Non-universal BBN bounds on electromagnetically decaying particles

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In Ref. \cite{1} we have recently argued that when the energy of a photon injected in the primordial plasma falls below the pair-production threshold, the universality of the non-thermal photon spectrum from the standard theory of electromagnetic cascades onto a photon background breaks down. We showed that this could reopen or widen the parameter space for an exotic solution to the “lithium problem”. Here we discuss another application, namely the impact that this has on non-thermal big bang nucleosynthesis constraints from \textsuperscript{4}He, \textsuperscript{3}He and \textsuperscript{2}H, using the parametric example of monochromatic photon injection of different energies. Typically, we find tighter bounds than those existing in the literature, up to more than one order of magnitude. As a consequence of the non-universality of the spectrum, the energy-dependence of the photodissociation cross-sections is important. We also compare the constraints obtained with current level and future reach of cosmic microwave background spectral distortion bounds.

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\section{I. INTRODUCTION}

Big Bang Nucleosynthesis (BBN) has been used since decades as a very powerful tool to constraint exotic particle physics (for reviews see for instance \cite{2,3}). In particular, meta-stable particles populating the plasma in the early universe could induce a non-thermal BBN phase via their decay products. Both hadronic and electromagnetic cascades typically contribute to these phenomena, although the former are much more model-dependent. On the contrary, electromagnetic cascades are known to lead to a quasi-universal spectrum, only dependent on the overall energy injected and the injection time: a monotonically decreasing, broken power-law (see e.g. Chapter VIII in \cite{4} for a basic derivation). Recently, we pointed out that in a particular regime, the commonly used universality of the spectrum achieved by photons as a result of electromagnetic cascades is violated \cite{1}. This corresponds to the situation when the photons injected at energy \(E_\gamma\) are not sufficiently energetic to induce pairs onto the background photons at temperature \(T\), and can be translated in the condition \(E_\gamma \lesssim 10 T_{\text{keV}}^{-1}\) MeV (we use natural units with \(c = k_B = 1\)). For \(T\) of order \(\mathcal{O}(\text{keV})\), characteristic of the period between the end of BBN to the formation of the cosmic microwave background (CMB), these energies are typically higher than the photo-disintegration thresholds of light nuclei, denoted by \(E_{\text{th}}\). As a result, the injection of photons whose energy falls in the couple of decades from few MeV to few hundreds MeV may have an impact different than the one estimated with the universal spectra, given by

\[
\frac{dN_\gamma}{dE_\gamma} = \begin{cases} 
K_0 \left( \frac{E_\gamma}{\epsilon_c} \right)^{-3/2} & \text{for } E_\gamma < \epsilon_X, \\
K_0 \left( \frac{E_\gamma}{\epsilon_X} \right)^{-2} & \text{for } \epsilon_X \leq E_\gamma \leq \epsilon_c, \\
0 & \text{for } E > \epsilon_c. 
\end{cases}
\]  

In the above expression, \(K_0 = \frac{E_0 \epsilon_X^2}{2 + \ln(\epsilon_c/\epsilon_X)}\) is a normalization constant that enforces the condition that the total energy is equal to the injected electromagnetic energy, \(E_0\); the characteristic energy \(\epsilon_c = m_\gamma^2/\epsilon_X^{\max}\) denotes the effective threshold for pair-production, \(\epsilon_X^{\max}\) being the highest energy of the photon background onto which pairs can be effectively created; \(\epsilon_X \lesssim \epsilon_c/3\) is the maximum energy of up-scattered inverse Compton (IC) photons (See \cite{5,6} for monte carlo studies leading to further justification of these parameters.)

In our previous publication \cite{1}, we have illustrated how this may reopen the possibility of purely electromagnetic new physics solution to the so-called lithium problem, but we anticipated that other domains may be affected. Here we outline the impact on the constraints in the abundances vs. lifetime plane for unstable early universe relics, decaying electromagnetically, and derived from the deuterium, helium-4 and helium-3 measurements. Our main conclusion is that the bounds are non-universal and that they may be significantly more stringent than commonly thought. In the following, we will compare the constraints obtained from different elements in the hypothesis of universal spectrum with the actual constraints obtained for monochromatic photon injections at different energies, below the pair-production threshold \(\epsilon_c\). This parameterization is used solely for the sake of clarity: the differences would persist for any spectrum (either primary photons or secondary due e.g. to up-scattering of background photons via IC by energetic \(\epsilon^\pm\) ) injected below the critical energy.

