Dark Matter annihilation energy output and its effects on the high-z IGM

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ABSTRACT
We study the case of DM self annihilation, in order to asses its importance as an energy injection mechanism, to the IGM in general, and to the medium within particular DM haloes. We explain the fact that for thermal relic WIMPS, the annihilation cross section is model independent (except for possible low velocity enhancements), and we consider two well motivated WIMP candidates, the SUSY neutralino and the first KK excited state of the $B$ electroweak boson (also called the KK-photon), that have masses in the range favoured by both, indirect searches and accelerator constraints ($m_\chi \gtrsim 100\, \text{GeV}$). We give an overview of the physical scenarios in which these particles arise, mentioning the motivations and general aspects of the underlying physical theories. We then analyse in detail the clustering properties of CDM, on all hierarchy levels, from haloes to subhaloes and DM density profiles within them. In this analysis, we use the Press-Schechter formalism to compute the DM halo mass function, we consider the subhalo mass function calibrated from N-body simulations, and we also consider the process of adiabatical compression by the presence of a SMBH for our computations. We then explicitly compute the energy output (or luminosity) of DM haloes due to annihilations, and compare the obtained luminosities with the standard AGN feedback process, concluding that DM annihilation does not provide the necessary output as to constitute an important feedback mechanism. We then compute the energy injection rate per baryon of annihilations on the IGM, in order to calculate the effects that it has on its temperature and ionization fraction. We find significant deviations in the evolutions of the temperature and ionization fraction of the IGM with respect to the adiabatic cooling case, in scenarios that take into account the clustering of DM at all levels, but mainly at $z \lesssim 20$, where the IGM evolution would be already driven by PopIII stars. We conclude that although the DM annihilation can have a significant impact on the properties of the IGM, it cannot be regarded as an alternative reionization scenario. Regarding the detectability of the WIMP through the modifications to the 21 cm differential brightness temperature signal ($\delta T_b$), we conclude that even considering the adiabatic compression and the substructure, and assuming a maximal absorption fraction of $f_{\text{abs}} = 1$, it is not likely that a $1\, \text{TeV}$ WIMP with thermal relic cross section may be detected. However, a lighter $10\, \text{GeV}$ WIMP could in principle be detected at the $\lesssim 10\, \text{mK}$ level at a redshift of $\sim 45$, corresponding to a frequency of around $30\, \text{MHz}$.

Key words: Cosmology - Dark Matter - Dark Ages - High-z IGM

1 INTRODUCTION

The Dark Matter (DM) is a component of the total density of the universe which can be measured only through gravitational interactions and (likely) weak interactions. It behaves as dust-like matter in the sense that it is collisionless and that its energy density decreases as the third power of the scale factor. There are many compelling evidences of its existence, although all of them are indirect. These are mainly dynamical evidences, which rely on the gravitational influence of DM. Some examples of these evidences are the rotation curves of late type galaxies (Rubin & Ford\textsuperscript{[1970]}), galaxy cluster dynamics (Zwicky\textsuperscript{[1933]}), and strong lensing (Paczynski\textsuperscript{[1987]}), among others. The effects of DM are very important not only for determining the dynamics of galaxies or galaxy clusters, but also for the evolution of the universe.
as a whole, and for the galaxy formation processes and structure formation scenarios.

For example, assuming that the DM consists of particles that were originally in thermal equilibrium with the radiation in the early universe, depending on the mass of the particles, the structure formation scenario will be top down or hierarchical. In the top down structure formation scenario, the first structures to collapse are large (on the mass scale that corresponds to present-day superclusters), because all overdensities at smaller scales are erased by free streaming. This is the case for a hot dark matter (HDM) model, in which the DM particles were relativistic at the time of decoupling. In the hierarchical structure formation scenario, small (and also large) structures start to collapse after decoupling, because DM particles were non-relativistic at the time of decoupling, and so the free streaming is negligible. This is the case for a cold dark matter (CDM) model (Blumenthal 1984). Because the ΛCDM model is currently the most favoured one (see for example the results by Planck (Ade et al. 2013), we study this particular scenario in this work.

The postulated CDM particle has been named a WIMP (or Weakly Interacting Massive Particle), because it weakly couples to normal matter (for introductions to the topic of DM astrophysics, see D'Amico, Kamionkowski & Sigurdson (2009) and Raffelt (2008)). In the standard scenario, this particle is considered to be a thermal relic. This means that its creation and annihilation processes from the early universe radiation field were in thermal equilibrium with each other, and thus, the present day DM number density is the relic abundance after the freeze-out of this processes, with annihilation being favoured after the universe had cooled enough (Kolb & Turner 1990). This fact alone allows for the computation of the interaction cross section for the creation-annihilation process given only the present day abundance of DM (its density parameter Ω_{m,0}).

An important feature of a self-annihilating thermal relic WIMP is that the annihilation products are standard model particles, and thus, they may inject energy to the IGM or to the medium where a particular DM halo is embedded. Globally, the annihilations of WIMPS are unlikely to occur because their abundance has long ago frozen. However, there may be regions where encounters among WIMPS are more probable due to an enhanced density. These regions are precisely the regions where structures form. Now, It is a direct consequence of the ΛCDM model that DM is clumpy, due to the hierarchical growth of structure. Thus, it is precisely in the inner regions of DM haloes and substructures that the annihilations are more probable. The energy output due to annihilations will be proportional to the square of the density, thus, these local density enhancements are of great importance (Natarajan, Croton & Bertone 2008).

The first hierarchy in density enhancements are the DM haloes that do not belong to any gravitationally bound structure at a particular redshift. In N-body DM only simulations, these are called main haloes. An analytic formalism to quantify the number density of these haloes per decade in mass is given by the Press-Schechter formalism (Press & Schechter 1974; Mo & White 1996; Bond et al. 1991; Sheth, Mo & Tormen 2001; Lacey & Cole 1993 among others). This formalism considers that the overdensity profile of the universe is a gaussian field, and thus, the properties of its fourier transform are known. One can define the matter power-spectrum as the square of the complex modulus of the overdensity in k space, and compute the variance of the matter fluctuations of scales of a given mass, as a window-smoothed average of the power spectrum on the corresponding volume. The Press-Schechter formalism then allows to obtain the number density of haloes considering that the gravitational collapse is a self-similar process, and thus, a given halo collapses if its mass scale variance surpasses a certain threshold for collapse at that redshift, and the probability that a halo has the required mass is distributed gaussianly with the variance corresponding to the associated mass.

The second hierarchy in density enhancements are the substructures embedded in a particular halo. As the gravitational collapse is a self-similar process, each halo will be the end node (the trunk) of an inverted halo merger tree at that redshift (Springel et al. 2005). As the DM is non-collisional, it relaxes much slower than normal matter, because it cannot radiate away energy to fall to the bottom of the gravitational potential well. Also, as friction is negligible, the angular momentum barrier will make the DM subhaloes to acquire stable orbits around the center of mass, but these subhaloes will retain their identities. Thus, inside each DM halo, it is expected that the full hierarchy of DM haloes is nested within, until the lowest mass haloes that formed. The number density of these subhaloes can be measured in Milky-Way sized halo simulations, and by approximating their abundance by a power-law behaviour (Via Lactea II simulation (Diemand et al. 2008), Acquarius project simulation (Springel et al. 2008)).

Apart from the number density of haloes and subhaloes, it is important to quantify the form of the density distribution within a particular structure, in order to compute the annihilation power output. (Navarro, Frenk & White 1997), when studying the clustering of DM in N-body simulations, found that haloes from different Cold Dark Matter cosmologies all have the same universal density profile (the NFW profile). This profile is characterized by two parameters, the scale radius and the scale density; and this parameters can be (on average) regarded as a function of the halo mass and the redshift. Also, this DM density profile may be further enhanced by the gravitational effect of baryons on DM through adiabatic compression (see, for example Sellwood & McGaugh (2005); and Natarajan, Croton & Bertone (2008)). Thus, the total resulting DM distribution in the universe can be computed at a particular redshift.

Knowing the global DM distribution, one can compute the resulting energy output per baryon, which can be further compared with what can be obtained assuming a perfectly homogeneous DM distribution (see Mapelli, Ferrara & Pierpaoli 2006; and Cumberbatch, Lattanzi & Silk (2010)). One can also compute the energy generated per unit time by a particular halo of a given mass and redshift, and one can calculate what is the energy received by a particular halo due to the rest of the universe.

This additional energy injection may have interesting effects on the global IGM, as well as on the environment within particular DM haloes. For example, several authors (for example Cumberbatch, Lattanzi & Silk (2010); Mapelli, Ferrara & Pierpaoli (2006); Ripamonti, Mapelli & Ferrara (2007a); Valdés et al. (2007); Pierpaoli (2004); Natarajan...
particle candidates (with masses in the $1\text{MeV} - 100\text{MeV}$ range), possible constraints on the mass and cross section coming from the CMB have been studied (see section 5 of Mapelli, Ferrara & Pierpaoli (2006); and Zhang, Chen, Lei & Si (2006)). However, the effect on the CMB by massive WIMPs, like the ones studied in this work, would be completely negligible. It can be seen that no clear consensus exists neither on the significance of the results as positive indirect detections of DM, nor on the importance of these detections for discriminating among different WIMP scenarios. What can be said nonetheless, is that evidence of possible new exotic physics is accumulating, and that WIMP DM annihilation represents a well motivated option. As for the evidence for Sommerfeld enhancement, it is specially controversial and largely model-dependent. However, the DM annihilation interpretation of the observational evidence supports WIMP masses in the range $\sim 10\text{GeV}$ and $\sim 1\text{TeV}$.

The main goal of this work is to compute this additional energy injection, both locally within haloes and globally, and to assess the importance of this injection on the different processes mentioned above. All the computations will be done considering a cosmology consistent with the Planck 2013 results (Ade et al. 2013). In particular, we will use $h = 0.6711$, $\Omega_{\text{m},0} = 0.3174$, $\Omega_{\Lambda,0} = 0.6825$, $f_{\text{bary}} \equiv \Omega_{\odot} \rho_{\odot}^{\text{bary}} / \rho_{\odot} = 0.17$, $n_s = 0.9624$ and $\sigma_8 = 0.8344$. The paper is organized as follows. In Section 2, the thermal relic WIMP, and the two above mentioned well motivated particular DM models will be explained in greater detail. In Section 3, the statistics of the DM clustering on all hierarchies will be developed and the NFW universal density profile, as well as the adiabatic compression mechanism will be explained. In Section 4, the DM annihilation energy output, both globally (per baryon) and locally within a particular DM halo will be computed. Also, the clumpiness factor, defined as the ratio of output between the smoothed and fully clumped cases, will be calculated. In Section 5, the energy injected on the medium of a particular halo, due to the contribution of the rest of the universe, will be obtained for different cosmological scenarios. In Section 6, the effects of the annihilation energy injection on the global IGM will be followed in detail, computing the IGM temperature, the ionization fraction and the resulting pressure, will be obtained for different cosmological scenarios. The general conclusions will be presented in Section 7.

2 WIMP MODELS AND THE THERMAL RELIC

As was previously mentioned in the introduction, the WIMP is a theorized particle that allows to account for the missing mass in the universe. The evidence of this missing mass is (as previously mentioned) mainly gravitational. This missing mass cannot come however from normal baryonic matter, because current Big-Bang nucleosynthesis constraints, like for example the primordial Helium abundance (Section 10.4 of Ryden (2003), set stringent limits on the normal matter content.

For a certain particle, motivated by some beyond-standard physical scenario, to be a viable WIMP candidate, it must have some basic properties. For example, it has to be charge-neutral and it can interact only weakly with ordinary matter. Also, it has to be collisionless and massive in
order to behave as Cold Dark Matter. A light DM particle will behave as Warm or Hot dark matter depending on if it was relativistic or not at the time of decoupling and also depending on its free streaming length (Angulo & White 2010), and the resulting structure formation scenario or minimum halo mass will be different to the CDM case. Also, a WIMP must be stable, or have a decay time longer than (or at least comparable to) the present Hubble time.

In some well motivated physical scenarios, the WIMP is a thermal relic. This means, as was mentioned before, that its creation and annihilation reaction were in thermal equilibrium with the radiation background in the early universe. Then, as the universe cooled down, the annihilation reaction was increasingly favoured over the creation reaction due to the rest-mass energy difference between both states. Finally, as the expansion rate of the universe became larger than the annihilation reaction rate, the global abundance of DM (its comoving number density) was frozen, and the probability of occurrence of an annihilation event became negligible.

