Hadronic Decay of Late-Decaying Particles and Big-Bang Nucleosynthesis

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We study the big-bang nucleosynthesis (BBN) scenario with late-decaying exotic particles with lifetime longer than \(\sim 1\) sec. With a late-decaying particle in the early universe, predictions of the standard BBN scenario can be significantly altered. Therefore, we derive constraints on its primordial abundance. We pay particular attention to hadronic decay modes of such particles. We see that the non-thermal production process of D, \(^3\)He and \(^\alpha\)Li provides a stringent upper bound on the primordial abundance of late-decaying particles with hadronic branching ratio.

It has long been recognized that the (standard) big-bang nucleosynthesis (BBN) provide a good probe for the early universe. With our current knowledge of nuclear reaction processes, we can precisely calculate abundances of the light elements (in particular, D, \(^3\)He, \(^4\)He, \(^6\)Li, and \(^7\)Li) as functions of the baryon-to-photon ratio \(\eta = n_B/n_\gamma\). Thus, comparing the theoretical predictions with the observations, we can obtain various informations about the evolution of the universe. Importantly, using the value of \(\eta\) suggested by the WMAP (\(\eta = (6.1 \pm 0.3) \times 10^{-10}\) [1]), the theoretical predictions show relatively good agreement with the observations.

In cosmological scenarios in the frameworks of physics beyond the standard model, however, the BBN may not proceed in the standard way. This is because, if we assume physics beyond the standard model, there exist various exotic particles. Those exotic particles may cause non-standard processes and spoil the success of the standard BBN.

In particular, if the exotic particle (called \(X\) hereafter) decays radiatively and/or hadronically after the BBN starts, the primordial abundances of the light elements may be significantly affected. Indeed, energetic particles produced by the decay of \(X\) may scatter off and dissociate the background nuclei. If such processes occur with sizable rates, predictions of the standard BBN scenario are changed.

In various models of particle physics, there exist long-lived (but unstable) particles and hence their effects on the BBN should be studied. In particular, many of those particles interact very weakly and hence it is difficult to study their properties by collider experiments. Thus, in some case, the BBN provides useful and important informations about those weakly interacting particles. Probably the most famous example of such late-decaying particle is the gravitino in the supergravity theory. Gravitinos are produced in the very early universe by scattering processes of the particles in the thermal bath. Their interaction is suppressed by the inverse powers of the (reduced) Planck scale \(M_P\), and hence the lifetime becomes very long. An order of magnitude estimate shows that, if the gravitino is lighter than \(\sim 10\) TeV, its lifetime becomes longer than \(\sim 1\) sec. In this case, the thermally produced gravitinos decay after the BBN starts.

Effects of the radiative decay of such long-lived particles have been extensively studied (see, e.g., [2–5]). However, in many cases, hadronic branching ratio may not be negligible and hence, in studying the BBN with late-decaying particles, it is necessary to consider effects of the hadronic decay processes. Even if \(X\) dominantly decays into the photon (and something else), hadronic branching ratio is expected to be at least \(10^{−(2−3)}\) since the emitted photon can be converted to a quark-antiquark pair. Of course, if \(X\) directly interacts with the colored particles, hadronic branching ratio may become larger.

In the past, the BBN with hadro-dissociation processes induced by hadronic decays of long-lived particles was studied in [6], which are effective for relatively long lifetime (\(\gtrsim 10^2\) sec).\(^1\) The analysis in Ref. [6] contains, however, a lot of room to be improved since many of nuclear reactions that they used were not accurate enough or not available at that time. After the study of [6], however, there have been significant theoretical, experimental and observational progresses in the study of the BBN. First of all, new data for the hadron reactions have become available and their qualities have been improved very much. Moreover, the primordial abundances of the light elements have been precisely determined with various new observations. In addition, it has been recently known that some of the non-standard processes induced by the decay of late-decaying particle \(X\), which were not taken into account in [6], may play important roles in the BBN. With these progresses, a new study of the late-decaying particles with hadronic branching ratio should be relevant.