This article is structured as follows: In Sec. \ref{sec:2} we describe the features of the e.m cascades and the breakdown of the universal non-thermal spectrum, as well as our method to solve the relevant Boltzmann equations. In Sec. \ref{sec:3} we describe the non-thermal nucleosynthesis formalism and the observational constraints being used in the following. In Sec. \ref{sec:4} we review the constraints coming from CMB, notably from its spectral distortions, to which we will compare the BBN ones. Our results are reported in Sec. \ref{sec:5}. Finally, Sec. \ref{sec:6} contains a discussion with our conclusions.
II. E.M. CASCADES AND BREAKDOWN OF UNIVERSAL NON-THERMAL SPECTRUM

In general, in order to compute the non-thermal photon spectrum which can photo-disintegrate nuclei, one has to follow the coupled equations of both photon and electron-positron populations. For the problem at hand, however, where we limit ourselves to inject photons incapable of pair-production, it is a good first approximation to ignore the non-thermal electrons: while the injected photons will in general Compton scatter and produce them, a further process, typically IC onto the photon background, is needed to channel back part of their energy in the photon channel. The energy of these photons is significantly lower than the injected photon one: whenever they are re-injected below nuclear photo-dissociation thresholds they are actually lost for non-thermal nucleosynthesis, otherwise they would contribute to strengthening the bounds, although only by a few percents, for the cases discussed below. For simplicity, let us also start by assuming that all photon interactions are destructive, i.e. photons are not down-scattered to a lower energy. Within this approximation, the Boltzmann equation describing the non-thermal photon distribution function \( f_\gamma \) reads:

\[
\frac{\partial f_\gamma(E_\gamma)}{\partial t} = -\Gamma_\gamma(E_\gamma, T(t)) f_\gamma(E_\gamma, T(t)) + S(E_\gamma, t),
\]

where \( S(E_\gamma, t) \) is the source injection term, \( \Gamma_\gamma \) is the total interaction rate, and we neglected the Hubble expansion rate since interaction rates are much faster and rapidly drive \( f_\gamma \) to a quasi-static equilibrium, \( \frac{\partial f_\gamma(E_\gamma)}{\partial t} = 0 \). Thus, we simply have:

\[
f_\gamma^0(E_\gamma, t) = \frac{S(E_\gamma, t)}{\Gamma_\gamma(E_\gamma, t)},
\]

where the term \( S \) for an exponentially decaying species with lifetime \( \tau_X \) and density \( n_X(t) \), whose total e.m. energy injected per particle is \( E_0 \), can be written as

\[
S(E_\gamma, t) = \frac{n^0_X \zeta_X (1 + z(t)) e^{-t/\tau_X}}{E_0 \tau_X} p_\gamma(E_\gamma, t),
\]

with \( z(t) \) being the redshift at time \( t \), and the energy parameter \( \zeta_X \) (conventionally used in the literature) is simply defined in terms of the initial comoving density of the X particle \( n^0_X \) and the actual one of the CMB, \( n^\gamma_X \), via \( n^\gamma_X = n^0_X \zeta_X / E_0 \). We shall use as reference spectrum the one for a two body decay \( X \to \gamma U \) leading to a single monochromatic line of energy \( E_0 \), corresponding to \( p_\gamma(E_\gamma) = \delta(E_\gamma - E_0) \). If the unspecified particle \( U \) is (quasi)massless, like a neutrino, one has \( E_0 = m_X / 2 \), where \( m_X \) is the mass of the decaying particle. Note that here, we will be interested in masses \( m_X \) between a few MeV and \( \mathcal{O}(100) \) MeV, and at temperatures of order few keV or lower, hence the thermal broadening is negligible and a Dirac delta spectrum as the one above is appropriate.