In two well studied scenarios, the WIMP particles are such that they are their own anti-particle. For example in the case of a Kaluza-Klein excited state, arising naturally in compactified extra-dimensional models (like Universal Extra Dimensions or UED, in which all particles are able to propagate in the extra dimension provided they have enough energy), the WIMP is the lightest neutral gauge boson excited state, corresponding to the first excited state of the B boson or $B^{(1)}$, and as the standard model B boson, it is its own antiparticle (the B boson is a superposition of the photon of Quantum Electrodynamics and the Z boson of the weak interaction). Also, in the case of the MSSM (the minimal supersymmetric standard model), considering conserved R-parity (and thus that superpartners are forbidden to decay to standard model particles), the WIMP is the lightest neutralino (i.e., the lightest mass eigenstate of the fermionic superpartners of the gauge bosons and the Higgses), and it is a Majorana fermion, and thus it is its own antiparticle.

In both these scenarios, the annihilation cross section is dependent only on the present day abundance of DM. This fact has been dubbed the WIMP miracle and is explained in what follows.

2.1 The WIMP miracle

The WIMP miracle refers to the fact that, given that the WIMP is a thermal relic, its creation/annihilation cross section can be computed without needing extra model-dependent input (except for the fact that the cross section must not depend on the energy of the WIMP, i.e., the s-wave cross section must be dominant). The cross section then only depends on the present day DM density parameter $\Omega_{m,0}$, and the present day value of the Hubble parameter $h$ (that gives the current expansion rate).

Intuitively, one would think that a higher cross section $\langle \sigma v \rangle$ would lead to less relic abundance of DM, thus resulting in a smaller value of $\Omega_{m,0}$. Both quantities are, indeed, inversely proportional to each other.

Obtaining this cross section is a standard derivation and it is done in, for example, Section 5.2 of Kolb & Turner (1990). For the value of the cross section, Hooper, Finkbeiner & Dobler (2007) use $\langle \sigma v \rangle = 3 \times 10^{-26}$ cm$^3$/s, and in the cosmology used in this work, that gives:

$$\langle \sigma v \rangle = 4 \times 10^{-27} \Omega_{X,0} h^2 \text{cm}^3/\text{s}. \quad (1)$$

In all that follows, we will use that value for the cross section of WIMP annihilation.

Now we will motivate and explain two scenarios that give rise naturally to a thermal relic WIMP with the above-given cross section.

2.2 The KK DM particle

In extra dimensional models (i.e., physical models that are built considering more than three spatial dimensions), the additional spatial degrees of freedom may be accessible to all or some of the particles. Also, these additional dimensions may be compactified differently and on different energy scales. If a standard model particle can propagate in this compactified (and thus finitely extended) dimensions, they will appear more massive in the four dimensional effective theory. In some cases, these excited states of standard model particles may have the desirable features of the WIMP. Here we will briefly discuss the lightest KK particle WIMP candidate and the extra-dimensional scenario in which it arises. We loosely follow the excellent review by Hooper & Profumo (2007).

One extra dimensional scenario that is interesting for being able to provide a suitable DM particle is Universal Extra Dimensions (UED). In this scenario, proposed by Appelquist et al. (2001), all standard model (SM) particles are able to propagate in the extra dimension. Because this dimension is compactified, and thus has a finite extent, the momentum that a particle can acquire in it will be discrete (the idea is similar to that of an electron propagating in a box of given size), thus giving an apparent extra rest-mass in the four dimensional world.

The spectra of different possible momentum states in the extra dimension will give a series of possible masses of a particle, quantized by the number of the excitation state. This masses will depend on the radius of compactification and will be given roughly by:

$$m^2_{X(n)} = \frac{n^2}{R^2} + m^2_{X(0)}, \quad (2)$$

where $m_{X(n)}$ is the mass of the $n^{th}$ excited Kaluza-Klein state of a SM particle $X$, $m_{X(0)}$ is its mass in the SM and $R$ is the compactification radius of the extra dimension, given in natural units (so it has dimension of Energy$^{-1}$). Then, the KK particles are the same SM particles, but considering their excitations in the extra dimension. Interactions between SM particles and KK particles are possible, but are suppressed (thus, these particles are weakly interacting). Moreover, the lightest KK particle is stable because of conservation of momentum in the extra dimension.

As was already mentioned, the lightest KK particle in the UED framework is the $B^{(1)}$, i.e., the first KK excited state of the B gauge boson (also called the KK-photon). This particle is charge neutral, and can annihilate with itself, and thus is a suitable WIMP candidate.
Some of the ways in which KK particles can interact with SM particles are through decays, annihilations and scatterings. For example, a heavier KK particle can decay into a lighter KK particle emitting standard model particles, until it has cascaded into the $B^{(1)}_i$ (also called $\gamma_1$, see Figure 4 of [Hooper & Profumo 2007]).

Also, more interestingly in view of the possible effects of WIMP annihilation on the IGM, two $B^{(1)}_i$s can annihilate to lepton pairs and photons (see Figure 16 of [Hooper & Profumo 2007]).

2.3 The Neutralino

Supersymmetry is a particular symmetry property of the lagrangian of a physical theory, such that said lagrangian remains invariant if a supersymmetric transformation is performed on the fields of the theory. The supersymmetric transformation is such that the change in a fermionic field is a bosonic field, and the change in a bosonic field is a fermionic filed. As in quantum field theory, the fields represent different particle contents (the particles are nothing more than the excited states of the fields), a supersymmetric theory naturally has one boson for each fermion present, and vice-versa.

The particle content of the MSSM (the Minimal Supersymmetric Standard Model) is the same as that of the SM, but for every SM particle, there is an associated supersymmetric partner (also, the Higgs field is different from the one in the SM, being two fields instead of just one). The naming scheme for the superpartners is to maintain the root of the name of the SM particle, but to change the end of the name with the suffix -ino in the case of fermionic superpartners of SM bosons and to add the prefix s- in the case of bosonic (scalar) superpartners of SM fermions. For example, the superpartners of the electron and the muon are the electron and the s-muon, and the superpartners of the photon and the Z boson are the photino and the Zino.

The WIMP candidate of SUSY is the lightest neutralino (unless the gravitino, the superpartner of the graviton in supergravity is lighter than the neutralino, but this scenario will not be considered here). The neutralinos are the mass eigenstates of the superpartners of the neutral gauge bosons and the Higgses, i.e., the Bino ($\tilde{B}$), neutral Wino ($\tilde{W}^0$) and the neutral Higgsinos ($\tilde{H}^0_1$ and $\tilde{H}^0_2$). They are labelled $N_i$, $i = 1, \ldots, 4$ in order of increasing mass, such that the $N_1$ corresponds to the lightest neutralino. The $N_1$ (the WIMP candidate) will be in general a linear combination of the gauginos (Bino and Wino) and the Higgsinos, and it will have a somewhat different phenomenology depending on its specific particle mixture. Although the precise differences between a mainly gaugino, mainly Higgsino or mixed gaugino-Higgsino type neutralino are beyond the scope of this work, it is sufficient to mention that, in general (in order to give the appropriate above-mentioned interaction cross section), a mixed gaugino-Higgsino composition is assumed.

In R-parity conserving SUSY scenarios, no superpartner can decay to SM particles only (SM particles have different R-parity than superpartners), and so the lightest neutralino ($N_1$) will be stable, whereas more massive SUSY particles can decay into the $N_1$ emitting standard model particles (as long as the standard conservation rules are not violated).

Also, the ($\tilde{N}_1$) is a Majorana fermion, and thus it can self-annihilate into SM particles. The SUSY particles can also interact (scatter off) standard model particles, but these scatterings are suppressed, thus the $N_1$ is weakly interacting. The main annihilation processes of neutralino into SM particles is into fermion pairs, although it can also annihilate into gauge bosons (W and Z) at tree-level. At 1-loop level, the neutralino can annihilate directly into gluons or photons, and although this is greatly suppressed compared to tree-level processes, the direct decay into photons is important because it signals a smoking gun of annihilation (the resulting two photons will have an energy very close to the neutralino rest-mass energy, and so it is a clear signal with no contaminants). The Feynman diagrams of the tree-level annihilation channels are shown in Figures 9, 10 and 16 of [Jungman, Kamionkowski & Griest 1996].

Finally, it is important to mention that the phenomenologies of KK particles (as they arise in UED) and of SUSY are very similar, and both of them could be tested, for example in particle colliders, by studying interactions above a certain energy scale. In the case of KK, this energy scale corresponds to the inverse of the radius of compactification of the extra dimension, whereas in SUSY, it corresponds to the mass of the neutralinos and superpartners. Current collider lower limits on the mass scale of this new physics (for example the LHC limits) are on the order of the 100 GeV and increasing. This mass limit will motivate the values that we adopt for the WIMP mass in the rest of this work. We use both a WIMP mass of 10 GeV and of 1 TeV.

3 THE DM CLUSTERING

As it was mentioned, the gravitational collapse process is hierarchical and self-similar, and so different levels of clustering are expected. In the following subsections we explain the prescriptions adopted to account for the DM clustering at all these levels, considering also the distribution of DM within a particular halo or subhalo. In Section 3.1, we follow the analytic approach, using the Press-Schechter formalism ([Press & Schechter 1974], [Mo & White 1996], [Bond et al. 1991]; Sheth, Mo & Tormen 2001; Lacey & Cole 1993 among others) to compute the number density of DM haloes per decade in mass. In Section 3.2, we account for the presence of substructure by considering the substructure mass function given by [Giocoli, Tormen & van den Bosch 2008].

In Section 3.3, we present the universal NFW density profile ([Navarro, Frenk & White 1997], that gives the DM distribution within each gravitationally bound halo, and explain how to calculate its parameters for different halo masses at different redshifts. Finally, in Section 3.4, we explain the effect that the gravitational collapse of the baryonic matter (in particular, the formation of a supermassive black hole (SMBH)) can have on the DM density profile through adiabatic compression.

3.1 The halo mass function

We use this formalism, as presented in [Mo & White 1996], but using the modified Sheth, Mo & Tormen 2001 halo mass function and bias factor, calibrated with the GIF simulations ([Kauffmann et al. 1999]).
We consider the overdensity $\delta(x)$ defined as $\delta(x) = \frac{\rho(x) - \bar{\rho}}{\bar{\rho}}$, and the overdensity threshold for collapse $\delta_c(z)$, defined as the density contrast (in the linear approximation) required for an overdensity in a certain region to already have formed a collapsed object (Navarro, Frenk & White 1997, equation A14). We consider also the mass fluctuation $\sigma(R)$, defined as the standard deviation of the matter overdensity field when smoothed on a scale of size $R$. Finally, the dimensionless mass parameter is defined as $\nu = \frac{\delta_c(z)}{\sigma(R)}$, and it is a measurement of the mass of a collapsed region (DM halo) relative to the mass of the structures that have recently collapsed at redshift $z$.

We will use the S&eth, Mo & Tormen (2001) probability distribution $f(\nu)$. The mass fluctuation $\sigma(R)$ is computed following Lacey & Cole (1993) and Mo & White (1996), starting from the power spectrum of matter overdensities, $P(k)$, at redshift $z$. When calculating the mass function we adopt a primordial power spectrum of the form $P(k) = A_k^{n_s}$, where $A$ is a normalization. We use the Planck 2013 (Ade et al. 2013) scalar index, $n_s = 0.9624$.

The effects on the power spectrum due to the gravitational collapse of the DM on sub-horizon scales after the radiation-matter equality epoch and the Barion Acoustic Oscillations (BAOs) are encoded in the transfer function. We adopt the functional form given by Bardeen et al. (the BBKS transfer function), which ignores the BAOs.

### 3.2 The substructure mass function

The second level of clustering of the DM corresponds to the self-bound haloes that retain their identity within bigger haloes. The presence of these substructures is a natural consequence of the hierarchical structure formation picture, because haloes at all redshifts are formed by the aggregation of smaller parent haloes. As the DM is collisionless, DM haloes are significantly less disrupted than barionic matter during the merging processes, and thus subhaloes may survive up to the smallest scales (due to the negligible, but nonzero, free-streaming of the WIMPS). Many authors (for example Kamionkowski, Kousshiappas & Kuhlen 2010, Giocoli, Pieri & Tormen 2009, Giocoli, Piet & Tormen 2008; Taylor & Babul 2004; Pieri, Bertone & Branchini 2008) have investigated the clustering properties of the subhaloes by means of fitting an analytic form to the subhalo mass function obtained in high-resolution numerical N-body simulations and re-simulations.

We use the substructure mass function proposed by Giocoli, Tormen & van den Bosch 2008, because it also considers the evolution of the substructure in time, due to the combined effects of gravitational heating and tidal stripping in the potential well of the main halo. These effects will tend to erode the subhaloes, which are the remnants of the haloes accreted by the host halo. As was found by these authors, the mass loss (to the smooth component of the main halo) of subhaloes can be approximated by an exponential decay of the subhalo mass on a characteristic timescale that is proportional to the dynamical time of the main halo.

