the end of helium recombination leads to a huge increase in the photons’ escape probability so that helium recombination is strongly accelerated. This makes it very hard to change the He I ionization fraction by DM annihilations, since practically every extra ionization is directly reversed by a recombination.

However, during the whole epoch of helium recombination, DM annihilations do lead to some extra ionizations. This drives loops of transitions, which start with the ionization of a neutral helium atom by DM annihilation and end with the release of extra photons at lower frequencies in the cascade of electrons from excited levels towards the ground state. This explains why it is possible to see the effect of extra ionizations in the helium recombination spectrum, while at the same time the effect on the ionization history is much smaller. However, due to the small abundance of helium in our Universe these extra photons are not as important for the total recombination radiation as those from hydrogen. Still for accurate computations of a spectral template one should also take these into account, but in this case also aspects related to details in the Lyman α radiative transfer (Chluba & Sunyaev 2009c,e), feedback (Chluba & Sunyaev 2009d) and electron scattering (Rubino-Martín et al. 2008; Chluba & Sunyaev 2008c) will become important.

In addition, one should mention that, for example, in the case of decaying particles it is possible that during helium recombination much more energy is released than during hydrogen recombination. In this situation the total emission from helium could be increased many times, although the changes to the hydrogen recombination spectrum may be still small. Similarly, the modifications to the ionization history might still be not important for computations of the CMB power spectra, so only the cosmological recombination spectrum will allow us to put constraints on possible extra energy release in this case (see Section 4 for more discussion).

3.3 Dependence of the cosmological recombination spectrum on ionization and excitation efficiencies $\tilde{g}_i$

In the previous sections we have assumed that all the $\tilde{g}_i$ are given by expression (10), which was initially suggested by Chen & Kamionkowski (2004). In this section, we demonstrate that the DM-induced contribution to the cosmological recombination spectrum is sensitive to the branching of the deposited energy into heating, ionizations and excitations. This emphasizes how important it is to refine the modelling of the effect of DM annihilations on the recombination process, especially when aiming at computing detailed templates for the recombination spectrum or accurately accounting for this process in connection with the CMB power spectra.

3.3.1 Changes in the recombination spectrum when excitations by DM annihilations are not included

As a first case, we will assume that excitations caused by DM annihilations are negligible. If additional excitations are mediated by photons, such approximation will be more appropriate. It is clear that this will reduce the amount of extra low frequency emission, since due to the strong coupling of electrons in the 2p state to the 3d and 3s states, and the continuum (Chluba & Sunyaev 2009a,e), every additional excitation to the second shell also leads to some additional ionizations and transitions among highly excited states. In Fig. 6, we show the DM-induced contribution to the total H I cosmological recombination spectrum when neglecting excitations by DM annihilations. At low frequencies the overall amplitude of the distortion is reduced by $\sim 15$–20 per cent, while at high frequencies the distortions are about two times lower than in the case which includes excitations. Also, the changes close to the Balmer and Paschen features reduced by a factor of $\sim 1.5$. For the same $f_\gamma$, one therefore finds a smaller admixture of the DM-induced signal to the normal cosmological recombination spectrum. This implies that the changes in the width and position of the recombination lines will be smaller when excitations are not efficient.

It is clear that part of the difference can be compensated by increasing the effective value $f_\gamma$, but since also the relative amplitudes of the features are affected (e.g. low to high frequency contributions), a differential signal remains. This in principle should allow us to determine how efficient excitations from DM annihilations are. As mentioned in Section 2.3.1, this will depend on the details of the DM model and the annihilation channels that are important.

3.3.2 Direct dependence on the redshift scaling of $\tilde{g}_i$

As pointed out by Chen & Kamionkowski (2004), and as also mentioned in Section 2.3.1, approximation (10) is very crude and does not capture most of the real dependencies of the branching of the deposited energy into heating, excitations and ionizations on the ionization degree, density of the plasma, helium abundance and expansion rate of the Universe. However, the modifications to the recombination spectrum do depend on the scaling of $\tilde{g}_i$ with redshift. As another example, we therefore also compute the changes...
in the recombination spectrum using approximation (10) instead of (11).