We calculate \( \Gamma_\gamma \) by summing the rates of processes that degrade the injection spectrum, namely:

- Scattering off thermal background photons, \( \gamma X \):
  \( \gamma + \gamma \to \gamma + \gamma \).
  This has been studied in \[7\] and the scattering rate of a \( \gamma \)-ray with energy \( E_\gamma \) over a blackbody with temperature \( T \) is given by:
  \[
  \Gamma_{\gamma\gamma} = -0.1513 a^4 m_e \left( \frac{E_\gamma}{m_e} \right)^3 \left( \frac{T}{m_e} \right)^6.
  \]

- Bethe-Heitler pair creation : \( \gamma + N \to e^\pm \pm N \).
  The cross-section for this process is given by \[8\]:
  \[
  \sigma_{BH} \approx \frac{3 \alpha}{8 \pi} \ln \left( \frac{2 E_\gamma}{m_e} - \frac{218}{27} \right) Z^2.
  \]

- Compton scattering over thermal electron : \( \gamma + e^- \to \gamma + e^\pm \). In the Klein-Nishina limit of interest here, \( E_\gamma \gg m_e \), we have \[3\]:
  \[
  \Gamma_{CS} = \bar{n}_e \frac{3}{8 E_\gamma} \ln \left( \frac{2 E_\gamma}{m_e} + \frac{1}{2} \right),
  \]

where \( \bar{n}_e \) is the number density of background electrons and positrons.

In reality, not all scattered photons will be “lost”: even ignoring the energy transferred to \( e^- \) and \( e^+ \), Compton scattering and \( \gamma \gamma \) scattering still leave lower energy photons in the final state. This effect can be accounted for by replacing the RHS of Eq. \[2\] by the following term:

\[
S(E_\gamma, t) \to S(E_\gamma, t) + \int_{E_\gamma}^{\infty} dx K_\gamma(E_\gamma, x, t) f_\gamma(x, t).
\]

The additional term whose kernel is \( K \) accounts for scattered photons and is obtained by summing the differential rates for the \( \gamma \gamma \) scattering off background photons and the Compton scattering over thermal electrons, respectively given by \[7\]

\[
\frac{d\gamma \gamma (E_\gamma, E'_\gamma)}{dE'_\gamma} = \frac{1112}{10125 \pi} \alpha^2 r_e^2 \frac{m_e^{-6}}{63} \times \frac{E'^2}{E_\gamma^2} \left[ 1 - E_\gamma / E'_\gamma + \left( \frac{E_\gamma}{E'_\gamma} \right)^2 \right]^2,
\]

\[
\frac{d\gamma e^- (E_\gamma, E'_\gamma)}{dE'_\gamma} = \pi r_e^2 \bar{n}_e \frac{m_e}{E_\gamma^2} \times \left[ \frac{E'_\gamma}{E_\gamma} + \frac{m_e}{E'_\gamma} + \left( \frac{m_e}{E_\gamma} - 1 \right)^2 - 1 \right].
\]

The integral in Eq. \[8\] now depends on \( f_\gamma \). We numerically solve this Boltzmann equation using an iterative method: we start from the Dirac distribution and the algebraic solution of Eqs. \[3\], plug in the result thus
It is clear that unless the injected photon energies are too high, this only brings a modest correction to the low energy tail of the spectrum, with the expected improvement in the constraints being even less prominent. The iterative solution technique adopted above would still perform correctly, although a detailed evaluation would render the calculation unnecessarily lengthy, and will not be pursued further here. We checked in fact that for the cases discussed in the following, four iterations for the photon spectrum are enough to obtain bounds accurate at the 10% level (and often better) and always on the conservative side.

Since the critical energy for pair-production is a dynamical quantity that increases at later times due to the cooling of the universe, it may happen that the primary photons energy \( E_0 \) is above threshold for pair-production at early times, and below it at late times (we do take into account that the decay is not instantaneous). In general, at each time we will compare \( E_0 \) with \( \epsilon_c \), and use the universal spectrum when \( E_0 > \epsilon_c \) or the monochromatic spectrum with the complete expression for \( S \) when \( E_0 < \epsilon_c \). This gives always a qualitatively correct solution, albeit it is somewhat approximate when \( E_0 \sim \epsilon_c \)

Since this is realized only in a very narrow interval of time, however, the final result are also quantitatively robust, barring peculiar mass-lifetime fine-tuning.

## III. NON-HEATLUM NUCLEOSYNTHESIS

### A. Review of the formalism

At temperatures of few keV or lower, the standard BBN is over, and the additional nucleosynthesis can be simply dealt with as a post-processing of the abundances computed in the standard scenario, for which we use the input values from Parthenope [9], with the updated value of \( \Omega_b \) coming from [10].