Explicitly, the unevolved subhalo mass function (that considers all subhaloes with the mass they had when they were accreted) is universal, and is given by:

$$\frac{dN}{d\ln(m_v/M_0)} = N_0 x^{-\alpha} e^{-6.283x^3} , \quad x = \frac{m_v}{m_{\text{crit}}},$$

with $\alpha = 0.8$ and $N_0 = 0.21$. These values were calibrated using the GIF Kaufmann et al. (1999) and GIF2 Gao et al. (2004) simulations, as well as resimulations done by Dolag et al. 2005. Here, $m_v$ represents the unevolved subhalo mass (the mass that the halo had when it was accreted) and $M_0$ represents the present-day host mass.

To account for time evolution, the authors find the following relation between the mass of a subhalo at time $t$ (given that it was accreted at time $t_a$), and its unevolved mass:

$$m_{sb}(t) = m_v \exp\left(\frac{t - t_a}{\tau(z)}\right),$$

where $m_{sb}$ is the evolved mass and $\tau(z)$ is the characteristic mass loss time, given by

$$\tau(z) = \tau_0 \left[\frac{\sigma(z, M_0)}{\sigma(z = 0, M_0)}\right]^{-1/2} \left[\frac{H(z)}{H_0}\right]^{-1},$$

and $\tau_0 = 2.0 \text{Gyr}$.

In this work, it was useful to approximate the accretion time of all haloes by the time at which the main halo collapsed, corresponding to a redshift $z_{coll}$, which we take as the redshift at which it acquired half of its mass (see Section 3.3). In terms of the look-back time to redshift $z$ ($LBT(z)$), equation 15 can be approximated by:

$$m_{sb} = m_v \exp\left[\frac{-LBT(z_{coll}) - LBT(z)}{\tau(z)}\right].$$

Now, the required subhalo mass function is the one that includes the mass loss of subhaloes in time, so it should be $\frac{dN}{dm} = \frac{dN}{dm_{\text{ev}}} \frac{dm_{\text{ev}}}{dm}$, using the chain rule. Expressing the subhalo mass function in terms of $m_{sb}$ instead of $m_v$, one obtains:

$$\frac{dN}{dm} = N_0 x^{-\alpha} e^{-6.283x^3} , \quad x = \frac{m K^{-1}(z, M_h)}{m_{\text{crit}}},$$

where $m$ is the present day (at redshift $z$), evolved mass of a subhalo, $M_h$ is the mass of the host halo and $K$ is defined by

$$K(z, M_h) = \exp\left[\frac{-LBT(z_{coll}(M_h)) - LBT(z)}{\tau(z)}\right].$$

It can be seen that the behaviour of the subhalo mass function is that of a power law with an exponential cutoff, and that all mass scales are displaced as a function of time since halo formation. The displacement is such that a present day subhalo of mass $m$, corresponded to a halo of greater mass (by a factor $K$), and so its abundance is decreased because it is the abundance corresponding to the greater mass at the time of halo formation.

These last two equations give the number of subhaloes present in a halo of mass $M$ at redshift $z$, that have a mass

$$N_{\text{ev}}(M, z) = \frac{dN}{dm} \frac{d^3m}{d^3v} \left[\frac{m_{\text{crit}}}{m}\right]^{-\frac{3}{2}}$$

and that all mass scales are displaced as a function of time since halo formation. The displacement is such that a present day subhalo of mass $m$, corresponded to a halo of greater mass (by a factor $K$), and so its abundance is decreased because it is the abundance corresponding to the greater mass at the time of halo formation.
between \( m \) and \( m + dm \), and so correspond to the subhalo mass function that we need. In the next section, we explore the smooth density profile of a particular halo, in order to account for the last level of the DM clustering hierarchy.

### 3.3 The NFW universal density profile

The last hierarchy corresponds to how the DM is smoothly distributed within a particular halo or subhalo. This DM distribution appears to be universal, as found by [Navarro, Frenk & White (1997)] using N-body simulations, and is given by the NFW density profile:

\[
\rho(r) = \frac{\rho_0}{\left(\frac{r}{r_s}\right)(1 + \frac{r}{r_s})^2},
\]

where \( \rho_0 \) is the characteristic density and \( r_s \) is the scale radius. Thus, the NFW profile depends on only two parameters, which on average only depend on the mass of a halo and the redshift at which it is observed. Before obtaining these parameters, it will be necessary to compute the redshift at which a halo with currently (at redshift \( z \)) has a mass \( M \) was formed.

For this quantity [Navarro, Frenk & White (1997)] use the redshift at which a halo which currently (at redshift \( z \)) has a mass \( M \) has just acquired half of its mass. The authors make use of the gaussianity of the overdensity field to compute the redshift of collapse as:

\[
\text{erfc} \left( \frac{\delta(z_{\text{coll}}(M,z)) - \delta_c(z)}{2\sqrt{\sigma^2(0.01M,z) - \sigma^2(M,z)}} \right) = 1/2.
\]

The above equation is an implicit definition of \( z_{\text{coll}}(M,z) \), and the election of the mass of the small substructure (of 0.01 M) is motivated by comparisons to N-body simulations considered by the authors.

With the obtained \( z_{\text{coll}}(M,z) \), we proceed to determine \( \rho_0 \) and \( r_s \) using the following relations:

\[
\rho_0 = \rho_{\text{crit}}(z)\delta_0(M,z),
\]

where \( \rho_{\text{crit}}(z) \) is the critical density of the universe at redshift \( z \), and \( \delta_0(M,z) \) is a characteristic overdensity (or density contrast) that depends on the mass and redshift of the halo whose NFW profile we want to obtain.

Also,

\[
r_s = \frac{r_{\text{vir}}(M,z)}{c(M,z)},
\]

where \( r_{\text{vir}}(M,z) \) is the virial radius of a halo of mass \( M \) at redshift \( z \), corresponding the the radius at which the average density within is equal to the density required for the halo to be a self-gravitating virialized structure (in the case of the Einstein-De Sitter cosmology, this density is 200 times the mean density of the universe). \( c(M,z) \) is the concentration parameter of the halo.

The virial radius is given by [Navarro, Frenk & White (1997)] as:

\[
r_{\text{vir}}(M,z) = 1.63 \times 10^{-2} \left( \frac{M}{h^{-1}M_\odot} \right)^{1/3} \times
\]

\[
\times \left[ \frac{\Omega_{m,0}}{\Omega_m(z)} \right]^{-1/3} (1 + z)^{-1} h^{-1} \text{kpc},
\]

where \( \Omega_m(z) \) is the matter density parameter at redshift \( z \), and the virial radius is given in proper kpc.

The characteristic density is given by:

\[
\delta_0 = 3.41 \times 10^4 \Omega_m(z) \left( \frac{1 + z_{\text{coll}}(M,z)}{1 + z} \right)^3.
\]

Finally, the concentration parameter is implicitly given in terms of the characteristic density contrast by the relation:

\[
\delta_0(M,z) = \frac{200}{3} \frac{c^3(M,z)}{\ln(1 + c(M,z)) - c(M,z)/(1 + c(M,z))}.
\]

Thus, the density profile of a particular halo or subhalo is, on average, completely specified by its mass and redshift.

In the next section, we proceed to study the consequences of the collapse of the baryons on the density profile of the DM haloes, and obtain a modified density profile for the latter considering the growth of a central SMBH, a process that appears to be ubiquitous in galaxy formation.

### 3.4 Adiabatic compression

We consider the modification in the NFW DM density profile resulting from the gravitational collapse of baryons. Baryons behave dynamically differently from DM in their collapse process because they can heat and radiate away their energy, as they are not collisionless and have nonzero internal pressure. Thus, as they lose energy, they fall towards the center of the gravitational potential well (this process is also called baryonic infall or cooling). This is the main mechanism that allows the formation of galaxies within the DM haloes, and it is also responsible for the adiabatic compression of the DM halo density profile. However, baryonic cooling will only take place in DM haloes with masses above a certain critical mass \( M_{\text{crit}} \). A first estimate of this critical mass is the Jeans mass, such that the baryons in haloes with smaller masses are pressure-supported (see [Barkana & Loeb (2001)]). However, the Jeans criteria assumes that the DM perturbation from which the halo is formed is still in the linear regime, which is not the case for already virialized haloes, and therefore, [Tegmark et al. (1996)] consider a slightly different \( M_{\text{crit}} \) criteria that accounts for the details of the collapse process. Finally, one could worry that the presence of DM itself may alter the \( M_{\text{crit}} \), but as shown by [Ripamonti, Mapelli & Ferrara (2007)], although the Jeans mass is altered, the \( M_{\text{crit}} \) as considered by Tegmark changes by an \( O(1) \) factor only (and this is considering LDM, because for the case of the massive WIMPS under study, it would be essentially unchanged). Therefore, in this work, we will consider that only those haloes (main and substructure) that attain a mass greater than \( M_{\text{crit}} \), as given in the Figure 6 of [Tegmark et al. (1996)], are able to undergo adiabatic compression. For haloes of smaller mass, we will simply consider
them to be of pure NFW form. (However, see $\gamma$, were the authors discuss a process that could allow for the formation of Intermediate Mass BHs at $z \sim 30$, which could generate adiabatic compression, but only for very rare haloes).

Different authors have studied the adiabatic compression mechanism, and have given algorithms to compute the compressed DM density profile, given the original, uncompressed profile and the newly grown density profile of the collapsed baryonic component. Two of the most used methods are the Blumenthal algorithm (Blumenthal 1984) and the Young algorithm (Young 1980). The Young algorithm assumes that the potential well has spherical symmetry, and so it corresponds to a central potential problem. The Blumenthal algorithm further assumes that all the orbits of the DM particles are circular. Sellwood & McGaugh (2005) conclude that the Blumenthal algorithm overpredicts the compression in the inner regions.

We use a different algorithm and account only for the adiabatic compression due to the formation of a point mass (i.e., a SMBH) in the center of the potential well. We implement this algorithm as described in Natarajan, Croton & Bertone (2008).

This algorithm assumes that the density profile is of the initial NFW form up to the radius of gravitational influence of the SMBH (defined as the radius up to which it is the main contributor to the contained mass), and that within this radius, the density profile is described by a power-law with a different (steeper) exponent. It also assumes (like the Blumenthal algorithm), that the original orbits of DM particles in the NFW profile were circular.

According to Natarajan, Croton & Bertone (2008), the value of the exponent of the density profile in the inner region is given by

$$\gamma_{\text{spike}} = 2 + \frac{1}{4 - \gamma},$$

where $\gamma$ is the exponent in the inner region of the uncompressed DM density profile ($\gamma = 1$ in the case of the NFW profile).

The derivation of the value of the exponent can be found in Quinlan, Hernquist & Sigurdsson (1995).

Finally, we need the maximum density and the radius of gravitational influence of the BH to have the complete density profile. For the maximum density, we use:

$$\rho_{\text{DM}}^\text{max} = \frac{m_X}{t_{\text{spike}} < \sigma v >},$$

where $m_X$ is the mass of the WIMP particle, and $< \sigma v >$ is the thermal relic annihilation cross-section (see Section 2.1). For $t_{\text{spike}}$, the time elapsed since the spike was formed, we use $t_{\text{spike}} = LBT(z_{\text{coll}}) - LBT(z)$, i.e., the look-back time at the redshift at which the halo collapsed (defined in Section 3.3) as measured from the current redshift.

For the gravitational influence radius of the black hole, we use:

$$r_{\text{BH}} = 0.2 \frac{GM_{\text{BH}}}{\sigma_{\text{sp}}^2},$$

where $M_{\text{BH}}$ is the black hole mass, and $\sigma_{\text{sp}}$ is the velocity dispersion of the spheroid component (the bulge in the case of haloes typical of late type galaxies).

For $\sigma_{\text{sp}}$, we use the observational relation between spheroid velocity dispersion and BH mass, as presented in Murray, Quataert & Thompson (2005):

$$M_{\text{BH}} = 1.5 \times 10^8 \left( \frac{\sigma_{\text{sp}}}{200 \text{[km/s]}} \right)^4 M_\odot. \quad (19)$$

The mass of the black hole depends, on average, only on the mass of the host halo and on the fact that the host is a main halo or a subhalo. The relation between both masses is given by Lagos, Cora & Padilla (2008):

$$\log \left( \frac{M_{\text{BH}}}{M_\odot} \right) = 0.84 \log \left( \frac{M_{\text{halo}}}{M_\odot} \right) - 2.1 \quad (20)$$

for the case of a main halo, and

$$\log \left( \frac{M_{\text{BH}}}{M_\odot} \right) = 0.84 \log \left( \frac{M_{\text{halo}}}{M_\odot} \right) - 2.9 \quad (21)$$

for the case of a subhalo.

