Big Bang nucleosynthesis and particle dark matter

Karsten Jedamzik$^{1,4}$ and Maxim Pospelov$^{2,3,4}$

1 Laboratoire de Physique Théorique et Astroparticules, UMR5207-CNRS, Université Montpellier II, F-34095 Montpellier, France
2 Department of Physics and Astronomy, University of Victoria, Victoria, BC V8P 1A1, Canada
3 Perimeter Institute for Theoretical Physics, Waterloo, ON N2J 2W9, Canada
E-mail: jedamzik@lpta.univ-montp2.fr and pospelov@uvic.ca

New Journal of Physics 11 (2009) 105028 (20pp)
Received 13 March 2009
Published 15 October 2009
Online at http://www.njp.org/
doi:10.1088/1367-2630/11/10/105028

Abstract. We review how our current understanding of light element synthesis during the Big Bang nucleosynthesis era may help to shed light on the identity of particle dark matter.
1. Introduction

In the late 1940s and throughout the 1950s a number of visionary scientists, including Alpher, Fermi, Follin, Gamow, Hayashi, Herman and Turkevich, attempted to explain nuclear abundance patterns observed in the nearby universe, such as the peculiar high helium mass fraction \( Y_p \approx 0.25 \). This initially speculative work on an era of nucleosynthesis (element formation) in an expanding universe at very high temperature \( T \sim 10^9 \) K developed slowly but steadily over the coming decades into what is now known as the standard model of Big Bang nucleosynthesis (BBN). The idea that the universe may have undergone a very hot and dense early phase was triggered by the observations of Hubble in the 1920s, of the recession velocity of galaxies being proportional to their inferred distance from the Milky Way, which were most elegantly explained by a universe in expansion. The ‘expanding, hot Big Bang’ idea received further support by the observation of the cosmic microwave background radiation (CMBR) by Penzias and Wilson in 1965, believed to be the left-over radiation of the early universe. Detailed observational and theoretical studies of BBN as well as CMBR and Hubble flow have developed into the main pillars on which present-day cosmology rests.

BBN takes place between eras with (CMBR) temperatures \( T \sim 3 \) MeV and \( T \sim 10 \) keV, in the cosmic time window \( t \sim 0.1-10^4 \) s, and may be characterized as a freezeout from nuclear statistical equilibrium of a cosmic plasma at very low, \( \sim 10^{-9} \), baryon-to-photon number ratio (cf section 2), conditions that are not encountered in stars. It produces the bulk of \(^4\)He and \(^2\)H (D), as well as good fractions of \(^3\)He and \(^7\)Li observed in the current universe, whereas all other elements are believed to be produced by either stars or cosmic rays. In its standard version, it assumes a universe expanding according to the laws of general relativity, at a
given homogeneously distributed baryon-to-photon ratio $\eta_b$, with only standard model (SM) particle degrees of freedom excited, with negligible $\mu_\text{lepton} \ll T$ lepton chemical potentials and in the absence of any significant perturbations from primordial black holes, decaying particles, etc. By a detailed comparison of observationally inferred abundances (cf section 3) with those theoretically predicted, fairly precise constraints/conclusions about the cosmic conditions during the BBN era may thus be derived. BBN has been instrumental, for example, in constraining the contribution of extra ‘degrees of freedom’ excited in the early universe to the total energy density, such as predicted in many models of particle physics beyond the standard model. Such contributions may lead to an enhanced expansion rate at $T \sim 1$ MeV, implying an increased $^4\text{He}$ mass fraction. It is now known that aside from baryons and other subdominant components, not much more than the already known relativistic degrees of freedom (i.e. photons $\gamma$, electrons and positrons $e^\pm$, and three left-handed neutrinos $\nu$) could have been present during the BBN era. BBN is also capable of constraining very sensitively any non-thermal perturbations as induced, for example, by the residual annihilation of weak-scale dark matter particles (section 5), or by the decay of relic particles (section 7) and the possible concomitant production of dark matter. Moreover, the sheer presence of negatively charged or strongly interacting weak mass-scale particles during BBN may lead to dramatic shifts in the yields of light elements through catalytic phenomena. BBN may therefore constrain the properties and production mechanisms of dark matter particles, and this paper aims at revealing this connection.