The non-thermal nucleosynthesis due to electromagnetic cascades can be described by a system of coupled differential equations of the type

\[
\frac{dY_A}{dt} = \sum_T Y_T \int_0^\infty dE_\gamma f_\gamma(E_\gamma,t)\sigma_{\gamma+T\rightarrow A}(E_\gamma) - Y_A \sum_P \int_0^\infty dE_\gamma f_\gamma(E_\gamma,t)\sigma_{\gamma+A\rightarrow P}(E_\gamma) \tag{11}
\]

where: \( Y_A \equiv n_A/n_0 \) is the ratio of the number density of the nucleus \( A \) to the total baryon number density \( n_0 \) (this factors out the trivial evolution due to the expansion of the universe); \( \sigma_{\gamma+T\rightarrow A} \) is the photodissociation cross section of the nucleus \( T \) into the nucleus \( A \), i.e. the production channel for \( A \); \( \sigma_{\gamma+A\rightarrow P} \) is the analogous destruction channel. Both cross sections are actually vanishing below the corresponding thresholds. In general one also needs to follow secondary reactions of the nuclear byproducts of the photodissociation, which can spallate on or fuse with

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**FIG. 1:** Spectrum computed in this article (solid black line) compared with approximated one used in the literature (long-dashed red), for the case \( E_0 = 70 \text{ MeV} \) at \( T = 100 \text{ eV} \). The short-dashed, blue lines show the contribution from the photon population as computed in our iterative treatment, with the number of iterations increasing from one to seven from bottom to top. The dot-dashed curve is the estimated contribution to the photon spectrum from the non-thermal population of electrons excited by the energy loss channels of the photons.
background thermalized target nuclei but none of that is relevant for the problem at hand: According to [11], the only significantation secondary production is that of $^6$Li. Despite extensive work in the past, the current observational status of $^6$Li as a reliable nuclide for cosmological constraints is doubtful, given that most claimed detections have not been robustly confirmed, and a handful of cases are insufficient to start talking of a “cosmological” detection, see [12]. We shall thus conservatively ignore this nuclide and the secondary reactions in the following.

With standard manipulations, namely by transforming Eq. (11) into redshift space, defining $H(z) = H_0^0(1 + z)^2$ as appropriate for a Universe dominated by radiation (with $H_0^0 = H_0\sqrt{\Omega_\gamma}$, $H_0$ and $\Omega_\gamma$ being the present Hubble expansion rate and fractional radiation energy density, respectively) one arrives at

$$\frac{dY_A}{dz} = \frac{-1}{H_0^0(z + 1)^3} \times \left[ \sum_f Y_f \int_0^\infty dE_\gamma f_\gamma(E_\gamma, z)\sigma_{\gamma} + T_{\rightarrow A}(E_\gamma) \right]$$

$$- Y_A \sum_p \int_0^\infty dE_\gamma f_\gamma(E_\gamma, z)\sigma_{\gamma} + p_{\rightarrow A}(E_\gamma) \bigg(12\bigg)$$

which is solved numerically for the cases of interest.

**B. Light element abundances**

Unless the $^2$H photodissociation channel is closed, which happens in the energy window between about 1.6 and 2.2 MeV, $^7$Li-based constraints are typically irrelevant. Actually, the interpretation of $^7$Li observations in terms of primordial yield is still a subject of debate, see [2] [3] [12], hence it would not be robust to derive constraints based on its observations. On the other hand, cascades could reconcile the envelope of $^7$Li observed values with a primordial origin, albeit a non-standard one, as shown in [1], to which we address for further details.

The most important nuclei for what follows are thus: $^4$He, which can only be destroyed by non-thermal BBN, hence we just impose the 2-$\sigma$ lower limit $^4$He $> 0.2368$ from [13]; $^2$H, for which we adopt the 2-$\sigma$ limit $2.56 \times 10^{-5}$ $< ^2$H $< 3.48 \times 10^{-5}$ from [14]; similar results would follow by adopting the combination value compiled in [2], namely $2.45 \times 10^{-5} < ^2$H $< 3.31 \times 10^{-5}$, which is also closer to the results of [15]; $^3$He, for which we impose no observational lower limit, but the 2-$\sigma$ upper limit from [16] $^3$He $< 1.5 \times 10^{-5}$. For the current application, the network of reactions used is reported in Fig. 2 and follows the parameterization in the appendix of [11]. (Actually, the reaction $^4$He$(\gamma, ^2$H)$^2$H is significantly suppressed with respect to the others and thus not shown in the figure, but is included in our numerical treatment.) Note that all cross-section share the same qualitative features: they rise fast just above threshold, go through a peak (the so-called Giant Dipole Resonance), eventually showing a decreasing tail at higher-energies. We shall compare the bounds thus obtained with the ones coming from CMB spectral distortions and entropy production, briefly recalled in the following section.