Thus, in this letter, we reconsider the BBN processes with long-lived exotic particle \(X\) paying particular attention to the effects of the hadronic decay modes. As a result, we will see that, with hadronic decay modes, the constraint on the primordial abundance of \(X\) becomes

\(^1\)BBN constraints from the interconversion process between neutrons and protons by hadronic decays were studied in Refs. [7,8], which is effective for shorter lifetime (\(\lesssim 10^2\) sec).
very severe compared to the case only with the radiative decay modes. In particular, we will see that non-thermal production of D, 3He and 6Li provides a stringent constraint.

We first introduce the framework of our study. Although we have several candidates of the late-decaying particles, we perform our analysis as model-independently as possible. Thus we parameterize the property of the late-decaying particle X using the following parameters: \( E_X \) (released energy from the single decay of X) which is equal to its mass \( m_X \) unless specially stated, \( E_{\text{jet}} \) (energy of the primary parton from the decay of X), \( \tau_X \) (lifetime), \( B_h \) (hadronic branching ratio), and the primordial abundance of X. We parameterize the primordial abundance by using the following “yield variable” \( Y_X = n_X/s \), which is defined at the cosmic time \( t \ll \tau_X \). Here, \( n_X \) is the number density of X while s is the total entropy density. We assume that X decays only into the particle in the “observable sector,” and that the branching ratio decaying into the hidden-sector particle vanishes.

In our analysis, we calculate the primordial abundances of the light elements for given sets of the parameters listed above taking account of dissociation processes due to hadronic (and electromagnetic) interactions. The outline of our calculation is as follows. In order to study the effects of the hadronic decay modes, we first calculate the energy spectrum of the (primary) hadrons (in particular, protons and neutrons) generated from the partons directly emitted from X. Then, we follow evolutions of the hadronic showers. As a result, the numbers of the light elements produced (or destroyed) by the decay of X is calculated. We include the hadro-dissociation processes (as well as the photo-dissociation ones) into the BBN reaction network and numerically follow the evolution of the abundances of the light elements. Details of our study will be described elsewhere [9].

When X decays hadronically, quarks or gluons (i.e., partons) are first emitted. Those partons are hadronized soon after the decay. Thus, we should consider how the nucleons (and the mesons) propagate in the thermal bath.

In considering the propagation of stable particles (i.e., proton, neutron, and heavier nuclei), there are two classes of important processes. The first is the scattering process with the background photons \( \gamma_{\text{BG}} \) and electrons \( e_{\text{BG}} \). By scattering off \( \gamma_{\text{BG}} \) and \( e_{\text{BG}} \), energetic nuclei lose their energy without affecting the abundances of light elements. The second class is the scatterings with the background nuclei. With such processes, first of all, the background nuclei become energetic after the scattering. Therefore the energetic nuclei are copiously produced. (We call this “hadronic shower.”) In addition, if inelastic scattering occurs, background nuclei are dissociated and the abundances of the light elements are changed. Thus, if the second class of reactions occur significantly, the abundances of the light elements deviate from the predictions of the standard BBN.

Including relevant hadronic scattering processes (as well as photo-dissociation processes), we have calculated the abundances of the light elements. In our study of the evolution of the hadronic showers, the basic framework is the same as that used in [6] although there are several modifications. The most important improvements are as follows. (i) We carefully take into account the energy loss processes for high-energy nuclei through the scattering with background photons or electrons. In particular, dependence on the cosmic temperature, the initial energies of nuclei, and the background \(^4\)He abundance are considered. (ii) We adopt all the available data of cross sections and transferred energies of elastic and inelastic hadron-hadron scattering processes. (iii) The time evolution of the energy distribution functions of high-energy nuclei are computed with proper energy resolution. (iv) The JETSET 7.4 Monte Carlo event generator [10] is used to obtain the initial spectrum of hadrons produced by the decay of X. (v) The most resent data of observational light element abundances are adopted. (vi) We estimate uncertainties with Monte Carlo simulation which includes the experimental errors of the cross sections and transferred energies, and uncertainty of the baryon to photon ratio \( \eta \). (We take \( \eta = (6.1 \pm 0.3) \times 10^{-10} \).