It is possible that the biggest contribution of BBN towards understanding the dark matter enigma has already been made. Before the advent of precise estimates of the fractional contribution of baryons to the present critical density, $\Omega_b \approx (0.02273 \pm 0.00062)/h^2$, where $h$ is the Hubble constant in units of $100$ km s$^{-1}$ Mpc$^{-1}$, by detailed observations and interpretations of the anisotropies in CMBR [1], BBN was the only comparatively precise means to estimate $\Omega_b$. As it was not clear if the ‘missing’ dark matter was simply in the form of brown dwarfs, white dwarfs, black holes (formed from baryons) and/or $T \sim 10^6$ K hot gas, various attempts to reconcile a BBN era at large $\Omega_b \sim 1$ with the observationally inferred light element abundances were made. These included, for example, BBN in a baryon-inhomogeneous environment, left over possibly due to a first-order quantum chromodynamics (QCD) phase transition at $T \approx 100$ MeV, or BBN with late-decaying particles, such as the supersymmetric gravitino (for reviews, cf [2]–[4]). Only continuous theoretical efforts of this sort, and their constant ‘failure’ to account for large $\Omega_b$, gave way to the notion that dark matter must be in the form of ‘exotic’, non-baryonic material, such as a new fundamental particle investigated in the present book.

2. Standard BBN (SBBN) theory

SBBN theory is well understood and described in detail in many modern cosmology textbooks. The essence of SBBN is represented by a set of Boltzmann equations that may be written in the following schematic form:

$$\frac{dY_i}{dt} = -H(T)T \frac{dY_i}{dT} = \sum (\Gamma_{ij}Y_j + \Gamma_{ikl}Y_kY_l + \cdots),$$

(1)

where $Y_i = n_i/s$ are the time $t$ (or temperature $T$)-dependent ratios between number density $n_i$ and entropy density $s$ of light elements $i = ^1\text{H}, ^2\text{H}, ^4\text{He}$, etc; the $\Gamma_{ij}$ are generalized rates for element interconversion and decay that can be estimated by experiments and/or theoretical
calculations, and $H(T)$ is the temperature-dependent Hubble expansion rate. The system of equations (1) assumes thermal equilibrium, e.g. Maxwell–Boltzmann distributions for nuclei, which is an excellent approximation maintained by frequent interactions with the numerous $\gamma$'s and $e^\pm$'s in the plasma. The initial conditions for this set of equations are well specified: for temperatures much in excess of the neutron–proton mass difference, neutron and proton abundances are equal and related to the baryon to entropy ratio, $Y_{\text{neutron}} \simeq Y_{\text{proton}} \simeq \frac{1}{2} n_{\text{baryon}}/s$, while the abundance of all other elements is essentially zero. At temperatures relevant to BBN, the baryonic contribution to the Hubble rate is minuscule, and $H(T)$ is given by the standard radiation-domination formula:

$$H(T) = T^2 \times \left( \frac{8\pi^3 g_* G_N}{90} \right)^{1/2}, \text{ where } g_* = g_{\text{boson}} + \frac{7}{8} g_{\text{fermion}},$$

(2)

with the $g$ denoting the excited relativistic degrees of freedom. This expression needs to be interpolated in a known way across the brief epoch of the electron–positron annihilation, after which the photons become slightly hotter than neutrinos and $H(T) \simeq T_0^2/(178\,s)$, where $T_0$ is the photon temperature in units of $10^9\,K$. A number of well-developed integration routines that go back to an important work of Wagoner et al [5] allow one to solve the BBN system of equations numerically and obtain the freezeout values of the light elements. A qualitative ‘computer-free’ insight to these solutions can be found in e.g. [6].