**IV. CONSTRAINTS FROM THE CMB**

It is well known that a late injection of photons in the thermal bath can lead to additional measurable cosmological alterations.

For instance, the injection of a significant amount of energy can lead to modification of the photon-baryon ratio $\eta$ or equivalently, to the increase of the co-moving entropy. Since the inferred values of $\Omega_b$ at BBN and CMB epoch are compatible, no major entropy release could have taken place between nucleosynthesis and decoupling. It can be shown that, in a Universe dominated by radiation and by considering that the decays has happened at $t \sim \tau$, we have for a small fractional change in entropy (see e.g. [17]):

$$\frac{\Delta S}{S} \simeq \ln \frac{S_f}{S_i} = 2.14 \times 10^{-4} \left( \frac{\zeta_{X \rightarrow \gamma}}{10^{-9} \text{GeV}} \right) \left( \frac{\tau_X}{10^8 \text{ s}} \right)^{1/2}$$

(13)

with a slight abuse of notation since $\zeta_{X \rightarrow \gamma}$ now has to be intended to include any electromagnetically interacting decay products, all of which contribute to modify the photon-baryon ratio. To derive a statistically sound constraint, one should combine BBN and CMB data, allowing for an entropy increase between the two epochs. Since, as we shall see, this constraint is typically much weaker than others, such an exercise would bring us far beyond the scope of this paper adding a lengthy and unnecessary complication. We shall thus limit ourselves to illustrate the constraint that would follow by allowing for a maximal 2% increase in the entropy between the two
periods. This is an educated guess of the order of the bounds that one can expect, roughly corresponding to the 2-σ error bars on Ω_b from Planck 2015 [10].

Furthermore, as reviewed in detail in [18], the spectrum of the CMB itself can also be affected through two types of deformation: a modification of the chemical potential µ and a modification of the Compton-y parameter, which is related to the energy gained by a photon after a Compton scattering. To first order, it is possible to distinguish the era of µ-distortion from the era of Compton-y distortion, because the rate of the processes which are responsible for one type of distortions dominates at very different time. Basically, the µ-distortion arises from rare process, implying a change in the number of photons such as γ + e → γ + γ + e, whereas the Compton-y distortions are due to the end of the equilibrium of Compton reactions, which happens much later in the history of the Universe, with a schematic µ - y transition happening at for z ≈ 4 × 10^5, i.e. τ ≈ 5 × 10^{10} s. For the relatively early time we focus on, the constraints come essentially from µ-type distortions. We follow here the results of Ref. [18], which contains improvements with respect to the ones given in [19], notably for z ≲ 2 × 10^6, while [19] is accurate enough at late times (see Fig. 16 in [18]). Hence we adopt

$$\mu \simeq 8.01 \times 10^{2} \left(\frac{\tau_X}{1 \text{ s}}\right)^{1/2} \times J \times \left(\frac{\xi_{X\rightarrow Y}}{1 \text{ GeV}}\right)$$

with

$$J = \begin{cases} \exp \left[-\left(\frac{\pi c}{\tau_X}\right)^2\right] & \text{for } z < 2 \times 10^6 \\ 2.082 \left(\frac{\pi c}{\tau_X}\right)^{10} \exp \left[-1.988 \left(\frac{\pi c}{\tau_X}\right)^{10}\right] & \text{otherwise} \end{cases}$$

where $\tau_{\text{IC}} = 1.46 \times 10^9 (T_0/2.7 \text{ K})^{-12/5} (\Omega_0 h^2)^{4/5} (1 - Y_p/2)^{4/5}$ is the “double Compton” interaction time in terms of the current CMB temperature $T_0$, with $Y_p \simeq 0.25$ the primordial mass fraction of Helium-4. We use the limit given by COBE on the chemical potential: $|\mu| \leq 9 \times 10^{-5}$ [20], but we will also show the sensitivity that should characterize the future experiment PIXIE, of the order of $\mu \gtrsim 5 \times 10^{-8}$, at 1-σ [21].