Now we have all the ingredients for computing the adiabatically compressed DM density profile for a given halo of mass $M$ at redshift $z$.

Putting all the pieces together, the adiabatically compressed density profile of a halo has four different regimes, and is given by:

$$\rho_{\text{DM}}(r) = \rho_{\text{NFW}}(r), \ r > r_{\text{BH}}, \quad (22)$$

$$\rho_{\text{DM}}(r) = \rho_{\text{NFW}}(r_{\text{BH}}) \left( \frac{r}{r_{\text{BH}}} \right)^{-\gamma_{\text{spike}}}, \ r_{\text{plateau}} < r < r_{\text{BH}}, \quad (23)$$

$$\rho_{\text{DM}}(r) = \frac{m_X}{t_{\text{spike}} < \sigma v >}, \ 4r_s < r < r_{\text{plateau}}, \quad (24)$$

$$\rho_{\text{DM}}(r) = 0, \ r < 4r_s, \quad (25)$$

where $r_{\text{plateau}}$ is the radius at which the density of the spike equals the maximum attainable DM density, and $r_s$ is the Schwarzschild radius of the BH. In Figure 1, we show the adiabatically compressed density profile of a DM halo of mass $5 \times 10^{13} M_\odot$ at redshift $z = 1$ (Typical of QSO systems). The four regimes are clearly distinguishable. The outer slope corresponds to the NFW $\gamma = 1$ and the inner slope corresponds to $\gamma_{\text{spike}} = \frac{7}{4}$.

### 4 DM ANNIHILATION ENERGY OUTPUT

Having considered the clustering properties of the DM at all scales and hierarchy levels, and having explained two particular DM WIMP candidates, motivating their expected masses and annihilation cross-sections, we can proceed to compute the expected energy output of annihilations in local structures (halos) and as an average for the whole IGM.

In this section, we obtain the luminosity due to annihilations of individual DM haloes, considering the clustering analysed
Dark Matter annihilation energy output and its effects on the high-z IGM

As the universe may be infinite in spatial extent (as would be the case for flat, i.e. $\Lambda$CDM and Einstein-De Sitter, and Open cosmologies), the relevant quantity is the energy injected to the IGM per baryon. In terms of the present day WIMP number density, this is calculated as:

$$\dot{E}_\chi^{smooth} = \frac{1}{2} \frac{m_\chi}{n_{b,0}} c^2 n_\chi^2 (1 + z)^3 < \sigma v > f_{abs}(z),$$

where $n_\chi,0$ is the present day DM number density, and is calculated as

$$n_\chi,0 = \Omega_{dm,0} \frac{\rho_{c,0}}{\mu m_H},$$

where $\Omega_{dm,0} = (1 - f_{bary})\Omega_{m,0}$ is the present day DM density parameter, $\rho_{c,0}$ is the present day critical density of the universe and $m_\chi$ is the WIMP mass. Also, $n_{b,0}$ is the present day baryon density, and is calculated as

$$n_{b,0} = \Omega_{b,0} \frac{\rho_{c,0}}{\mu m_H},$$

where $\Omega_{b,0}$ is the present day baryonic density parameter (equal to $f_{bary}\Omega_{m,0}$), $m_H$ is the hydrogen mass and $\mu$ is the mean atomic weight of the baryon content. The mean atomic weight, considering a universe with only hydrogen and helium (a good approximation, particularly before star formation), can be computed as $\mu = f_H + 4 f_{He}$, where $f_H$ is the fraction of hydrogen and $f_{He}$ is the fraction of helium by number. Assuming a value of $X = 0.74$ and $Y = 0.26$ (the hydrogen and helium fractions by mass, respectively), gives $f_H = 0.92$ and $f_{He} = 0.08$. Thus, $\mu = 1.24$ will be used in this work.

Knowing the absorption fraction of the energy $f_{abs}(z)$, we can readily compute the energy injected to the IGM per baryon through DM annihilations. Different authors use various prescriptions for $f_{abs}(z)$. Natarajan & Schwarz (2009) consider that only the photons resulting as secondary annihilation products can inject energy to the IGM, and calculate the photon energy spectrum for different neutralino models explicitly. They consider that as the photons propagate, they lose energy, and compute the probability that they scatter off the IGM atoms. They find that $f_{abs}(z)$ is in the range of 0.1 – 0.2 for $z < 50$ (they do not consider higher redshifts).

Cumberbatch, Lattanzi & Silk (2010), compute $f_{abs}(z)$ for the neutralino (see also Ripamonti, Mapelli & Ferrara [2007a]), calculating first the number of photons and electrons produced per neutralino annihilation event, and they obtain $f_{abs}(z)$ in the range of 0.01 – 0.1, being higher for higher $z$ (and considering redshifts up to z = 1500).

Other authors, like Mapelli, Ferrara & Pierpaoli (2006), simply assume that $f_{abs}(z) = 1$, and compute the maximal effects that their WIMP candidates can have on the IGM. We will consider that $f_{abs}$ is a constant that does not depend on redshift, and we will use the values $f_{abs} = 0.01, 0.1$ and 1, in order to obtain results comparable to the different authors.

In Figure 2, we show the results obtained for the smooth energy injection per baryon, for two different WIMP masses (1TeV and 10GeV), as explained in the caption. We only consider the maximal absorption fraction ($f_{abs} = 1$), be-

---

**Figure 1.** The adiabatically compressed density profile of a halo with mass $5 \times 10^{13} M_\odot$, at redshift $z = 1$.
cause for other values, the resulting curves should simply be rescaled by $f_{\text{abs}}$. It can be seen that the effect of a lower WIMP mass is, as expected, to boost the energy injection rate by a factor $\propto m_\chi^{-1}$.

### 4.2 The DM annihilation luminosity of a halo

Having considered the global case of the energy injection to the IGM, we now consider the energy output per unit time (or luminosity) of a particular halo of mass $M$ at redshift $z$. Following [Natarajan, Croton & Bertone (2008)], from the general formula the energy output integrated on the volume of the halo is given by

$$L_\chi = \frac{<\sigma v > c^2}{2m_\chi} \int_V \rho_{DM}(x,t) d^3x,$$

(32)

where $\rho_{DM}$ is the complete DM density profile of the halo, accounting for the smooth main halo, the substructure and the adiabatic compression in both (see Section 3.2 - Section 3.4). This luminosity can be decomposed into the main halo luminosity and the substructure luminosity such that:

$$L_\chi = L_{\chi,\text{main}} + L_{\chi,\text{subs}},$$

(33)

where

$$L_{\chi,\text{main}} = L_\chi(M_{\text{main}}, z) = \frac{<\sigma v > c^2}{2m_\chi} \int_V \rho_{DM}(r, M_{\text{main}}, z) dV,$$

(34)

and

$$L_{\chi,\text{subs}} = L_{\chi,\text{subs}}(M_{\text{main}}, z) = \int_{m_{\text{free}}}^{M_{\text{main}}} \frac{dN}{dm} L_\chi(m, z) dm.$$  

(35)

Note that in the above equations, $L_\chi(M,z)$ is the DM annihilation luminosity of a smooth, adiabatically compressed halo or subhalo, the only difference among them is that for the halo or the subhalo, the mass of the SMBH depends differently on the mass of the halo. It is important to mention that in order to account for the substructure luminosity, an integration of the substructure mass function multiplied by the luminosity of an individual subhalo in all the substructure mass range must be performed. This integration should be done between $m_{\text{free}}$, the free streaming mass of the WIMP (the minimum mass that a virialized DM halo can have), and $M_{\text{main}}$, the mass of the main halo (no substructure can be as massive as the main halo). The free streaming mass will depend naturally on the WIMP mass, as $m_{\text{free}} \propto m_\chi^{-3}$. Based on the analysis of the effects of the free streaming length on the matter power spectrum by [Angulo & White (2010)], we consider a free streaming mass of $m_\chi = 10^{-3} M_\odot$ for a WIMP mass of $m_\chi = 1\text{TeV}$, and a free streaming mass of $m_\chi = 10^{-4} M_\odot$ for a WIMP mass of $m_\chi = 10\text{GeV}$.

We explicitly compute the luminosity of DM haloes, in the $\Lambda$CDM cosmology, considering two different WIMP masses ($10\text{GeV}$ and $1\text{TeV}$) and the presence or absence of substructure and adiabatic compression. In Figure 3, we show the results in the form of the light to mass ratio ($L_\chi/M$) for the annihilation luminosity, as a function of halo mass.

It can be seen that, for the same halo mass, the luminosity of the halo increases with redshift. This is a result of the structure formation process, because a DM halo that formed earlier will be more concentrated and have a higher characteristic density. Thus, its annihilation rate, and in consequence its luminosity, will be increased. Also, it can be seen for the cases that include adiabatic compression, that for each redshift there is a discontinuity in luminosity at the characteristic density. Thus, its annihilation rate, and in consequence its luminosity, will be increased. Also, it can be seen for the cases that include adiabatic compression, that for each redshift there is a discontinuity in luminosity at the characteristic density. Thus, its annihilation rate, and in consequence its luminosity, will be increased. Also, it can be seen for the cases that include adiabatic compression, that for each redshift there is a discontinuity in luminosity at the characteristic density. Thus, its annihilation rate, and in consequence its luminosity, will be increased.
Figure 3. The luminosity to mass \((L_{\chi}/M)\) ratio of the internal DM annihilation luminosity of individual haloes, as a function of halo mass (in the range from \(10^{-4}\) to \(10^{17} M_\odot\)). The different curves in the four plots correspond to different redshifts; in particular, the dotted curves correspond to \(z = 0\), the short-dashed curves correspond to \(z = 5\), the dot-dashed curves correspond to \(z = 10\), the triple-dot-dashed curves to \(z = 50\) and the long-dashed curves to \(z = 100\). The different plots correspond to different clustering scenarios and WIMP masses; in particular, the top-left plot corresponds to the case with substructure and adiabatic compression, with a WIMP mass of 1TeV; the top-right plot corresponds to the case with substructure and 1TeV WIMP mass, but with no adiabatic compression; the bottom-left plot corresponds to the case with a WIMP mass of 1TeV and adiabatic compression, but no substructure; and the bottom-right plot corresponds to the case with adiabatic compression and substructure, but a WIMP mass of 10GeV. In all the plots we also include the \(\nu = 1\) mass values for the 3 lowest redshifts in the corresponding line-styles.

compare the obtained light to mass ratios for the annihilation luminosity with the typical mass to light ratios in haloes at different mass scales. For haloes with masses in the \(10^3 M_\odot - 10^{17} M_\odot\) range, the logarithm of the mass to light ratio \((\log_{10}(\mathcal{L}))\) at redshift zero (near universe) varies between 1.5 and 3, as shown by Eke et al. (2004) and also by Marinoni & Hudson (2002). Also, \(\mathcal{L}\) has a minimum for halo masses around \(10^{12} M_\odot\), and it increases for lower masses due to SNe feedback and for higher masses due to AGN feedback (see Lagos, Cora & Padilla (2008) and also Mutch, Croton & Poole (2013)). Thus, we note that the corresponding light to mass ratios for the haloes \((\mathcal{L}/M)\) are much higher than the \(L_{\chi}/M\) curves that we obtain for redshift zero. Furthermore, the minimum in \(\mathcal{L}\) occurs for masses where the adiabatic compression is already present, and therefore the \(L_{\chi}/M\) is maximum, and thus both mass to light ratios correlate positively instead of negatively as it would be expected in the case of star-formation-quenching feedback. Therefore, we can conclude that the DM annihilation is irrelevant as a feedback mechanism. (However, see Ascasibar (2006).) It should be remembered that, although we plot halo luminosities up to haloes of mass of \(10^{17}\), the most massive gravitationally bound haloes today (at \(z = 0\)), have masses of the order of \(10^{15} M_\odot\), so more massive haloes are virtually non-existent; and further more, the masses of the haloes that have a particular number density at a certain redshift (given by the characteristic dimensionless mass \((\nu)\)), falls rapidly with redshift. For clarity, we also include in Figure 3, the lines corresponding to \(\nu = 1\) at different redshifts.