In a nutshell, SBBN may be described as follows. After all weak rates fall below the Hubble expansion rate, the neutron-to-proton ratio freezes to $\sim 1/6$, subject to a slow further decrease to $\sim 1/7$ by $T_0 \simeq 0.85$ via neutron decay and out-of-equilibrium weak conversion. At this point, to a good approximation, all neutrons available will be incorporated into $^4\text{He}$, since it is the light element with the highest binding energy per nucleon. Synthesis of $^4\text{He}$, and all other elements, has to await the presence of significant amounts of $D$ (the ‘deuterium bottleneck’). This occurs rather late, at $T_0 \simeq 0.85$, since at higher $T_0$ the fragile $D$ is rapidly photodisintegrated by the multitude of CMBR photons. At $T_0 \lesssim 0.85$ the fairly complete nuclear burning of all $D$ then results in only trace amounts $O(10^{-5})$ of $D$ (and $^3\text{He}$) being left over after SBBN has ended. Elements with nucleon number $A > 4$ are even less produced due to appreciable Coulomb barrier suppression at such low $T_0$, resulting in only $O(10^{-10})$ of $^7\text{Li}$, and abundances of other isotopes are even lower. SBBN terminates due to the combination of a lack of free neutrons and the importance of Coulomb barriers at low $T$. In the following, a few more details are given.

$O(0.1)$ abundances: $^4\text{He}$. The $^4\text{He}$ mass fraction $Y_p$ is dependent on the timing of major BBN events, such as the neutron-to-proton freezeout at $T \simeq 0.7\,\text{MeV}$, post-freezeout neutron depletion before the deuterium bottleneck and the position of this bottleneck itself as a function of temperature. Consequently, $Y_p$ is dependent on such well-measured quantities as Newton’s constant, neutron–proton mass difference, neutron lifetime and deuterium binding energy, and to a much lesser degree on less precisely known values for the nuclear reaction rates. This sensitivity to the timing of the BBN events makes $^4\text{He}$ an important probe of the Hubble expansion rate, and of all possible additional non-standard contributions that could modify it. The SBBN predicts $Y_p$ with an impressive precision, $Y_p = 0.2486 \pm 0.0002$, where we use the most recent evaluation [7].

$O(10^{-5})$ abundances: $D$ and $^3\text{He}$. Deuterium and $^3\text{He}$ BBN predictions are more sensitive both to nuclear physics and to $\eta_b$ input. Reactions involving these elements are well measured, and with the current WMAP input SBBN is capable of making fairly precise predictions of these abundances: $D/H = 2.49 \pm 0.17 \times 10^{-5}$; $^3\text{He}/H = (1.00 \pm 0.07) \times 10^{-5}$.
$O(10^{-10})$ abundances: $^7$Li. Among all observable BBN abundances, $^7$Li is the most sensitive to the $\eta_b$ and nuclear physics inputs. The actual observable that BBN predicts is the combined abundance of $^7$Li and $^7$Be, as later in the course of the cosmological evolution $^7$Be is transformed into $^7$Li via electron capture. At the CMBR-measured value of the baryon-to-photon ratio $\eta_b$, more than 90% of primordial lithium is produced in the form of $^7$Be in the radiative capture process, $^4$He + $^3$He $\rightarrow$ $^7$Be + $\gamma$. As the rate for this process per each $^3$He nucleus is much slower than the Hubble rate, the output of $^7$Be is almost linearly dependent on the corresponding $S$-factor for this reaction. With recent improvement in its experimental determination [8]–[10], the current $\sim$15% accuracy prediction for $^7$Be + $^7$Li stands at $5.24^{+0.71}_{-0.65} \times 10^{-10}$ [7].