V. RESULTS

One of the most peculiar features of the spectral non-universality of photons injected below the pair production threshold is that the final outcome reflects the energy distribution of the injected photons with respect to the shape of the relevant photo-disintegration cross sections, shown in Fig. 2. This motivated us to choose in the following, for each nuclide, the results for two representative examples of monochromatic injection: one close to the resonant peak, and another one well after it. The markedly different outcomes obtained in the two cases should thus convincingly argue that constraints of actual models are going to be determined not only by the decay time and the overall energy injected, but also by the energy range at which the bulk of the photons lies.

A. Constraints from $^4$He

The simplest situation is certainly the one concerning $^4$He: being the only abundant nucleus subject to photodisintegration, its non-thermal e.m. production is irrelevant and one only has to care about its destruction, i.e. only the term proportional to $Y_A$ at the RHS of Eq. (12) is important. The results obtained by using a monochromatic injection at 70 MeV (hatched/light shaded red), at 30 MeV (dark shaded red), and the universal spectrum are shown in Fig. 3. The vertical lines indicate the time at which the threshold energy for pair production $\epsilon_c$ starts exceeding the corresponding injected energy. One might naively expect that this is the time at which the constraints obtained from the incorrect use of the universal spectrum start to deviate from the actual ones. However, when taking into account the fact that the decay is not instantaneous, it turns out that constraints already start to deviate at $\tau \sim \tau_X/5$, and the closer to the post-threshold cross-section resonance we inject energy, the earlier deviations appear.

![FIG. 3: Constraints from Helium-4 depletion in the standard case (black line) and for a non-universal spectrum with $E_0 = 30$ MeV (dark shaded red) and $E_0 = 70$ MeV (hatched/light shaded red). We also show the sensitivity to entropy variation constraint (green dashed line), current constraints from CMB spectral distortions (excluded above the short-dashed blue line) and the sensitivity reach of the future mission PIXIE [21] (above the red dot-dashed line).](image-url)
both cases bounds can differ from the ones derived with the universal spectrum by a large factor, up to an order of magnitude if the energy of the photons is around the peak of the photodissociation cross-section. For higher-injected energies, they tend to become closer to the universal spectrum constraints, as it should. In fact one can envisage fine-tuned situations in which they become slightly weaker, albeit this conclusion does depend on the extrapolation of the photodisintegration cross sections, whose reliability at high-energy has never been quantitatively assessed in the context of BBN applications.

B. Constraints from $^2$H

For deuterium, the situation is more complicated because several regimes are present. At low $\tau_X$, $\epsilon_c$ is below $^4$He photodissociation threshold (and in some cases also below $A = 3$ nuclei photodissociation thresholds, which are however less relevant). Hence only constraints from over-destructions are present. At high-$\tau_X$, however, what dominates is the over-production from $^4$He destruction.

Fig. 4 shows the illustrative case where production channels are turned off: this is exact for $E_0 \lesssim 8$ MeV, but a good approximation till $E_0 \lesssim 20$ MeV. Note the qualitative similarity to the $^4$He case, apart for the modifications due to the different features of the respective cross-sections.

Whenever production channels from $^4$He are open, which requires $E_0 \gtrsim 20$ MeV, the constraints are significantly stronger at large $\tau_X$, as shown in Fig. 5. Once again, a violation of universality (and a sensitivity to the energy dependence of the cross section) is clearly manifest by the two cases shown, $E_0 = 30$ MeV and $E_0 = 70$ MeV.

It is also worth noting that for deuterium the constraints are significantly better than the ones coming from CMB spectral distortions. By improving the sensitivity to $\mu$ down to $\mu \gtrsim 5 \times 10^{-8}$, the sensitivity expected by the future mission PIXIE [21] (shown by the red, dot-dashed curve) would greatly strengthen these constraints, with the exception of very small lifetimes where deuterium over-destruction would still provide the dominant bounds.