In Fig. 7, we present the comparison of these two approximations for the standard ionization history computed with \textsc{recfast} (Seager et al. 1999). As one can see, the ionization efficiency rises much slower when using expression (11), which is more closely based on the work of Shull & van Steenberg (1985), than equation (10). At $z \sim 1100$, i.e. close to the maximum of the Thomson visibility function, $g_{\text{H,Shell}}$ is practically two times smaller than $g_{\text{H,Chen}}$. This implies that ionizations will become important significantly later, implying that the maxima of the DM-induced spectral features should also be shifted towards higher frequencies (see Section 3.2.1 for explanation). Furthermore one expects that the overall amplitude of the additional distortions should be smaller, since the energy deposition rate $dE_d/dt$ decreases with redshift.

In Fig. 6, we present the DM-induced distortion when using $g_{\text{H,Shell}}$ for all $g_i$. Indeed, the maxima of the distortions are shifted by $\Delta \nu/\nu \sim 10$–20 per cent and the overall amplitude reduced by 20–50 per cent at different frequencies. This shows that the cosmological recombination spectrum is rather sensitive to the detailed time dependence of the ionization and excitation efficiencies. Given the large uncertainty in these functions in the current computations, it will be important to refine the modelling of the energy deposition by DM annihilations in this respect. This will also be important in connection with precise computations of the ionization history (see Fig. 8) and CMB power spectra, and the obtained limits on models of annihilating DM using current CMB data (e.g. Huet et al. 2009; Slatyer et al. 2009; Galli et al. 2009a).

### 3.4 Effect of DM annihilations including the time dependence of $f_d$

In the previous sections, we have seen that the shape of the additional spectral distortion did not change very much when increasing the DM annihilation efficiency (see Fig. 4). However, neglecting the effect of excitations or changing the redshift dependence of $g_i^a$ did lead to some notable modifications in the shape of the distortions (cf. Fig. 6). Similarly, one expects that including the additional time dependence in the energy deposition rate will affect the distortions. Here we want to compare the distortions for different models of annihilating DM; however, as we show below the model dependence of $f_d$ introduces only rather small differences.

Inspecting the results of Slatyer et al. (2009) for the energy deposition efficiencies, $f_d(z)$, for different models of annihilating DM, one can see that in most cases only the overall amplitude of $f_d$ is changing, while the shape is very similar, resembling the one for $\chi \chi \rightarrow e^+e^-$ DM annihilation with $M_\chi = 100$ GeV. This leads to a strong degeneracy, since such changes in the overall amplitude can be compensated when allowing for appropriate (constant) boost factors to the annihilation cross-section, e.g. motivated by the effect of Sommerfeld enhancement. Therefore, we expect that in all these cases the shape of the DM-induced spectral distortion will be very similar and that only the amplitude will depend on the specific model via an overall efficiency factor.

The largest differences in the shape of $f_d(z)$ can be found for DM models that annihilate via the channels $\chi \chi \rightarrow e^+e^-$ with $M_\chi = 1$ GeV and $M_\chi = 100$ GeV, and $\chi \chi \rightarrow \mu^+\mu^-$ DM annihilation with $M_\chi = 1$ GeV (see fig. 4 in Slatyer et al. 2009). The functions $f_d$ which we used to represent these cases are shown in Fig. 9. The curves were computed applying the fitting formulae for the different redshift ranges given by Slatyer et al. (2009) and smoothly connecting them at $z \sim 170$ and linearly between $z \sim 1470$ and $z \sim 2500$. One can see that around the time of maximal photon production by helium ($z \sim 2200$), all of these functions are more or less constant at a level of ~1 per cent. Similarly, during the time of photon release by hydrogen ($z \sim 1300$–1400) all the functions $f_d$ are only weakly dependent on time. The largest time dependence is seen for the case $\chi \chi \rightarrow e^+e^-$ with $M_\chi = 100$ GeV. In all the cases shown, the strongest variations appear at much lower redshifts ($z \lesssim 600$) so that from the differences in $f_d$ one does not expect any important changes to the shape of the CMB distortions induced by the different DM annihilation models. The dependence