$O(10^{-14})$ and less abundances: $^6$Li and $A \geq 9$ elements. $^6$Li is formed in the BBN reaction

$$^4\text{He} + \text{D} \rightarrow ^6\text{Li} + \gamma, \quad Q = 1.47 \text{ MeV},$$

which at BBN temperatures is approximately four orders of magnitude suppressed relative to other radiative capture reactions such as $^4$He + $^3$H $\rightarrow$ $^7$Li + $\gamma$, and $\sim$seven/eight orders of magnitude suppressed relative to other photonless nuclear rates. The reason for the extra suppression is in a way accidental: it comes from the same charge to mass ratio for $^4$He and D, which inhibits the E1 transition, making this radiative capture extremely inefficient. This results in $O(10^{-14})$ level prediction for primordial $^6$Li, which is well below the detection capabilities. Heavier elements with $A \geq 9$ such as $^9$Be, $^{10}$B and $^{11}$B are never made in any significant quantities in the SBBN framework, and the main reason for this is the absence of stable $A = 8$ nuclei, as $^8$Be is underbound by 92 keV and decays to two $\alpha$.

3. Observed light element abundances

In the following, we will briefly discuss the observationally inferred light elements, element by element. Here the discussion will also include isotopes, or isotope ratios, such as $^6$Li, $^7$Be and $^3$He/$^2$H, which are not always considered in SBBN but are very useful to constrain deviations from SBBN.

3.1. $^4$He

The primordial $^4$He/H ratio is inferred from observations of hydrogen- and helium-emission lines in extragalactic low-metallicity HII regions and compact blue galaxies, illuminated by young star clusters. Two particular groups have performed such analysis for years now, with their most recent results $Y_p \approx 0.2477 \pm 0.0029$ [11] and $Y_p \approx 0.2516 \pm 0.0011$ [12]. These estimates are (surprisingly) considerably larger than earlier estimates by both groups (i.e. 0.239 and 0.242, respectively), explained in large part by a new estimate for HeI emissivities [13]. Other differences with respect to older studies, and/or between the two new studies themselves, are the adopted rates for collisional excitation of H- (He-) emission lines, corrections for a temperature structure in these galaxies (‘temperature variations’), corrections for the presence of neutral $^4$He (‘icf—ionization corrections’), as well as corrections for troughs in the stellar spectra at the position of the $^4$He- (H-) emission lines (‘underlying stellar absorption’). All of these may have an impact on the $\geq 1\%$ level. This, as well as the comparatively large change from earlier estimates (coincidentally going into the direction of agreement with the SBBN prediction of $Y_p \approx 0.248$), implies that a conservative estimate $Y_p \approx 0.249 \pm 0.009$ [14].

New Journal of Physics 11 (2009) 105028 (http://www.njp.org/)
(see also $Y_p \approx 0.250 \pm 0.004$) [15] of the error bar may be more appropriate when constraining perturbations of SBBN.

3.2. $D$

For the observational determination of primordial $D/H$ ratios, high-resolution observations of low-metallicity quasar absorption line systems (QALS) are employed (cf [16]–[20]). QALS are clouds of partially neutral gas that fall on the line of sight between the observer and a high-redshift quasar. The neutral component in these clouds yields absorption features, for example, at the redshifted position of the Lyman-$\alpha$ wavelength. For the very rare QALS of sufficiently simple velocity structure, one may compare the absorption at the Lyman-$\alpha$ position of H with that of D (shifted by 81 km s$^{-1}$) to infer a $D/H$ ratio. Here the low metallicity of these QALS is conducive to make one believe that stellar D destruction in such clouds is negligible. Currently, there exist only about 6–8 QALS with $D/H$ determinations. When averaged they yield typically $2.68-2.82 \times 10^{-5}$ [19]–[22] for the central value, with inferred statistical 1$\sigma$ error bars of $0.2-0.3 \times 10^{-5}$, comparing favorably to the SBBN prediction of $2.49 \pm 0.17 \times 10^{-5}$ [7] at the WMAP inferred $\eta_b$. Nevertheless, as an important cautionary remark, the various inferred $D/H$ ratios in QALS show a spread considerably larger than that expected from the above-quoted error bars only. This is usually a sign of the existence of unknown systematic errors. Until these systematics are better understood, primordial values as high as $D/H \approx 4 \times 10^{-5}$ should not be ruled out.