C. Constraints from $^3$He

First of all, a premise is necessary: There are in fact two nuclei with $A = 3$, $^3$He and $^3$H, the latter being unstable to beta decay into $^3$He with a half-time of over 12 years, or about $4 \times 10^8$ s. Practically, however, for the purposes of the constraints discussed here one can sum the equations for $^3$He and $^3$H and treat them as a single effective nucleus with $A = 3$. The reason is twofold: first, we only require $^3$He not to be overproduced with respect to the observational upper limit. Hence the key reactions are the production channels by single nucleon photodisintegration from $^4$He, which are only open above 20 MeV, rather than the destruction ones. Second, $^3$He and $^3$H are “mirror nuclei” under the isospin symmetry $n \leftrightarrow p$, and their nuclear properties are in fact very similar: the corresponding thresholds in nuclear cross sections, for instance, only differ by some 0.8 MeV (compare the two curves in Fig. 2). From Fig. 6, where we report our results, it is clear that the photodisintegration cross section for single nucleon emission from $^3$He, when open, is so important that very stringent nucleosynthesis con-
VI. CONCLUSIONS

We have argued that the universality of the photon spectral shape in electromagnetic cascades has often been used in cosmology even beyond its regime of applicability. When the energy of the injected photons falls below the pair-production threshold, i.e. approximately when $E_\gamma \lesssim m_e^2/(22 T) \sim 10 T^{-1} \text{keV}$, the universal form breaks down. In [1], we showed how this could potentially open the possibility of a purely electromagnetic decay solution to the so-called “lithium problem”. In this article, we showed how important the modifications to the photon spectrum in this regime are for the constraints from non-thermal BBN. This required the numerical solution of the relevant Boltzmann equations, which we attacked by an iterative scheme.

The constraints we obtained, for illustrative cases of monochromatic energy injection at different epochs, are often much stronger than the ones presented in the literature (up to an order of magnitude), notably when the injected photon energy falls close to the peak of the photodisintegration cross-section of the relevant nucleus. In fact, the breaking of the non-universality is non-trivial and is essentially controlled by the energy behavior of the cross-sections: in the universal limit, most of the photons lie at relatively low-energies, so that the cross-section behaviour at the resonance just above threshold is what matters the most. In the actual treatment, the photons may be also sensitive to the high-energy tail of the process. Future studies aiming at assessing the nuclear physics uncertainties affecting these types of bounds would benefit from this insight. It cannot be excluded that in some cases constraints weaken a bit with respect to what considered in the literature.

We also compared BBN bounds with constraints coming from CMB spectral distortions. A summary plot of the “best constraints” is reported in Fig. 7 for two choices of the monochromatic photon energy. We concluded that BBN limits are improving over current constraints from COBE via the requirement not to under-produce $^2\text{H}$ (at low injection lifetime $\tau_X$), or not to over-produce $^3\text{He}$ (at high $\tau_X$), while $^4\text{He}$ is never competitive. The bounds from a future CMB spectral probe as PIXIE would not only greatly improve current CMB constraints, but would also reach the level of current constraints from $^3\text{He}$ (often improving over them) allowing for an independent consistency check. This is reassuring, since the cosmological reliability of $^3\text{He}$ constraints does stand on some astrophysical assumptions. Below $\tau_X \sim 5 \times 10^5 \text{s}$, however, $^2\text{H}$ constraints would probably remain the most stringent ones for a long time to come. Fortunately they are: i) quite robust, relying on the single, well-known cross-section $^2\text{H}(\gamma,n)p$; ii) easy to compute, since no coupled network of equations needs to be solved, the problem reducing to the numerical evaluation of a single integral (same situation leading to Eq. (6) in [1].) This is also the region where constraints coming from hadronic decay modes (not revisited here) are quite

FIG. 6: Constraints from the Helium-3 production in the standard case (black line) and for a non-universal spectrum with $E_0 = 30 \text{ MeV}$ (dark shaded red) and $E_0 = 70 \text{ MeV}$ (hatched/light shaded red). All other constraints/sensitivities shown as in Fig. 3.

FIG. 7: Global BBN best constraints in the standard case (black line) and for a non-universal spectrum with $E_0 = 30 \text{ MeV}$ (dark shaded red) and $E_0 = 70 \text{ MeV}$ (hatched/light shaded red). All other constraints/sensitivities shown as in Fig. 3.
stringent. A synergy between BBN and CMB is thus going to be necessary for this kind of physics even in the decades to come.

In conclusion, our work suggests that models in the literature that fulfilled the BBN constraints with less than an order of magnitude margin should perhaps be reconsidered. In particular, those characterized by soft gamma-ray emissions and/or at relatively late times should have been more prone to incorrect conclusions about their viability. Our study also suggests that actual bounds should be derived via a case-by-case analysis. Finally, we provided further arguments supporting the usefulness of an improved constraint from CMB spectral distortion of the \( \mu \) type, since it would not manifest the unexpected “spectral sensitivity” that we have uncovered.

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