4.3 The global, clumped, DM annihilation luminosity per baryon

Having computed the DM annihilation luminosity of a particular halo of mass \(M\) at redshift \(z\), we now return to the problem of computing the global energy injection rate to the IGM per baryon, but considering the clumpiness of the DM on all clustering scales. To do this, we simply consider that, knowing the luminosity of a particular halo (that already includes substructure and adiabatic compression), and knowing the halo mass function (calculated in Section 3.1), the energy output rate density (or luminosity density) due to DM annihilations is given simply by:

\[
\epsilon_{\chi}^{\text{clumped}}(z) = (1+z)^3 \int_{m_{\text{free}}}^{m_{\text{max}}} L_{\chi}(M, z) \frac{dn}{dM}(M, z) dM, \tag{36}
\]

where \(m_{\text{free}}\) is the free-streaming mass, \(m_{\text{max}}\) is the maximum mass of a DM halo, \(L_{\chi}\) is the total DM annihilation luminosity (main halo plus substructure) of a DM halo of mass \(M\) at redshift \(z\), \(\frac{dn}{dM}\) is the previously calculated halo mass function (in comoving coordinates), and the factor \((1+z)^3\) is to convert the comoving density to a proper density. We consider \(m_{\text{max}} = 10^{17} M_\odot\) since the number density of collapsed DM haloes with mass above this limiting mass is negligible.

Finally, the energy injection rate to the IGM per baryon can be directly computed as:

\[
\dot{E}_{\chi}^{\text{clumped}}(z) = \frac{\epsilon_{\chi}^{\text{clumped}}(z)}{n_b(z)}, \tag{37}
\]

where \(n_b(z)\) is the proper number density of baryons at redshift \(z\) (equal to \((1+z)^3 n_{b,0}\) in terms of the present day baryon number density computed in Section 4.1).

It is important to mention that the energy injection rate to the IGM per baryon computed in this way, considers only the contribution of the DM haloes and not of the smooth component. In the case of a zero free-streaming mass, all the DM would be in haloes since there would be no minimum mass scale for the formation of collapsed structures. However, due to the non-zero free streaming mass, all haloes with masses less than that value, that should exist according to the halo mass function, are simply not allowed to form. As the clumped energy injection should be always greater than the smooth energy injection (computed in Section 4.1), and as the contribution due to DM haloes will be declining rapidly with increasing redshift after the truncation of low mass haloes due to the free-streaming limit becoming important, a good approximation is to use as the clumped DM injection rate the greater value among the one calculated in equation (47), and the one calculated in Section 4.1. This is the approach that was used for the case of a 10GeV WIMP.

The results obtained for the clumped energy injection rate per baryon, considering different clustering scenarios and WIMP masses, are shown in Figure 4. We only consider the maximal absorption fraction \((f_{abs} = 1)\), because for
The clumped energy injection rate per baryon \( \dot{E}_\chi(z) \) as a function of redshift, for different scenarios, in thick lines. The dotted curve corresponds to the case with substructure and adiabatic compression for a WIMP mass of 1TeV, the long-dashed curve corresponds to the case with 1TeV WIMP mass and substructure but no adiabatic compression, the dot-dashed curve corresponds to the case of 1TeV WIMP with adiabatic compression but no substructure, and the short-dashed curve corresponds to a WIMP mass of 10GeV with substructure and adiabatic compression. Also included as thin lines are the previous curves of Figure 2, corresponding to the smooth cases, for a 1TeV WIMP (dotted) and a 10GeV WIMP (short-dashed).

Other values, the resulting curves should simply be rescaled by \( f_{abs} \).

It can be seen that the case considering a WIMP with mass of 10GeV is the one that gives the higher injection rate, in agreement with what was discussed above. However, the curve for this case has a cutoff at \( z \sim 50 \) instead of at \( z \sim 80 \) like for the cases considering a 1TeV WIMP. This can be understood as a consequence of the higher free-streaming mass (and thus, higher minimum mass that haloes can have), and the fact that the mass of the haloes that have a characteristic \( \nu = 1 \) (the haloes that have just recently collapsed) decreases with increasing redshift, such that with a higher free-streaming mass, less haloes are allowed to form at higher redshifts. We can see that the presence of adiabatic compression is also very important for boosting the annihilation energy injection. We also note that the curves considering adiabatic compression have a secondary broad peak at lower redshift, besides the peak adjacent to the cutoff. This secondary peak is present due to the existence of \( M_{\text{crit}} \), because for higher redshifts, the fraction of haloes that have \( M > M_{\text{crit}} \) decreases due to the redshift dependence of the halo mass function, and so does the energy output in DM annihilations. The case considering both adiabatic compression and substructure is indistinguishable from the case that considers only adiabatic compression and no substructure, which is in agreement with what was discussed in Section 4.2 with respect to the interpretation of Figure 3.

In the next section, we obtain the clumpiness factor, \( C(z) \), implied by our energy injection calculations, and compare our results with those obtained by Cumberbatch, Lattanzi & Silk (2010).

### 4.4 The clumpiness factor

The clumpiness factor is a useful tool when accounting for DM clustering in global DM annihilation energy injection computations, like the ones required for calculating the heating and ionization of the IGM due to DM (see Section 6.1 and Section 6.2). It simply relates the clumped energy injection to the smooth one by:

\[
C(z) = \frac{\dot{E}_\chi^{\text{clumped}}(z)}{\dot{E}_\chi^{\text{smooth}}(z)}.
\]

We show the values obtained for the clumpiness factor \( C(z) \) in Figure 5. The meaning of the curves are explained in the caption.

In general, it can be seen that for all the cases considered, the clumpiness is a decreasing function of redshift. This is reasonable, because as the redshift increases, the universe will tend to be more homogeneous and uniform, because...
the gravitational collapse and formation of structures would have had less and less time to occur. Also, for all the cases considered, the universe will be practically completely homogeneous by redshift \((z \sim 80)\), except for the case with \(10\text{GeV}^{-1}\) DM WIMP, in which the universe becomes homogeneous at redshift \((z \sim 50)\). Note that in this case, although \(E_{\chi_{\text{damped}}}^{}(z)\) is higher than for the other cases by almost an order of magnitude at redshift zero, \(E_{\chi_{\text{smooth}}}^{}(z)\) is higher as well, and so the clumpiness \(C(z)\) is not. It can also be seen that in agreement with what was discussed before, the adiabatic compression mechanism is important for low redshifts, and the presence of substructure is irrelevant for increasing the clumpiness provided that the adiabatic compression is present (for haloes and subhaloes with \(M > M_{\text{crit}}\)), as discussed in the previous sections.

We show also the clumpiness factors \(C(z)\) obtained by Cumberbatch, Lattanzi 
& Silk (2010), in Figure 5, using different line-styles. The two sets of curves can be readily compared as the definitions of the clumpiness factors are the same. There are, however, differences in the methods for computing them. In particular, Cumberbatch, Lattanzi 
& Silk (2010) compute the halo mass function using the P-S theory, starting from an analytic fit to \(\sigma(M, z)\), and then consider that substructure and sub-substructure follow a pure power-law mass function with exponent \(-2\). We refer the reader to the authors for further detail about their computation procedure. In general, it can be seen that we predict less clumpiness than the authors at low redshift, but comparable clumpiness than the authors at high redshift. This can be understood because the authors also include sub-substructure (sub-subhalo within subhalo), that may be more important near redshift zero, but we include adiabatic compression (and they do not), which may be important at higher redshifts than substructure, but low enough so that the presence of haloes with \(M > M_{\text{crit}}\) is not disfavoured by the redshift dependence of the halo mass function, and at higher redshifts, by the free-streaming mass.

We do not include Sommerfeld enhancement in our calculations, because it is highly dependent on unknown physics (the mass of the new force carrier and the resulting potential, as well as the physical scenarios in which this new force could arise), however, its effect is to directly modify \(<\sigma v>_t\) by a constant factor. As we already mentioned in the Introduction, the presence of Sommerfeld enhancement is controversial and not supported by all observational evidence.

Having computed the DM annihilation luminosity generated within a DM halo for different scenarios, in the next section we study the luminosity received by a DM halo due to the rest of the universe.

### 5 EXTERNAL DM ANNIHILATION ENERGY INJECTION ON A HALO

On previous sections, we have considered particular WIMP models, calculated the DM clustering on all scales explicitly, and computed the DM annihilation luminosity generated within a particular DM halo, and the energy injection per baryon on the IGM. In this section, we calculate the DM annihilation luminosity received by a halo due to the energy generated in the rest of the universe. For this, we will use the annihilation luminosity of a halo that was obtained in Section 4.2, together with the bias factor and the two-point correlation function that characterize the spatial clustering of haloes, and the cosmological radiative transfer formalism as presented in Section 2.1 of Haardt 
& Madau (1996). From the point of view of any given halo, the radiation contributed by all the structures external to it will be characterized in terms of a specific intensity field, and then, the received luminosity will depend on the effective cross-section of this halo to the presence of the radiation field. In the following subsections we first characterize the spatial clustering of haloes by giving the prescriptions used for the computation of the correlation function and the bias, and then, based on these, we obtain the radially clustered halo mass function about any given halo of mass \(M\). Then, finally, we compute the intensity of the radiative background and obtain the external luminosity received by the halo.

#### 5.1 The radial clustering of haloes

We start the discussion of this section by considering the two-point correlation function. The correlation functions we will consider take into account redshift evolution.

Although the correlation function is completely specified by the power spectrum, the \(P(k)\) that we calculated is valid in the linear theory approximation, and thus, it is not a good estimator of the DM halo spatial clustering on small scales. Therefore, we will use a standard analytic fit to the zero redshift correlation function, as a pure power law, with its exponent and normalization motivated by observations (for example the Sloan Digital Sky Survey, as analysed by Zehavi et al. (2011)) and simulations:

\[
\xi_0(r) = \left( \frac{r}{6h^{-1}\text{Mpc}} \right)^{-1.8},
\]

where \(r\) is in comoving Mpc.

Now, we proceed to explain the bias factor, that represents the excess in clustering of haloes above different mass scales. DM haloes with different masses (above a mass scale \(M\)), provide biased measurements of clustering. The bias factor of mass scale \(M\) can be defined by:

\[
b^2(M, z) = \frac{\xi_{\text{DM}}(r, z)}{\xi_M(r, z)},
\]

with \(\xi_{\text{DM}}(r, z)\) the correlation function of the DM haloes with masses above \(M\), and \(\xi_M(r, z)\) the unbiased correlation function (labelled simply as \(\xi(r, z)\), and considered previously). This bias factor is independent of \(r\) in a first order approx.

We adopt the bias factor given by Sheth, Mo 
& Tormen (2001) that takes into account ellipsoidal collapse and fits well numerical simulation measurements.

Finally, we consider the radial halo mass function. This mass function, that will be given by the mass function computed with P-S theory, but considering the clustering of DM around the particular halo of mass \(M\), can be defined as
\[
d\frac{dn}{dm}(M, m, z, r) = \left(1 + b(\nu(M, z))b(\nu(m, z))\xi(r, z)\right) \times \frac{dn}{dm}(m, z) \tag{41}
\]

Now, having characterized the radial clustering of haloes about a given halo, we proceed to compute the radiation field around the halo and the power absorbed by it.

### 5.2 The radiative transfer of the DM annihilation output and the external luminosity

In section 2.1 of [Haardt & Madau (1996)]{#15}, the solution to the problem of radiative transfer, in terms of the isotropic specific intensity, in the context of a dynamic cosmological background was presented. The authors considered the radiative transfer equation,

\[
\left(\frac{\partial}{\partial t} - \nu \frac{\partial}{\partial \nu}\right) J = -3\frac{\dot{\nu}}{a} J - c\kappa J + \frac{c}{4\pi} \epsilon, \tag{42}
\]

with formal solution

\[
J(\nu_0, z_0) = \frac{1}{4\pi} \int_{z_0}^{\infty} dz' \frac{\partial}{\partial z'} \left(\frac{1 + z_0}{1 + z'}\right)^3 \epsilon(\nu, z) e^{-\tau_{eff}(\nu_0, z_0, z)}, \tag{43}
\]

where \(J\) is the specific intensity (in units of \(\text{ergs cm}^{-2}\text{s}^{-1}\text{Hz}^{-1}\text{sr}^{-1}\)), \(\epsilon\) is the specific emissivity, \(\nu\) is the redshifted frequency of the radiation, \(\frac{\partial}{\partial \nu}\) is the distance travelled by a single photon of radiation per redshift interval and \(\tau_{eff}\) is the effective optical depth for the radiation travel.

In our case, because we are considering a model-independent WIMP scenario, we do not have a specific spectral energy distribution (SED) for the radiation resulting from the annihilation events. Therefore, instead of using a frequency dependent specific intensity, we modify this formalism to consider an intensity integrated in frequency. In this case, the emissivity will be the usual power density, which is the distance travelled as a secondary photon of radiation per redshift interval and \(\tau_{eff}\) is the effective optical depth for the radiation travel.