3.3. $^3He/D$

Observational determinations of $^3He/H$ ratios are possible within our galaxy, which is chemically evolved. The chemical evolution of $^3He$ is, however, rather involved, with $^3He$ known to be produced in some stars and destroyed in others. Furthermore, any D entering stars will be converted to $^3He$ by proton burning. The net effect of all this is an observed approximate constancy of $(D + ^3He)/H \approx 3.6 \pm 0.5 \times 10^{-5}$ [23] over the last few billion years in our galaxy. Whereas the relation of galactic observed $^3He/H$ ratios to the primordial one is obscure, the ratio of $^3He/D$ as observed in the presolar nebulae is invaluable in constraining perturbations of SBBN. This ratio $0.83^{+0.53}_{-0.25}$ [23] (where the error bars are obtained when using the independent 2$\sigma$ ranges of $^3He/H$ and $D/H$) provides a firm upper limit on the primordial $^3He/D$ [24]. This is because $^3He$ may be either produced or destroyed in stars, while D is always destroyed, such that the cosmological $^3He/D$ ratio may only grow in time.

3.4. $^7Li$

$^7Li/H$ ratios may be inferred from observations of absorption lines (such as the 6708A doublet) in atmospheres of low-metallicity galactic halo stars. When this is done for stars at low metallicity [Z], $^7Li/H$ ratios show a well-known anomaly (with respect to other elements), i.e. $^7Li/H$ ratios are constant over a wide range of (low) [Z] and some range of temperature (the ‘Spite plateau’). As most elements are produced by stars and/or cosmic rays, which themselves produce metallicity, the $^7Li$ Spite plateau is believed to be an indication of a primordial origin of this isotope. This interpretation is strengthened by the absence of any observed scatter in $^7Li$ abundance for such stars. There have been several observational determinations of $^7Li$ abundance on the Spite plateau. Most of them fall in the range $^7Li/H \approx 1-2 \times 10^{-10}$ such as
1.23^{+0.68}_{-0.33} \times 10^{-10} \ [25, 26] \text{ and } 1.1-1.5 \times 10^{-10} \ [27], \text{ with some being somewhat higher such as } 2.19 \pm 0.28 \times 10^{-10} \ [28]. \text{ Here differences may be due to differing methods of atmospheric temperature estimation. These values should be compared to the SBBN prediction } 5.24^{+0.71}_{-0.67} \times 10^{-10} \ [7] \text{ (with } 1\sigma \text{ error estimates), clearly indicating a conflict that is often referred to as the ‘lithium problem’. It is essentially ruled out that this problem be solved by either an erroneous atmospheric temperature determination or significant changes in } ^7\text{Li producing/destroying nuclear SBBN rates. There remain only two viable possibilities of a resolution to this statistically significant (4–5)σ problem. Firstly, it is conceivable that atmospheric } ^7\text{Li has been partially destroyed in such stars due to nuclear burning in the stellar interior. Although far from being understood, one may indeed construct (currently ad hoc) models that deplete } ^7\text{Li by a factor } \approx 2 \text{ in such stars, while respecting all other observations } [29, 30]. \text{ Secondly, it is possible that the lithium problem points directly towards physics beyond the SBBN model, possibly connected to the production of dark matter (cf section 7).}