![Figure 7. Ionization efficiency, $g_i^a$, for the standard ionization history.](image)

![Figure 8. Effect of DM annihilation on the free electron fraction for $f_{\Delta \nu_0} = 2 \times 10^{-22}$ and different choices of $g_i^a$. We included 20 shells for hydrogen and helium in our computations. The curves in the upper group show the relative difference in the free electron fraction in comparison to the reference model (Rubino-Martín et al. 2008). The shaded area indicates the region around the Thomson visibility function, which defines the last scattering surface. For the dashed–dotted line we used $g_i^a$ as given by equation (10), for both hydrogen and helium. In the case ‘no excitations’, we excluded the DM-induced excitations of hydrogen and helium. For the solid curve, we use expression (11) for all $g_i^a$.](image)
J. Chluba

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2009); however, we expect that more detailed computations of the

ionization and excitation efficiencies could reveal more pronounced

dependencies on the specifics of the DM annihilation model.

4 DISCUSSION

In the previous sections, we have focused on illustrating the main
dependencies and effects in connection with the changes that are
introduced by DM annihilations during cosmological recombination.

In this section, we want to go slightly beyond a purely theoretical
study and include current limits on the possible DM annihilation
rate into our considerations. We also wish to extend the discussion
to more general cases of energy release and explain why the cosmo-
logical recombination spectrum could allow us to learn something
about non-standard thermal histories in cases to which the ionization
history and CMB power spectra are not sensitive.

4.1 CMB power spectra versus the cosmological

recombination spectrum

For models of annihilating DM which still seem to be allowed by
current CMB data from Wilkinson Microwave Anisotropy Probe
(WMAP; Komatsu et al. 2009), at 95 per cent confidence level one
has \( f_{\chi \chi} \sigma v \lesssim 3.6 \times 10^{-25} \) cm\(^3\) s\(^{-1}\) (Slatyer et al. 2009; Galli et al. 2009a). In our parametrization, according to equations (3) and
(4), this translates into \( f_{\chi(z)} \sim 2 \times 10^{-22} \) for this value, the
changes in the cosmological recombination spectrum are expected
to be of the order of per cent only (cf. Figs 3 and 4). Unfortunately,
this would make it rather hard to learn something in addition about
DM annihilations from the cosmological recombination spectrum.
So why should one try to measure the cosmological recombination
spectrum when the constraints obtained with the CMB temperature
and polarization power spectra are already so strong?

First, one should emphasize that for the allowed values of \( f_{\chi \chi} \sigma v \) also the changes in the CMB temperature and polarization
power spectra are fairly small: for \( f_{\chi(z)} \sim 2 \times 10^{-22} \) they reach
\( \sim 4-5 \) per cent at high multipolos (see Fig. 11), where the main trend is
completely featureless, and the variable part has an amplitude of
only \( \sim 1 \) per cent. In this context, it is particularly important
that to cosmic variance per cent level correction to the CMB power
spectra only becomes statistically significant at high multipolos.

\(^2\) We arbitrarily chose this frequency.
In cosmological recombination, DM-induced changes to the CMB temperature and polarization power spectra could be significant. The CMB power spectra are strongly degenerate with the values of the scalar spectral index $n_s$, the baryon density $\Omega_b$ and $\sigma_8$ (e.g. see Padmanabhan & Finkbeiner 2005; Slatyer et al. 2009; Galli et al. 2009a). Here also previously neglected physical processes (see Fendt et al. 2009; Sunyaev & Chluba 2009, for overview) are important, since for Planck they can lead to significant biases to the values of $n_s$ and $\Omega_b$ (Rubino-Martín et al. 2009) or otherwise another confusion to the possible signatures of DM annihilations.