Thus, we write

\[
J(z_0) = \frac{1}{4\pi} \int_{z_0}^{\infty} dz' \frac{\partial}{\partial z'} \left(\frac{1 + z_0}{1 + z'}\right)^3 \epsilon_{\nu}(z) e^{-\tau_{eff}(\nu_0, z', z_0, z)}, \tag{44}
\]

Now, we give the prescriptions used for the terms included in this equation.

For the line element, we use

\[
\frac{dl}{dz} = c \frac{dLBT(z)}{dz}. \tag{45}
\]

Also, in order to take into account the radial clustering statistics about a halo of mass \(M\), we compute the comoving distance travelled by a photon emitted at redshift \(z\) and received at \(z_0\), as

\[
r(z_0, z) = \frac{c}{H_0} \left(u \left(\frac{1}{1 + z}\right) - u \left(\frac{1}{1 + z_0}\right)\right), \tag{46}
\]

where

\[
u(x) = \int_0^1 \frac{d\chi}{\chi^2 \left(\frac{\Omega_M}{\chi^2} + \Omega_\Lambda + (1 - \Omega_M - \Omega_\Lambda)\right)^{1/2}}. \tag{47}
\]

Now, we write the emissivity of the annihilation luminosity in the vicinity (at a comoving distance \(r\)) of a halo of mass \(M\) as

\[
\epsilon_{\nu}(M, z, r) = \epsilon_{\text{esc}} \int_{M_{\text{free}}}^{\infty} \frac{dn}{dm}(M, m, z, r)(1 + z)^3 L_{\nu}(m, z), \tag{48}
\]

and it already considers the correlation function and the bias of DM haloes. Also, \(L_{\nu}(m, z)\) corresponds to the total annihilation luminosity generated by a halo, considering both the main halo and the substructure therein.

Now, for the effective optical depth, we use

\[
\tau_{eff}(z_0, M, z) = \int_{M_{\text{free}}}^{\infty} \int_{z_0}^{\infty} dm' \frac{dn}{dm'} (M, m, z, r) (1 + z')^3 \sigma_{\text{bary}}(m), \tag{49}
\]

where we specify \(\sigma_{\text{bary}}\) as follows.

The baryonic cross-section of a halo to the DM annihilation radiation was computed as

\[
\sigma_{\text{bary}}(m) = f_{\text{bary}} \left(\frac{m}{m_p}\right) \left(1 - \frac{Y_p}{2}\right), \tag{50}
\]

where \(f_{\text{bary}} = \frac{\Omega_b}{\Omega_m} = 0.17\) is the fraction of baryons by mass, \(Y_p = 0.26\) is the primordial abundance of He by mass and \(\sigma_T\) is the Thompson cross-section. This corresponds to considering the total cross section of a DM halo as the number of electrons present in the halo times the Thompson cross-section. Although this prescription is crude, it is also completely model-independent. The implicit assumption here is that the photons of the radiation are energetic enough so that their frequency is higher than the frequency of typical atomic transitions in neutral hydrogen and Helium (\(\sim 10\) eV), so that we are away from the Rayleigh regime and therefore bound electrons can be considered as free; and also that they are soft enough so that their energies are small with respect to the rest-mass energy of the electrons.

The first assumption is certainly true, whereas the second assumption seems dubious because the mass of an electron is on the MeV scale whereas the mass of a WIMP is on the GeV-TeV scale. However, one should keep in mind that the radiation background received by the halo is not necessarily the same radiation that was obtained as a secondary product of the annihilations, because the photons have been redshifted as they travelled through the universe, and also, they could have been scattered off other haloes before. In any case, the best reason for considering this prescription is its model-independence, and also the fact that, because for all considered halo masses and redshifts we are always in the optically thin regime (\(\tau_{eff} << 1\)), the absorbed external luminosity will be simply proportional to this cross-section, and therefore, the use of the Thompson cross-section as the cross-section of each individual electron amounts to obtaining an upper bound on the external luminosity, because for
higher energy photons, the correct Klein-Nishina correction factor only decreases the cross-section.

Finally, the specific intensity will be given by

\[ J(z_0, M) = \frac{1}{4\pi} \int_{z_0}^{\infty} dz \frac{dI}{dz} \left( \frac{1+z}{1+z} \right)^4 \epsilon_x (M, z, r(z, z)) \times \epsilon^{-\tau_{eff}(z_0, M, z)} . \] (51)

Note that we have already taken into account the radial dependence of the clustering statistics about a given halo of mass M. Also, we have taken into account the fact that haloes of different mass have different luminosities and different biases with respect to the target halo, by performing an extra mass integration in the computation of \( \epsilon_x \).

Now, given the specific intensity of the radiation field, we compute the flux that is incident on the surface of the halo, and the energy that is absorbed by it. By surface, we mean that we assign a radius to the halo that corresponds to \( R = \sqrt{M_{\text{bary}}(z)} \).

The flux incident on any surface element of the halo will be given by

\[ F = \int_{\Omega} J \cos \theta d\Omega , \] (52)

where in this case, \( \Omega \) corresponds to the solid angle of the half-sphere, because any surface element receives radiation from the entire half-space that corresponds to the outside of the halo.

Thus, because \( J \) is isotropic, the flux can be written as

\[ F = \pi J . \] (53)

Finally, the power absorbed by a halo (our external luminosity) will be given by

\[ L_x^{\text{ext}} = \epsilon_{\text{abs}} 4\pi^2 R^2 J , \] (54)

where \( R \) is the radius of the halo. This can also be written as

\[ L_x^{\text{ext}} = 4\pi \epsilon_{\text{abs}} \sigma_{\text{bary}}(M) J . \] (55)

We see that the final result is proportional to \( \sigma_{\text{bary}}(M) \), and therefore, for the reasons already discussed, it should only be considered as an upper bound on the external luminosity. Also, we choose \( \epsilon_{\text{abs}} = 0.9 \), because an absorption fraction of 0.9 means that 10% of the photons produced inside DM haloes manage to reach the IGM, and therefore, this choice is somewhat equal to choosing \( f_{\text{abs}}(z) = 0.1 \) for the energy injection per baryon in Section 4.3, as motivated by Cumberbatch, Lattanzi & Silk (2010) and Natarajan & Schwarz (2009), although both fractions are not directly comparable.

In Figure 6, we show our results for the external DM annihilation light-to-mass ratio (defined as \( L_x^{\text{ext}}/M \)) as a function of halo mass at different redshifts, considering only the \( \Lambda \)CDM cosmology, in the case of a WIMP with mass 1TeV and considering both substructure and adiabatic compression.

In general, comparing these results with the ones obtained for the DM luminosity to mass ratio generated within a particular halo (see Figure 3a), one can see that the external annihilation luminosity is lower than the internal annihilation luminosity at all redshifts and that the difference is of at least 2 orders of magnitude, as in the case of \( z = 2 \) and for masses less than \( 10^7 M_\odot \). Therefore, considering that the luminosity computed here is only an upper bound, we can conclude that the external DM annihilation luminosity is not relevant as a feedback mechanism. We can also see that the curves for the received luminosity imply an almost constant light to mass ratio at each redshift, which is reasonable considering that the halo cross-section is proportional to the number of electrons, and thus proportional to the mass.

6 GLOBAL IGM HEATING, IONIZATION AND 21 CM SIGNAL DUE TO DM

Different authors have studied the effects of DM annihilations (or decays) on the global properties of the IGM, due to their extra energy injection (for example, see Pierpaoli (2004); Mapelli, Ferrara & Pierpaoli (2006); Ripamonti, Mapelli & Ferrara (2007a); Valdés et al. (2007); Natarajan & Schwarz (2009); Cumberbatch, Lattanzi & Silk (2010); Furlanetto, Oh & Pierpaoli (2006); Chuzhoy (2008) among others). In this section, we study the effects on the IGM of the energy injection found on Section 4.3. We divide the section into different subsections, each corresponding to one of the different physical quantities of interest that are affected by this extra energy injection. We explain the formalism and the definition of the quantities, as well as how to compute them, in the beginning of the corresponding subsections, and then we show the results we obtained and compare them.
with different authors. The first two sections (the ionization fraction and the IGM temperature) refer to general properties of the IGM, whereas the next sections have to do with the signal of this extra-heating on the 21 cm line of the hyperfine structure of neutral hydrogen (HI).

We start, in the next subsection, with the computation of the IGM ionization fraction.

### 6.1 The IGM ionization fraction

The ionization fraction ($x_e$) is defined as the ratio of free electrons to hydrogen atoms $x_e = n_e / n_H$. Following Valdes et al. (2007), the variation of the ionization fraction is given by:

$$\frac{dx_e}{dz} = \frac{1}{H(z)(1+z)} \left[ R_e(z) - I_e(z) - I_\chi(z) \right], \quad (56)$$

where $I_e$ and $R_e$ are the standard ionization and recombination rates per baryon, and $I_\chi(z) = \chi(z) \frac{\dot{\chi}}{\chi}$ is the contribution to the ionization rate due to the energy injection of DM annihilation. $E_\chi$ is the energy injection rate per baryon computed in Section 4.3. $E_0$ is the ionization energy of Hydrogen (13.6 eV), and $\chi(z)$ is the ionization efficiency and is given by $\chi(z) = \frac{1}{1 + \frac{E_\chi}{E_0}}$ (i.e., if all matter is ionized, no energy is used for ionization, and if all matter is neutral, $1/3$ of the injected energy goes to ionizations). The factor $\frac{1}{H(z)(1+z)}$ is simply a change of variables, because the differential of cosmic time $dt$ and the redshift differential $dz$ are related by $dt = \frac{-dz}{H(z)(1+z)}$.

We define the DM annihilation contribution to the ionization fraction as $\delta x_e = x_e - x_e^{no-DM}$, where $x_e^{no-DM}$ is the ionization fraction in the standard scenario. It is important to mention that this standard ionization fraction is a consequence of the physics of recombination from the epoch of the last scattering surface, and is simply the relic free-electron abundance from the freeze-out of the recombination reaction. Thus, the $x_e^{no-DM}$ does not account for the reionization of the universe by Population III stars, that according to Planck 2013 (Ade et al. 2013), occurred at redshift 11.35. With this definition of $\delta x_e$, we have:

$$\frac{d\delta x_e}{dz} = \frac{-I_e(z)}{H(z)(1+z)}. \quad (57)$$

We solve this differential equation, using the $x_e^{no-DM}$ given in Figure 3 of Mapelli, Ferrara & Pierpaoli (2006), by the Euler method, using as boundary condition that $\delta x_e(z = 1000) = 0$ (at redshift $z > 1000$, as we are near the recombination epoch, the ionization fraction is near 1 and is almost entirely due to the standard cooling scenario).

We show the obtained ionization fraction ($x_e(z)$) curves in Figure 7 and in Figure 8. The first figure shows the $x_e(z)$ curves obtained for the different DM clustering scenarios and WIMP masses, but considering the lumped DM case, and with an absorption fraction of $f_{abs} = 1$ (more details on the caption of the figure). The second figure shows the ionization fractions, assuming a WIMP mass of 1TeV and substructure and adiabatic compression, but for different values of the absorption fraction ($f_{abs} = 0.1$ or 0.01), or simply not considering DM clustering at all, and using the homogeneous DM energy injection computed previously.

It can be seen from Figure 7 that the ionization fraction evolves very differently from the case with no DM energy injection, as long as the DM clumpiness is considered. Also, for the case of a WIMP with a mass of 10GeV, the ionization fraction can be increased up to $x_e \sim 0.03$ at redshift zero, whereas for the other cases that consider a 1TeV WIMP, this partial re-ionization can only bring $x_e \sim 0.005$ at redshift zero, hence leading to a difference of one order of magnitude of the final ionization fraction. Note, however, that the computation of the ionization fraction does not account for the PopIII stars (for a review, see Section 4 of Barkana & Loeb (2001)), that are thought to have re-ionized the universe around redshift $z \sim 11.35$ according to the results by Planck (Ade et al. 2013); thus, this results are only reliable down to that redshift at most. However we use a more conservative limit of $z = 20$ in the plots, to account for the fact that the formation of the first stars started earlier than this quoted redshift of reionization. Note, nonetheless, that we should keep in mind that the ionization process is by no means an instantaneous process, and that this redshift
is irrelevant with respect to the presence of adiabatic compression.

From Figure 8, it can be seen that even without considering DM clumpiness, the case with a 10GeV WIMP considerably modifies the ionization fraction of the universe. It has the effect of raising the relic ionization at redshift 0 (without accounting for stellar populations) by a factor of \( \sim 20 \). Also, the evolution of the ionization fraction is particularly sensitive to the absorption fraction \( f_{\text{abs}} \), such that reducing \( f_{\text{abs}} \) from 1 to 0.1 results in a completely different evolution for \( z < 10 \) (and in a ionization fraction at \( z=0 \) an order of magnitude lower). The cases with \( f_{\text{abs}} = 0.1 \) and no clumpiness give evolutions of the ionization fraction that are very similar to the case were no DM energy injection is present.