We are interested in the situation where the number density of X is small enough so that the energetic particles in the hadronic shower dominantly scatter off the background particles. (Otherwise, the abundances of the light elements are so affected that the results become inconsistent with the observations.) If the energetic nuclei lose most of their energy by scattering off \( \gamma_{\text{BG}} \) and \( e_{\text{BG}} \), the number densities of the (background) nuclei is almost unaffected. On the contrary, if the scattering rate with \( \gamma_{\text{BG}} \) and \( e_{\text{BG}} \) becomes negligible, number densities of the light elements are significantly changed by the hadronic processes.

For charged nuclei, energy-loss rates due to the scatterings with the \( \gamma_{\text{BG}} \) and \( e_{\text{BG}} \) depend on the velocity of the nuclei \([7]\). For the temperature \( T \lesssim 20 \text{ keV} \), relativistic nuclei do not lose their energy by the scatterings with \( \gamma_{\text{BG}} \) and \( e_{\text{BG}} \). For non-relativistic nuclei with velocity \( \beta_N \), energy-loss is dominated by the scattering process with background electrons with velocity \( \beta_e < \beta_N \). (As pointed out in [7], energy loss via the scatterings with high-velocity electron is extremely suppressed and is negligible in our case.) Thus, we use the energy-loss rate

\[
\frac{dE_N}{dt} = -\frac{4\pi\alpha^2 Z^2 n_e}{m_e\beta_N} I(\beta_N/\sqrt{2T/m_e}) \ln \left( \frac{\Lambda m_e \beta^2}{\omega_p} \right),
\]

where \( \omega_p \) is the plasma frequency, \( m_e \) is the electron mass, \( Z \) is the charge of the nuclei, \( \alpha \) is the fine structure constant, \( \Lambda \) is a constant of \( O(1) \). (In our numerical calculations, we take \( \Lambda = 1 \).) In addition,

\[
I(r) = \frac{4}{\sqrt{\pi}} \int_0^r dx x^2 e^{-x^2}.
\]
Notice that the number density of the background electron with $\beta_e < \beta_N$ is given by $n_e I (\beta_N / \sqrt{2T/m_e})$. For $\beta_N \gg \sqrt{T/m_e}$, $I \approx 1$ and the energy-loss is very efficient. On the contrary, once the velocity of the nuclei becomes smaller than the thermal velocity of the electron $\langle \beta_e \rangle \sim \sqrt{T/m_e}$, $I \sim \mathcal{O}(\beta_N / \sqrt{2T/m_e})^3$ and energy loss becomes less effective.\(^2\)

For the neutral particle (i.e., neutron), on the contrary, the scattering process with $\epsilon_{BG}$ is not important for $T \lesssim 0.1$ MeV. Thus, once energetic neutrons are produced, they may induce hadronic showers by scattering off the background nuclei and change their abundances. If the temperature becomes lower than $\sim 0.3$ keV, however, most of the (energetic) neutrons decay into protons before scattering off the background particles. Thus, if $\tau_X \gtrsim 3 \times 10^7$ sec, the most important constraint on $Y_X$ is from the photo-dissociation processes of the light elements.

In our analysis, we also take account of the photo-dissociation processes induced by high-energy photons generated from the decay products of $X$. Once $X$ decays, energy as large as $(1 - B_h)Y_X$ is directly deposited into the radiation while the rest of the released energy $B_{hX}$ first goes into the hadronic sector. Importantly, however, spectrum of the high energy photon primarily depends on the total amount of the injected energy (as well as the temperature). Thus, we approximate that all the emitted energy in the decay process is eventually converted to the form of radiation, calculate the photon spectrum, and estimate the photo-dissociation rates of the light elements. For details of the treatment of the radiatively decaying particles, see [3].

Moreover, pions and kaons produced in the hadronization are also considered. For $T \gtrsim 0.1$ MeV ($t \lesssim 10^2$ sec) they can scatter off the background nuclei before their decay, which leads to increase the neutron to proton ratio and hence the abundance of D and $^4$He. In this letter, we have conservatively omitted the similar effects induced by $n\bar{n}$ and $p\bar{p}$ pairs which are effective for $t \gtrsim 10^2$ sec [8] because of insufficient data for the cross sections of energetic $\bar{n}$ and $\bar{p}$. If we include them, the constraint can become severer [8].