Furthermore, at high multipoles the signals related to Sunyaev–Zel’dovich (SZ) clusters (e.g. see Molnar & Birkinshaw 2000; Refregier et al. 2000; Holder & Carlstrom 2001; Komatsu & Seljak 2002) and experimental systematics (e.g. due to calibrations of the beam; Colombo, Pierpaoli & Pritchard 2009) become very important.

Disentangling all the aforementioned components is a major problem; this is where the cosmological recombination spectrum could help in addition: just looking at the amplitude of the DM-induced changes the cosmological recombination spectrum, it is clear that in this respect the recombination spectrum in principle is similarly sensitive as the CMB power spectra (see Fig. 11). Also, the cosmological recombination spectrum does not depend on the value of the scalar spectral index $n_s$, so such confusion is already excluded. In addition, the cosmological recombination spectrum is not limited by cosmic variance: no statistical comparison of the measured energy spectrum with some ensemble of Universes is involved. One would investigate the recombination spectrum for our particular realization of the Universe, only encountering small fluctuations of the cosmological parameters in different directions of the sky. It is even possible to take these corrections into account in the computations of the cosmological recombination spectrum, but for the standard cosmological model the effects are expected to introduce changes that are smaller than $\Delta I_\nu / I_\nu \sim 10^{-4} - 10^{-3}$.

To observe the cosmological recombination spectrum, no absolute measurement is necessary (Chluba & Sunyaev 2008a; Sunyaev & Chluba 2009). Due to its very peculiar frequency dependence, it should therefore in principle be possible to separate the cosmological recombination spectrum from foregrounds and instrumental signals. One can use particularly clean patches on the sky, observing with wide-angle horns. Also the cosmological recombination signal should be polarized at a very small level only, providing another possibility of discriminating it from other signals. By measuring the positions and width of several spectral features in the cosmological recombination spectrum with high precision, one should therefore in principle be able to extract valuable information on possible sources of extra ionizations and excitations during recombination.

4.2 Why the cosmological recombination spectrum could teach us something in addition?

Until now, we have only considered the possibility of continuous, very extended energy release caused by DM annihilations. However, as an example, for energy injection due to long-lived decaying particles additional aspects become important. First, in these cases most of the energy will be released over a characteristic time $\Delta t / t \sim 30 - 40$ per cent. This implies that for decaying particles, the possible changes to the ionization history or cosmological recombination spectrum could be significantly more sharp. Secondly, if the lifetime of the particle is shorter than $\sim 380,000$ yr, then most of the energy will be released before the maximum of the Thomson visibility function. In that case the effect on the CMB temperature and polarization power spectra could be much smaller or even negligible, while the changes in the cosmological recombination spectrum could still be very strong.

For example, if energy were released at $\sim 18,000$ yr after the big bang (corresponding to the time of He III $\rightarrow$ He II recombination) then the CMB power spectra would practically remain unchanged, while at the same time due to extra ionizations and excitations the number of photons produced by He II could increase significantly. Similarly, if the energy is released between the two epochs of helium recombination or during the recombination of neutral helium, the CMB power spectra will mostly be unaltered, while a lot of extra

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Figure 11. DM-induced changes to the CMB temperature and polarization power spectra (upper panel) and the H I cosmological recombination spectrum (lower panel) for $f_{\text{DM}}^0 \sim 2 \times 10^{-23}$. We used CMBeasy to compute the changes in the CMB power spectra (Doran 2005).
emission by He II could be induced. The cosmological recombination spectrum therefore may allow us to check for extra sources of ionizations or excitations at times where the CMB temperature and polarization power spectra are not sensitive.