### 6.2 The IGM kinetic temperature

We now consider the kinetic temperature of the IGM. Following [Valdés et al. (2007)](Valdes2007), the change in this temperature is given by:

\[
(1 + z) \frac{dT_h}{dz} = \frac{l_x}{H(z)(1 + Y + x_e)} (T_h - T_{\text{CMB}}) - \frac{2 \chi_h E_h}{3 k_h H(z)(1 + Y + x_e)},
\]

where \( l_x = \frac{(8 \sigma_T_n R^2_{\text{CMB}})}{\langle 3m_n c^2 \rangle} \) is the analogue to a thermal conductivity between the CMB and the IGM, \( \chi_h \) is the heating efficiency and is given by \( \chi_h = \frac{4 \pi^2}{15} \) and \( Y \) is the helium fraction by mass (taken as \( Y = 0.26 \)).

It is expected (see for example [Cumberbatch, Lattanzi & Silk (2010)](Cumberbatch2010)) that the IGM kinetic temperature tracks the CMB temperature at high redshifts, so we used the CMB temperature at redshift \( z = 1000 \) as the boundary condition for the IGM kinetic temperature, and solved the differential equation by the Euler method. The interpretation of this cooling equation is straightforward, as the first term on the right hand side corresponds to the adiabatic cooling of the IGM in absence of heat exchange or energy injection, and proceeds as \( T \propto (1 + z)^2 \). The second term corresponds to the heat exchange between the CMB and the IGM, and the third term allows for the extra heating due to the DM annihilation energy injection.

We show the obtained IGM kinetic temperature curves in Figure 9 and in Figure 10. The first figure shows the temperatures obtained for the different DM clustering scenarios and WIMP masses, but considering the clumped DM case, and with an absorption fraction of \( f_{\text{abs}} = 1 \) (more details on the caption of the figure). The second figure shows the temperatures, assuming a WIMP mass of 1TeV and substructure and adiabatic compression, but for different values of the absorption fraction \( f_{\text{abs}} = 0.1 \) or 0.01, or simply not considering DM clustering at all, and using the homogeneous DM energy injection computed previously.

In can be seen on Figure 9 that for all the cases that consider the clumpiness of the DM (assuming \( f_{\text{abs}} = 1 \)), the cooling is very different from the standard adiabatic scenario. Significant re-heating occurs in the case of a 10GeV WIMP, where the IGM may be heated even to temperatures of \( \lesssim 1000K \) by redshift zero (although, remember that this

![Figure 8](image-url)  
*Figure 8.* The ionization fraction of the IGM for different cases, assuming a WIMP mass of 1TeV and substructure and clustering, or no DM clumpiness at all. Our results are shown as thick lines. The different line-styles of the thick lines correspond to: (dotted) the standard clustering scenario considering both substructure and adiabatic compression with a 1TeV WIMP and \( f_{\text{abs}} = 1 \), (short-dashed) the same clustering scenario but considering \( f_{\text{abs}} = 0.1 \), (long-dashed) the same clustering scenario but with \( f_{\text{abs}} = 0.01 \), (short-dot-dashed) the same WIMP mass but considering no DM clumpiness, (long-dot-dashed) a 10GeV WIMP with no DM clumpiness, and (solid) the standard cooling scenario, with no DM energy injection. For comparison, we also include the results obtained by Mapelli, Ferrara & Pierpaoli (2006) for the case of the 100GeV SUSY neutralino, taken from the bottom panel of Figure 3 of the authors and we show them as thin lines. The thin solid line corresponds to \( < \sigma v > = 2 \times 10^{-26} \text{cm}^3 \text{s}^{-1} \) and the thin dotted line corresponds to \( < \sigma v > = 1 \times 10^{-24} \text{cm}^3 \text{s}^{-1} \). Also included is the line corresponding to \( z = 20 \), that indicates when the effects of PopIII stars may become important.

Corresponds only to the best fit for the six-parameter base \( \Lambda \text{CDM} \) model under the assumption of instantaneous reionization (as explained in [Ade et al. (2013)](Ade2013), but that large regions of neutral Hydrogen may have still been present up to \( z > 7 \), with a volume filling factor of \( \gtrsim 0.1 \) (see section 5.1 of [Ripamonti, Mapelli & Zaroubi (2008)](Ripamonti2008)). Also, we note that as discussed by [Ripamonti, Mapelli & Zaroubi (2008)](Ripamonti2008) and by ?, the reionization may have been caused not only by the first stars, but also by the X-ray emission from accretion onto early Black Holes.

It can also be seen in the figure that adiabatic compression is considerably less important than WIMP mass in determining the evolution of the ionization fraction, and that it becomes relevant only for \( z \sim 10 \). The case with substructure but no adiabatic compression results in \( x_e \sim 0.0003 \) at \( z = 0 \), whereas the case with adiabatic compression but no substructure is indistinguishable from the case that considers both processes. Therefore, in agreement with what was mentioned in previous sections, the presence of substructure...
cooling scenario does not account for energy injection due to stellar components and early black holes, and so it is reliable only down to redshift \( z \sim 20 \). Also, adiabatic compression is more important for the cooling history than substructure, and the extra heating due to this compression can amount to almost two order of magnitude of difference by redshift zero, but the difference is relevant only for \( z \lesssim 10 \).

Also, it can be seen on Figure 10 that the absorption fraction is critical for determining the particular cooling history, and that the low-redshift heating due to DM annihilations can vary enormously with different choices for \( f_{\text{abs}} \). It can be seen that, in general, in the cases considering adiabatic compression, the IGM kinetic temperature reaches a quasi-stationary state, and the final temperature attained is almost proportional to the absorption fraction.

We also include the results for the ionization fraction and IGM kinetic temperature obtained by Mapelli, Ferrara & Pierpaoli (2006), in order to compare them with our results.

It can be seen that for the cases with no DM clumpiness, we reproduce the results obtained by the authors. Also, the boost they considered in the cross section (for example, accounting for Sommerfeld enhancement), has almost exactly the same effect than lowering the mass of the WIMP particle. Furthermore, we find significant effects not accounted for by the authors when considering the clumpiness of the DM on all its levels, but only at \( z < 10 \) for 1TeV WIMP or at \( z < 30 \) for 10GeV WIMP.

6.3 The neutral hydrogen spin temperature

We now study the 21 cm signal from neutral hydrogen, and how it is modified by DM annihilation. We follow Cumberbatch, Lattanzi & Silk (2010) and Valdés et al. (2007) in the explanation of the physics and the processes involved, and refer the reader to the above papers for further details.

The 21 cm line corresponds to the transition energy between the singlet and triplet hyperfine levels \((n = 0, j = 0\) and \(n = 0, j = 1\) of the hydrogen atom. The spin temperature of the neutral hydrogen is defined as the temperature that relates the populations of singlet and triplet states:

\[
\frac{n_1}{n_0} = 3 e^{-\frac{T_s}{T^*}},
\]

where \( T_s \) is the spin temperature, \( T^* \) is the equivalent temperature of the 21 cm transition (corresponding to \( T^* = \frac{hc}{\lambda k_B} \), where \( \lambda = 21\text{cm} \)), \( n_1 \) is the triplet state occupation number, \( n_0 \) is the singlet state occupation number and the factor of 3 is the degeneracy of the triplet state (i.e., 3).

In general, the spin temperature is a weighted average of the IGM kinetic temperature and the CMB temperature. In absence of additional heating mechanisms, both the IGM kinetic temperature and the spin temperature are coupled to the CMB at high redshifts \((z \gtrsim 300)\), whereas at lower redshifts, the IGM cools faster than the CMB. Thus, the spin temperature will be lower than the CMB and the IGM will be visible in 21 cm absorption.

There are two main mechanisms that couple the spin temperature and the kinetic temperature: Spin exchange collisions and Ly-alpha pumping (also called the Wouthuysen-Field (W-F) effect). Spin exchange collisions...
correspond to the collisions of neutral hydrogen (responsible for the 21 cm line) with other neutral hydrogen atoms, with electrons or with protons. These collisions will tend to thermalize the populations of both levels (singlet and triplet) to the kinetic temperature of the IGM, and thus will couple $T_s$ with $T_k$. Lyman-alpha pumping, on the other hand, corresponds to the mixing (and thus thermalisation) of the singlet and triplet levels by Lyman-alpha transitions (from energy levels 1 to 2 and vice-versa) to the 2p level. Therefore, the mixing of hyperfine states through the W-F effect will depend on the intensity of the Lyman-alpha photon field, which in turn will depend on the kinetic temperature of the IGM and on the direct energy injection due to DM WIMPS. Thus, the W-F will also couple $T_s$ with $T_k$.

At low redshift ($z < 50$), the IGM density is sufficiently rarefied so as to make the spin exchange collisions too improbable, and thus $T_s$ will tend to follow the temperature of the CMB again.

As the mechanisms coupling $T_s$ and $T_{k IG M}$ depend on the ionization fraction $x_e$, and the DM energy injection rate per baryon $E_x$, it is expected that the presence of an annihilating WIMP will alter the evolution of $T_s$ with redshift, and thus, the detection of this signal may help to discriminate between different WIMP scenarios.

The spin temperature at redshift $z$ can be computed as:

$$ T_s = \frac{T_{CMB} + y_0 T_k + y_e T_k}{1 + y_0 + y_e}, \quad (60) $$

where $y_0$ is the W-F coupling, and $y_e$ is the collisional coupling to the IGM temperature. The coupling coefficients can be computed as:

$$ y_0 = \frac{P_{10} T_s}{A_{10} T_k}, \quad (61) $$

and

$$ y_e = \frac{C_{10} T_s}{A_{10} T_k}, \quad (62) $$

where $A_{10} = 2.85 \times 10^{-15} \text{s}^{-1}$ is the spontaneous emission coefficient for the transition, and $P_{10}$ and $C_{10}$ are the radiative and collisional transition coefficients respectively (they are rates, with units of $[\text{s}^{-1}]$). The collisional transition coefficient is given by:

$$ C_{10} = k_{10} n_{HI} + n_e \gamma_e, \quad (63) $$

where $n_{HI}$ and $n_e$ are the neutral hydrogen number density and the free electron number density respectively, and $k_{10}$ and $\gamma_e$ are the specific (one particle) transition coefficients for collisions with neutral hydrogen atoms and electrons respectively. We have neglected the contribution due to collisions with protons, following Cumberbatch, Lattanzi & Silk (2010). The number densities $n_e$ and $n_{HI}$ are easily obtained in terms of the above-defined ionization fraction $x_e$, the hydrogen and helium abundances by number ($f_H$ and $f_{Hee}$) and the baryonic density of the universe $n_b$, also defined previously. We use the following formulas for the coefficients $k_{10}$ and $\gamma_e$:

$$ k_{10} = 3.1 \times 10^{-11} T_k^{0.357} \exp(-32 K/T_k), \quad [\text{cm}^3 \text{s}^{-1}] \quad (64) $$

and

$$ \log(\gamma_e / \text{cm}^3 \text{s}^{-1}) = -9.607 + 0.5 \log(T_k/1 K) \exp(-[\log(T_k/1 K)]^{1.5}/1800), $$

$$ 1K < T_k < 10^5 K, \quad (65) $$

$$ \log(\gamma_e / \text{cm}^3 \text{s}^{-1}) = -9.607 + 0.5 \log(T_k/1 K), T_k < 1K, \quad (66) $$

$$ \gamma_e(T_k > 10^4 K) = \gamma_e(10^4 K). \quad (67) $$

The fit to $k_{10}$ is given by Kuhlen, Madau & Montgomery (2006) and the fit to $\gamma_e$ is given by Liszt (2001).

We now calculate $P_{10}$ as:

$$ P_{10} = \frac{16 \pi J_0 \sigma_\alpha}{27 \hbar \nu_\alpha}, \quad (68) $$

where $J_0$ is the isotropic specific intensity of the radiation field in the Ly-alpha line, $\sigma_\alpha$ is the Ly-alpha transition cross-section in the monochromatic approximation, and $\nu_\alpha$ is the frequency of the Ly-alpha transition photon. According to Laor (1996), the $\sigma_\alpha$ monochromatic cross-section is given by:

$$ \sigma_\alpha = \frac{\pi \epsilon^2}{m_e c} f_{12} f_{se} \frac{n_1}{n_1 + n_2}, \quad (69) $$

where $f_{se}$ is the stimulated emission correction and is given by:

$$ f_{se} = (1 - e^{-\frac{\hbar \nu_\alpha}{k_B T_k}}), \quad (70) $$

the first and second energy levels population ratio $\frac{n_2}{n_1}$ of the hydrogen atom is given by:

$$ \frac{n_2}{n_1} = \frac{g_2}{g_1} e^{-\frac{\hbar \nu_\alpha}{k_B T_k}}, \quad (71) $$

the $n-th$ level degeneration is given by $g_n = 2n^2$, and $f_{12}$ is the oscillator strength of the transition ($f_{12} = 0.416$). Note that and the factor $\epsilon^2$ in $\sigma_\alpha$ should be replaced by $\frac{\epsilon^2}{4 \pi \sigma_0}$ when using MKS units.