With the procedure explained above, we have calculated the abundances of the light elements for a given set of the model parameters. By comparing the results with the observations, we derived upper bounds on $Y_X$. As observational constraints on the light element abundances, we adopt the following values, D/H = $(2.8 \pm 0.4) \times 10^{-5}$ [11], $^4$He mass fraction $Y_p = 0.238 \pm 0.002 \pm 0.005$ by Fields and Olive (FO) [12] and $Y_p = 0.242 \pm 0.002(\pm 0.005)_{\text{syst}}$ by Izotov and Thuan (IT) [13], $\log_{10}(^7\text{Li}/H) = -9.66 \pm 0.056(\pm 0.3)_{\text{syst}}$ [14], $^6\text{Li}/^7\text{Li} < 0.07(2\sigma)$ [15], and $^3\text{He}/D < 1.13(2\sigma)$ [16]. The above errors are at 1 $\sigma$ level unless otherwise stated. Here we added the systematic error to $Y_p$ (IT) which is as same as the value in Fields and Olive. For $^7\text{Li}/H$ we also added the additional systematic error for the possibilities that $^7\text{Li}$ in halo might have been depleted in stars or supplemented by production in cosmic-ray interactions [17,18]. For $^7\text{Li}/H$ we conservatively use only the upper bound since the experimental data for the non-thermal $^7\text{Li}$ production by energetic $^4$He, especially for the transfered energy to $^4$He through the inelastic collisions between nucleon and $^4$He, are insufficient, by which we do not include them in our computation.

In Figs. 1 and 2, we show the upper bound on $m_X Y_X$ as a function of the lifetime $\tau_X$ for $B_h = 10^{-3}$ and 1. (Here, we used $m_X = 1$ TeV although, even if we vary $m_X$ within $0.1 - 100$ TeV, the bounds do not change significantly [9].) In deriving the bound we estimate the confidence levels by the $\chi^2$ fitting including both the theoretical and the observational errors. As we mentioned, for the case with very long lifetime, upper bound is almost the same as the one for the case with radiatively decaying particles [3]. The most stringent constraint comes from the overproduction of $^3$He.\(^3\)

For shorter lifetime $\tau_X \sim 10^8$ sec, overproductions of D and $^6$Li provide constraints on $m_X Y_X$. D is mainly produced by the hadro-dissociation of $^4$He. Notice that the hadro-dissociation of $^4$He is possible with energetic $p$ and $n$ even at such early epoch although the photo-dissociation of $^4$He can be effective only at $t \gtrsim 10^6$ sec ($T \lesssim 1$ keV) since high energy photons lose their energy by scatterings off the background photons before interacting with $^4$He at earlier epoch. The overproduction of $^6$Li is mostly due to the non-thermal processes with energetic T and $^4$He which are generated by the hadro-dissociation of $^4$He. These T and $^4$He can be sufficiently energetic and may scatter off the background $^4$He to produce $^6$Li via the processes $T + ^4$He $\rightarrow ^6$Li + n and $^3$He + $^4$He $\rightarrow ^6$Li + p.

One might think that the produced $^6$Li may be energetic and scatter off background nuclei to be destroyed. In our numerical calculations, we obtain the energy distribution of the non-thermally produced $^6$Li and calculated the surviving rate of such $^6$Li. We have found that, for $\tau_X \gtrsim 100$ sec, the surviving rate is almost 1. Therefore, one does not have to take account of the destruction of the produced $^6$Li. If the temperature is higher than

\(^2\)In the earlier version of this letter, we did not take into account the velocity distribution of $\epsilon_{BG}$ and used the formula for the case when all the background electrons have the same velocity $\langle \beta_e \rangle$. Thus, for $\beta_X = \langle \beta_e \rangle$, the energy-loss rate we used was incorrect and was underestimated. This resulted in an overestimate of the non-thermally produced $^6$Li.