It is also important to mention that any other modification of the recombination process, e.g. caused by the variation of fundamental constants such as the fine-structure constant $\alpha$ or Newton’s gravitational constant $G$ (see Galli et al. 2009b; Scóccola, Landau & Vucetich 2009, for references), will also lead to changes in the shape and positions of features in the cosmological recombination spectrum. Observing such changes may therefore provide another very clean and direct way to test non-standard physics at early stages of our Universe.

4.3 Towards detailed templates for the cosmological recombination spectrum

There is no principal difficulty in computing the cosmological recombination spectrum with $\sim 0.1$ per cent accuracy, also including the effects of possible energy or particle release in the recombination epoch, even for more general cases. However, in this case, one would have to take previously neglected physical processes (see Fendt et al. 2009; Sunyaev & Chluba 2009, for overview) into account in addition. Here in particular those connected with the Lyman $\alpha$ radiative transfer (Chluba & Sunyaev 2009a,c; Hirata & Forbes 2009) and two-photon transitions (Hirata 2008; Chluba & Sunyaev 2008b; 2009c) will be very important, because they should affect the recombination spectrum at the level of $\sim 10$ per cent (see comments in Chluba & Sunyaev 2009c). Also the effect of electron scattering will be important (Rubino-Martín et al. 2008; Chluba & Sunyaev 2008c), in particular for the contributions from helium.

Furthermore, the computations of the heating, ionization and excitation efficiencies will also have to be refined in order to study detailed model dependencies (for additional comments, see Section 2.3), and as we have seen in Section 3.3, the difference in the recombination spectrum can be large. Also one should include more shells for hydrogen and helium into the computations, since the total amplitude of the additional distortions is still expected to increase, especially at low frequencies.

In addition, the emission due to He II should be taken into account. Given that there is a very extended period between the recombination of He II at $z \sim 6000$ and He I at $z \sim 2200$, the total amount of He II emission induced by DM annihilations is expected to be very important, possibly even exceeding that from neutral helium. Such calculations are beyond the scope of this paper, but we plan to investigate these aspects in more detail in the future.

However, we note that currently it will probably be more important to investigate more general observational prospects in connection with a measurement of the cosmological recombination spectrum, including possible foregrounds and systematics. Also, it will be very important to understand which frequency bands will be most useful and sensitive to changes in the cosmological parameters or energy injection. It is obvious that measuring the cosmological recombination spectrum will be very challenging, but on the other hand there could be a lot to learn from this. In particular, a combination of the CMB temperature and polarization anisotropies with the CMB energy spectrum could open a way to further tighten CMB-based constraints on standard and non-standard aspects of our cosmological model.

5 CONCLUSIONS

Assuming that DM is annihilating throughout the history of our Universe, we have demonstrated that the cosmological recombination spectrum in principle is sensitive to the branching of the deposited energy into heating, ionizations and excitations. If energy only goes into heating of the medium (without leading to some significant primordial $\mu$- or $\gamma$-type CMB distortion), the recombination spectrum is practically not affected, while extra ionizations and excitations lead to modifications in the contributions from both hydrogen (cf. Fig. 3) and helium (cf. Fig. 5).

We have shown that the overall amplitude of the DM-induced spectral distortions depends on the total amount of ionizations and excitations (Fig. 4) at $z \sim 1200–1300$ for hydrogen and $z \sim 2000–2400$ for helium. Furthermore, the relative importance of DM-induced excitations and ionizations, and the time dependence of their efficiencies determine the exact shape and position of the additional CMB spectral distortions from the recombination epoch (see Fig. 6). Since these efficiencies depend on the model considered for the annihilating DM or, more generally, the process that produced the additional ionizations and excitations (e.g. decaying particles), by measuring the cosmological recombination spectrum in several spectral bands in principle one should be able to place additional constraints on possible energy release during cosmological recombination. Given the rather strong dependence of both the changes to the cosmological recombination spectrum and the CMB temperature and polarization power spectra on these efficiencies, for precise predictions in connection with possible extra release of energy during recombination it will be important to refine the computations of these efficiencies in the cosmological context (see Section 2.3 for more comments).