Finally, the specific intensity of the Ly-alpha radiation field is given by:

$$ J_\alpha(z) = \frac{n_b^2 h c}{4 \pi H(z)} \left[ x_e x_p \alpha_{eff}^{eH} + x_e x_{HI} \gamma_e h \right] \left[ \chi \right] \frac{\dot{E}}{n_b \hbar \nu_\alpha}, \quad (72) $$

where $x_p = (1 - x_{HI})$ is the proton fraction with respect to the total hydrogen number density, $\alpha_{eff}^{eH}$ is the radiative recombination coefficient to the $n = 2, l = 1$ level, given in Table 1 of Pengelly (1964), $\gamma_e h$ is a collisional excitation
rate of HI atoms involving electrons, and \( \chi_\alpha = \chi_e^2 \), where \( \chi_e \) is the excitations efficiency given by \( \chi_e = 2 - 2z \) (here, it can be seen that \( \chi_e + \chi_i + \chi_h = 1 \), and thus, all the energy injected to the IGM by DM annihilation goes to either heating, ionization or excitation). For \( \gamma_{eH} \), we use:

\[
\gamma_{eH} = 2.2 \times 10^{-8} \exp \left[-11.84/(T/10^4 K) \right] \quad [cm^3 s^{-1}],
\]

(73)
given by Cumberbatch, Lattanzi & Silk (2010).

Now, we can explicitly compute the spin temperature \( T_s \), as a function of \( z \), and use it in the following calculations.

### 6.4 The 21 cm differential brightness temperature

Having obtained the spin temperature \( T_s \), we proceed to compute the 21 cm differential brightness temperature \( \delta T_b \), defined as:

\[
\delta T_b = \frac{T_s - T_{CMR}}{1 + z} \tau (z),
\]

(74)
where \( \tau \) is the optical depth of the neutral IGM at 21(1+z) cm, and can be calculated as:

\[
\tau (z) = \frac{3c^3 h A_{10}}{2 \pi k B_0^2 T_s(z) H(z)} \eta_{HI}(z),
\]

(75)
where \( \nu_0 = 1420 MHz \) is the frequency corresponding to the 21 cm transition in the rest frame of the source. The differential brightness temperature is important because it gives a difference in brightness between the CMB and the neutral hydrogen averaged in the beam of the instrument, and will be given by the above formula when the beam is large enough that the signal equals the average global one. Thus, this is the quantity that can be measured observationally in the low frequency range, and will be used for the next section.

### 6.5 The difference of differential brightness temperatures

Finally, we compute the difference in differential brightness temperature, defined as:

\[
\Delta \delta T_b (z) = |\delta T_b (z) - \delta T_{b,0}(z)|,
\]

(76)
where \( \delta T_{b,0}(z) \) corresponds to the differential brightness temperature obtained without the extra energy injection by DM annihilations. According to Valdés et al. (2007), this signal could be measured with current and future radio telescopes, such as LOFAR, 21CMA, MWA and SKA, in principle, up to the \( \approx 1 mK \) level at redshifts where baryonic energy injection processes (such as PopIII star formation) are not important (\( z \approx 20 \sim 30 \)). However, the authors explain that the various foregrounds (i.e., Galactic free-free and synchrotron emission, unresolved extragalactic radio sources, free-free emission from ionizing sources, synchrotron emission from cluster radio haloes and relics) are much stronger than the cosmological signal, and are also very difficult to remove. The authors also explain some possible methods for effectively removing the foregrounds and for clearly identifying the extra energy injection signal; however, that is beyond the scope of this work and we refer the reader to their paper.

In Figure 11 and in Figure 12, we show the results we obtained for the \( \Delta \delta T_b(z) \) curve, for the different cases we considered. In Figure 11 we assume the maximal absorption fraction \( f_{abs} = 1 \), and consider the different DM clustering scenarios and WIMP masses. In Figure 12, we consider the standard scenario for the DM clustering (adiabatic compression with substructure), and we vary the absorption fraction, the WIMP mass, or consider the perfectly smooth case (for more details about the plots, see their captions).

It can be seen, from Figure 11, that the most important factor that determines the \( \Delta \delta T_b \) signal is the WIMP mass. For a 10GeV WIMP, the signal is around \( \sim 40mK \) at redshift 100, very much above the \( \sim 1mK \) sensitivity of current and future instruments quoted by Valdés et al. (2007). However, the authors also mention that current and future radio frequency arrays will probe up to a minimum frequency of 30MHz, thus corresponding to a 21cm emission at \( z \approx 46 \). Therefore, the sensitivity window should be considered as \( \Delta \delta T_b > 1mK \), and also \( \sim 20 < z < 46 \). The 10GeV case is barely within this window, in the range of \( z \) between \( \sim 37 \) and 46 (with an amplitude of \( \Delta \delta T_b \) that varies from \( 1mK \) to
It can be seen from the comparison between the curves given by the authors and ours, that the dynamical range of the values for the obtained $\Delta \delta T_b(z)$ curves are similar, but that the particular shapes of the curves are different. This may be because the authors use an analytic approximation to $\sigma(M, Z)$, compute $f_{abs}$ self-consistently and assume a pure power-law shape for the substructure, without considering time evolution of substructures within particular haloes, and so their curves will be more soft, having less numerical noise from the integration and equation solving procedures. Also, all our curves are peaked around $z \sim 100$ (with a secondary peak in the cases with $10 GeV$ WIMPs), whereas the results from the authors are peaked around $z \sim 50 - 60$.

7 CONCLUSIONS

In this paper, we develop formalisms for calculating the generated and received DM self annihilation luminosity of individual halos, as well as the energy injection rate per baryon in the universe as a whole, accounting for the clustering of the DM at all scales but without considering Sommerfeld enhancement. We also compute the effects on the IGM and on the 21cm signal of this extra energy injection due to the DM.

It is important to mention that we take a numeric approach for performing the calculations of the barionic clumpiness, instead of the analytic approach used by previous authors. Also, we incorporate adiabatic compression, which had not been done so far. The prescriptions used for the computations are the primordial power spectrum, the transfer function, the window function, the substructure mass function and its corresponding time evolution, the computation of the NFW parameters as a function of halo mass and redshift and the computation of the adiabatically compressed density profile for each halo, considering the critical mass criteria as given by [Hegger et al. (1996)].

We now give some general conclusions about the results obtained.

As discussed in Section 4.2, the luminosity generated within individual halos, considering $1 TeV$ WIMPs, is not enough to self-regulate the star formation of the halo, and thus, it cannot contribute significantly to the standard feedback mechanisms ( AGN feedback for haloes of mass greater than $10^{12} M_\odot$ and SNe feedback for haloes of mass smaller than $10^{10} M_\odot$).

DM clustering in general is determinant in order to assess the effects of DM annihilation on the IGM and for computing the energy injection per baryon. Evidence of this are the clumpiness factors obtained in Section 4.4, which can be in excess of $10^6$ at redshift zero (the energy injection rate is thus boosted by this factor with respect to the case of a perfectly smooth universe). Also, in favour of the importance of clustering are the results obtained for the ionization fraction $x_e(z)$ and the IGM kinetic temperature $T_K(z)$ discussed in Section 6.1 and Section 6.2, where it is clearly seen that, for the cases assuming a perfectly homogeneous universe, and for a WIMP mass of $1 TeV$, the obtained evolutions of these quantities are virtually indistinguishable from the standard case with no energy injection due to DM. For example, not accounting for substructure, led [Mapelli, Ferrara]...
to conclude that the effects of a heavy ($\sim 100\, \text{GeV}$) neutralino on the IGM were negligible, and it was shown by Cumberbatch, Lattanzi & Silk (2010) that this may not be the case by accounting for the clustering of DM. However, we found that even for our maximally clustered case (Adiabatic compression plus substructure), for a 1$\text{TeV}$ WIMP, the deviations with respect to the smooth case only become significant at $z \sim 10 - 20$.

We find that the presence of annihilating and clumped DM may result in significant deviations in the evolutions of the inferred temperature and ionization fraction for $z \lesssim 20$. For example, in the scenario that takes into account the clustering of DM at all levels, as well as the adiabatic compression, a 1$\text{TeV}$ WIMP with thermal relic annihilation cross-section may increase the temperature of the IGM to the $\sim 100\, \text{K}$ level, or contribute to the ionization fraction rising to a value of $\sim 5 \times 10^{-3}$ at redshift zero (not considering PopIII stars). We conclude that although the DM annihilation can have a significant impact on the inferred properties of the IGM, it cannot be regarded as an alternative reionization scenario, because even in the case of a 1$\text{GeV}$ WIMP (with lower mass than favoured by direct and indirect detections and accelerator constraints), the value of $x_e$ at $z < 100$ does not rise over 0.03. This is in agreement with the results obtained by Mapelli, Ferrara & Pierpaoli (2006), where for all the different models of decaying or annihilating DM considered by the authors, they obtained values of $x_e$ at $z < 100$ of at most 0.3 (for the case of axino-type Light Dark Matter with a mass of 10$\text{MeV}$ and without considering PopIII stars). We note however that among the models that they considered, only the neutralino DM model is comparable to our results for the no clumpiness case. Both results are in good agreement, however we do not use exactly the same mass as they do, and we do not consider cross-section enhancement. Varying the mass of the WIMP and considering a constant enhancement in cross-section can have enormous effects on the properties of the IGM as the DM energy injection rate scales as $\propto \frac{\langle \sigma v \rangle}{m_\chi}$ (disregarding the variation in the presence of substructure when changing $m_\chi$ due to free-streaming effects).

About the detectability of the 21$\text{cm}$ $\Delta T_b$ signal, defined as the absolute value of the difference of the differential brightness temperatures in 21$\text{cm}$ with and without DM energy injection (see Section 6.5), we can conclude considering the results that we obtain that, according to the sensitivity limits mentioned by Valdés et al. (2007), and considering a maximal absorption fraction $f_{abs} = 1$, a thermal relic WIMP with mass of 1$\text{TeV}$ is not likely to be detected from the global signal alone (see Section 6.5). A lower mass WIMP may be detectable, although a mass on the order of $\sim 10\, \text{GeV}$ or lower is not favoured by neither accelerator constraints (see Section 2.3), nor by direct and indirect detection searches (see Section 4.6). Comparing with the results obtained by Cumberbatch, Lattanzi & Silk (2010), we notice that for their different SUSY neutralino models, the authors obtain $\Delta T_b$ curves that cover a similar dynamic range than our curves. However, in the redshift range of $20 < z < 46$ that is interesting for detection, they tend to predict a stronger signal than our 1$\text{TeV}$ and 10$\text{GeV}$ WIMP models, because their curves are centred around $z \sim 50 - 60$, whereas ours are centred around $z \sim 100$. This also applies to our $f_{abs} = 0.01$ case, to which we should compare their results to (because the authors already incorporate an explicit $f_{abs}$ treatment and they have values $f_{abs} \sim 0.01$ in that redshift range; see their figure 12 and figure 13). Note that the authors use a lower mass 50$\text{GeV}$ WIMP, but they obtain the energy and number of photons and electron-positron pairs that result from the cascades produced by the primary products of the annihilation and use these in order to compute $f_{abs}$. According to their results, all of their SUSY models would be in principle detectable, except for their model 4 (see their Table 1 for more details).

Finally, as future work, we should mention that it is in principle possible to use the DM annihilation luminosity received by a halo (Section 5), together with the luminosity generated by the halo itself (Section 4.2), in order to compute the effects that the extra energy injection will have on the standard physical processes that take place in the baryonic component of the halo (Intra Cluster Medium or ICM in the case of galaxy cluster-sized haloes), like for example its neutral hydrogen abundance, Bremsstrahlung luminosity, etc... Also, being able to compute the ionization fractions and temperatures of individual haloes at high redshifts, as well as their spatial clustering, may provide a way of computing the angular fluctuations of the differential brightness temperature ($\Delta T_b$) signal for different WIMP scenarios, in order to have another test besides the predicted global signal (for example, Furlanetto, Oh & Pierpaoli (2006) study this, but without considering substructure or adiabatic compression).

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