\(^3\)However, even for $\tau_X \gtrsim 10^8$ sec, the hadronic processes give some contribution to the D production, which can be seen from comparison between Fig. 1 and Fig. 2.
95% C.L. for the case of \( B_{\beta} = 10^{-3} \) reporting the hadro-dissociation processes. In this case, background electron and lose their energy without com-
tron (and positron) is still abundant and, in this case, decays when the number density of the background elec-
cay of \( \sim 10 \) keV, however, non-thermally produced \( ^{6}\text{Li} \) is destroyed by the thermal process \( ^{6}\text{Li}(p, ^{3}\text{He})^{3}\text{He} \), as was also pointed out in [19]. Thus the constraint from the \( ^{6}\text{Li} \) is weakened for \( m_{\chi} X \).

FIG. 1. Upper bounds on \( m_{\chi} Y_{X} \) as a function of \( \tau_{X} \) at 95% C.L. for the case of \( B_{\beta} = 10^{-3} \). The name of the element which gives the constraint is written by each line. We assume that two hadron jets are produced by single de-
cay of \( \chi \) with the energy \( E_{\text{jet}} = m_{\chi}/2 \). Here we consider \( m_{\chi} = \epsilon_{X} = 1 \) TeV. Note that \( Y_{X} = n_{X}/s \).

FIG. 2. Same as Fig. 1, except for \( B_{\beta} = 1 \).

\[ m_{\chi}/2 \approx 10 \text{ GeV} \]

\[ \tau_{X} \approx 10^{4} \text{ sec} \]

If the lifetime of \( \chi \) is short (i.e., \( \tau_{X} \ll 10^{2} \) sec), \( \chi \) decays when the number density of the background electron (and positron) is still abundant and, in this case, energetic nucleons and nuclei are likely to scatter off the background electron and lose their energy without commit-
mittig the hadro-dissociation processes. In this case, the production of D and \( ^{6}\text{Li} \) is suppressed and the upper bound on \( m_{\chi} Y_{X} \) is not stringent.

Finally, we apply the above results to the primordial abundance of gravitinos. In the inflationary universe, gravitinos are produced by the scattering processes of the thermal particles. The yield variable of the gravitino is proportional to the reheating temperature \( T_{R} \) after the inflation, \( \tau_{X} = 1.5 \times 10^{-12} \times (T_{R}/10^{10} \text{ GeV})^{3} \) [3]. In addition, assuming the (massless) gauge boson and gaugino as the final state, lifetime of the gravitino is given by \( \tau_{3/2} \approx 4 \times 10^{8} \text{ sec} \times N_{G} (m_{3/2}/100 \text{ GeV})^{-3} \), where \( N_{G} \) is the number of the generators of the gauge group, and \( m_{3/2} \) is the gravitino mass. As examples, we consider two typical cases. The one is the case where the grav-
itino dominantly decays into the photon and photino, producing two hadron jets with \( E_{\text{jet}} = \frac{1}{3} m_{3/2} \); in this case, we take \( B_{\beta} = 10^{-3} \), \( N_{G} = 1 \), and \( \epsilon_{X} = \frac{1}{3} m_{3/2} \). The other is the case where the gravitino dominantly decay into the gluon and gluino, producing one hadron jet with \( E_{\text{jet}} = \frac{2}{3} m_{3/2} \); in this case we take \( B_{\beta} = 1 \), \( N_{G} = 8 \), and \( \epsilon_{X} = \frac{2}{3} m_{3/2} \). For these cases, we read off the upper bound on the reheating temperature for several values of the gravitino mass. The results are shown in Table I. It is seen that the constraint on \( T_{R} \) is much more stringent than that obtained for gravitino without hadronic decay.

Note added: While finalizing this letter, we found the paper by K. Jedamzik [20] which have some overlap with our analysis.

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| \( m_{3/2} \) (GeV) | \( B_{\beta} = 10^{-3} \) | \( B_{\beta} = 1 \) |
|-------------------|-----------------|-----------------|
| 100 GeV           | \( 2 \times 10^{6} \) GeV | \( 3 \times 10^{7} \) GeV |
| 300 GeV           | \( 3 \times 10^{6} \) GeV | \( 3 \times 10^{7} \) GeV |
| 1 TeV             | \( 3 \times 10^{7} \) GeV | \( 3 \times 10^{8} \) GeV |
| 3 TeV             | \( 2 \times 10^{7} \) GeV | \( 7 \times 10^{7} \) GeV |

TABLE I. Upper bounds on \( T_{R} \) for several values of \( m_{3/2} \).

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