Although for currently allowed values of the effective DM annihilation rate the changes to the cosmological recombination spectrum are of the order of per cent, we have argued that there are several reasons to believe that one could learn something in addition by studying the signals from the recombination era (see Section 4). In particular, the cosmological recombination spectrum is expected to be sensitive to cases of energy or particle release (e.g. before the maximum of the Thomson visibility function), by which the CMB temperature and polarization power spectra are not affected (see Section 4.2). Beyond a check that recombination has occurred as we think it has, observing the cosmological recombination radiation would therefore allow us to directly check our understanding of standard and non-standard physical processes happening at about 260000, 130000 and 18000 yr after the big bang.

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Understanding annihilating dark matter

APPENDIX A: ESTIMATES FOR THE PRIMORDIAL CMB SPECTRAL DISTORTIONS FOR THE CONSIDERED MODELS OF ANNIHILATING DM

A significant fraction of the total energy released due to DM annihilation is going into heating of the medium, in particular at redshifts well before the recombination epoch, where \( f_{\gamma} \approx f_{\gamma} \) and \( f_{\gamma} \approx 0 \). It is clear that this will lead to some primordial distortion of the CMB, since at redshifts \( z \lesssim 2 \times 10^6 \) the thermalization process stops being 100 per cent efficient (e.g. Illarionov & Syunyaev 1975b; Burigana, Danese & de Zotti 1991; Hu & Silk 1993).

One can estimate the expected primordial distortion of the CMB by computing the total change in the energy density of the CMB, \( \Delta \rho_{\gamma}/\rho_{\gamma} = \int \Delta n_{\gamma}/n_{\gamma} \, dt \), that is caused by the deposition of energy in the DM annihilation process. For a \( \mu - \gamma \) type spectral distortion, one then expects \( \mu \approx 1.4 \Delta \rho_{\gamma}/\rho_{\gamma} \) and \( \gamma \approx \frac{1}{2} \Delta \rho_{\gamma}/\rho_{\gamma} \) (e.g. Illarionov & Syunyaev 1975b). For \( s \)-wave annihilation (\( \left< \sigma v \right> = \) constant), one finds (cf. McDonald, Scherrer & Walker 2001)

\[
\mu \approx 3 \times 10^{-10} f_{\gamma, \text{lim}} \left[ \frac{M_{\chi} c^2}{100 \text{GeV}} \right]^{-1} \left[ \frac{\Omega_{\chi h^2}^2}{0.13} \right] \times \frac{\left< \sigma v \right>}{\frac{3 \times 10^{-26}}{\text{cm}^3 \text{s}^{-1}}},
\]

where \( f_{\gamma, \text{lim}} \) is the limiting energy deposition efficiency at large \( z \). This is only a rough estimate, but since the final value for \( \mu \) is so far below the current upper limit \( |\mu| < 9 \times 10^{-8} \) obtained with Cosmic Background Explorer (COBE)/Far Infrared Absolute Spectrophotometer (FIRAS; Mather et al. 1994; Fixsen & Mather 2002), for our purpose it is sufficient. The effective \( \gamma \)-parameter (Zeldovich & Sunyaev 1969) is of the same order.

Although for \( y \approx 10^{-7} \)–10^{-6} a significant number of photons is generated in the pre-recombinational epochs of hydrogen and helium (Chluba & Sunyaev 2009b), inducing interesting narrow spectral features in the cosmological recombination radiation, at the level of \( \mu \) or \( y \approx 10^{-10} \), this process can be completely neglected. In our computations, we can therefore assume that the ambient CMB radiation field is given by a blackbody with present-day temperature \( T_0 = 2.725 \text{ K} \) (Fixsen & Mather 2002).

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