3.5. ^6\text{Li and } ^9\text{Be}

The isotope of ^6\text{Li is usually not associated with BBN, as its standard BBN production } ^6\text{Li}/^7\text{Li} \sim 10^{-14} \text{ is very low. However, the smallest deviations from SBBN may already lead to important cosmological } ^6\text{Li abundances. It is therefore interesting that the existence of } ^6\text{Li has been claimed in about } \sim 10 \text{ low-metallicity stars } [27], \text{ with, nevertheless, each of these observations only at the } 2–4\sigma \text{ statistical significance level. Asplund et al } [27] \text{ infer an average of } ^6\text{Li}/^7\text{Li} \approx 0.044 (\text{corresponding to } ^6\text{Li}/^7\text{Li} \approx 6 \times 10^{-12}) \text{ for their star sample, whereas Cayrel et al } [34] \text{ infer } ^6\text{Li}/^7\text{Li} \approx 0.052 \pm 0.019 \text{ for the star HD84937 } [31]. \text{ Such claims, if true, would be of great interest, as the inferred } ^6\text{Li}/^7\text{Li} \text{ in very low metallicity stars is exceedingly hard to explain by cosmic ray production } [32], \text{ although } \text{in situ} \text{ production in stellar flares may be conceivable } [33]. \text{ Moreover, the } ^6\text{Li observations seem to be consistent with a plateau structure at low metallicity, as expected when originating right from BBN. However, recent work } [34] \text{ has cast a significant shadow over the claim of elevated } ^6\text{Li}/^7\text{Li} \text{ ratios at low } [Z]. \text{ Similar to } ^7\text{Li, } ^6\text{Li is inferred from observations of atmospheric stellar absorption features. Unlike in the case of } ^6\text{D and } ^6\text{H in QALS, the absorption lines of } ^7\text{Li and } ^6\text{Li are always blended together. } ^6\text{Li}/^7\text{Li} \text{ ratios may therefore be obtained only by observations of a minute asymmetry in the 6708 line. Such asymmetries could be due to } ^6\text{Li, but may also be due to asymmetric convective motions in stellar atmospheres. The analysis in Cayrel et al } [34] \text{ prefers the latter explanation.}

Unlike the case of } ^6\text{Li, the detection of } ^9\text{Be in many stars at low metallicities is not controversial. Observations of } ^9\text{Be } [35, 36] \text{ are far above the } O(10^{-18}) \text{ SBBN prediction, and exhibit linear correlation with oxygen, clearly indicating its secondary (spallation) origin } [37]. \text{ The lowest level of detected } ^9\text{Be}/^7\text{Li} \text{ is at } \sim \text{few } \times 10^{-14}, \text{ which translates into the limit on the primordial fraction at } 2 \times 10^{-13} \ [38], \text{ assuming no significant depletion of } ^9\text{Be in stellar atmospheres.}

4. Cascade nucleosynthesis from energy injection

The possibility that BBN may be significantly perturbed by the presence of energetic, non-thermal SM particles in the plasma first received detailed attention in the 1980s [40]–[47]. Although much of the pioneering work had been done, the first fully realistic calculations of coupled thermal nuclear reactions and non-energetic phenomena have been presented only
recently [48]–[51]. Energetic particles may be injected as products of the decay or annihilation of primordial black holes or supersymmetric Q-balls. The injected energetic photons $\gamma$, electron/positrons $e^\pm$, neutrinos $\nu$, muons $\mu^\pm$, pions $\pi$, nucleons and antinucleons $N$ and $\bar{N}$, gauge bosons $Z$ and $W^\pm$, etc may be considered as ‘cosmic rays’ of the early universe. In contrast to their present-day counterparts, and with the exception of neutrinos, these early cosmic rays thermalize rapidly within a small fraction of the Hubble time $H^{-1}(T)$ for all cosmic temperatures above $T \sim 1$ eV. This, of course, happens only after all unstable species (i.e. $\pi$, $\mu$, $Z$ and $W^\pm$) have decayed leaving only $\gamma$, $e^\pm$, $\nu$ and $N$. Many of the changes in BBN light-element production occur during the course of this thermalization. One often distinguishes between hadronically ($\pi$, $N$ and $\bar{N}$) and electromagnetically ($\gamma$ and $e^\pm$) interacting particles, mainly because the former may change BBN yields at times as early as $\tau \gtrsim 0.1$ s (i.e. $T \lesssim 3$ MeV), whereas the latter only have an impact for $\tau \gtrsim 10^3$ s (i.e. $T \lesssim 3$ keV). In the following, we summarize the most important interactions and outline the impact of such particles on BBN. For hadronically interacting particles these effects include:

1. $\pi^\pm$’s may cause charge exchange, i.e. $\pi^- + p \rightarrow \pi^0 + n$, at 1 MeV $\gtrsim T \gtrsim 300$ keV, thereby creating extra neutrons after $n/p$ freezeout and increasing the helium mass fraction $Y_p$.
2. Antinucleons $\bar{N}$ injected in the primordial plasma preferentially annihilate on protons, thereby raising the effective $n/p$ ratio and increasing $Y_p$.
3. At higher temperatures, neutrons $n$ completely thermalize through magnetic moment scattering on $e^\pm$ ($T \gtrsim 80$ keV), whereas protons $p$ do so through Coulomb interactions with $e^\pm$ and Thomson scattering off CMBR photons ($T \gtrsim 20$ keV). Any extra neutrons at $T \sim 40$ keV may lead to an important depletion of $^7$Be.
4. At lower temperatures, both energetic neutrons and protons may spall $^4$He, e.g. $n + ^4$He $\rightarrow ^3$H + $p + n + (\pi)$ or $n + ^4$He $\rightarrow D + p + 2n + (\pi)$. Both reactions are important as they may either increase the $^2$H abundance or lead to $^6$Li formation via the secondary non-thermal reactions of energetic $^3$H($^3$He) on ambient $\alpha$.

The main features of electromagnetic injection are:

1. Energetic $\gamma$ may pair-produce on CMBR photons, i.e. $\gamma + \gamma_{\text{CMBR}} \rightarrow e^- + e^+$, as long as their energy is above the threshold $E_C \approx m_{\gamma}^2 / 22T$ for this process. The created energetic $e^\pm$, in turn inverse Compton scatter, i.e. $e^\pm + \gamma_{\text{CMBR}} \rightarrow e^\pm + \gamma$, to produce further $\gamma$. Interactions with CMBR photons completely dominate interactions with matter due to the exceedingly small cosmic baryon-to-photon ratio $\eta$.
2. Only when $E_\gamma \lesssim E_C$ do interactions with matter become important. These include Bethe–Heitler pair production $\gamma + p(^4$He) $\rightarrow p(^4$He) + $e^+ + e^-$ and Compton scattering $\gamma + e^- \rightarrow \gamma + e^-$ off plasma electrons, as well as photodisintegration (see below).
3. A small fraction of $\gamma$ with $E_\gamma \lesssim E_C$ may first photodisintegrate D at $T \lesssim 3$ keV, when $E_C$ becomes larger than $E_b^D \approx 2.2$ MeV (the D binding energy), and later $^4$He at $T \lesssim 0.3$ keV since $E_b^{^4}$He $\approx 19.8$ MeV. Such processes may first cause D destruction and, later, D and $^3$He production, and more importantly $^3$He/D overproduction. They may also lead to $^6$Li production.

In the context of non-thermal energy injection related to particle dark matter, there are two very important processes that have a profound impact on $^6$Li and $^7$Li abundances and
Figure 1. Constraints on the abundance $\Omega_X h^2$ of relic particles decaying at $\tau_X$ assuming $M_X = 100$ GeV for the particle mass. The most stringent limits are given, from early to late times, by $^4$He, D, $^6$Li and $^3$He/D overproduction, respectively. The various lines are for different $\log_{10} B_h$, as labeled, where $B_h$ is the hadronic branching ratio. From [39].

deserve further comments. Energetic $^3$He and $^3$H produced via electromagnetic or hadronic energy injection (i.e. via spallation or photodisintegration) provide the possibility of efficient production of $^6$Li via non-thermal nuclear reactions on thermal $^4$He:

$$^3\text{H} + ^4\text{He} \rightarrow ^6\text{Li} + n, \quad Q = -4.78 \text{ MeV}$$

$$^3\text{He} + ^4\text{He} \rightarrow ^6\text{Li} + p, \quad Q = -4.02 \text{ MeV}. \quad (4)$$

For energies of projectiles $\sim 10$ MeV, the cross sections for these non-thermal processes are of the order of 100 mbn, and indeed $10^7$ times larger than the SBBN cross section for producing $^6$Li. This enhancement figure underlines $^6$Li sensitivity to non-thermal BBN, and makes it an important probe of energy injection mechanisms in the early universe.

Another important aspect of the non-thermal BBN is the possibility of alleviating the tension between the Spite plateau value and the predicted abundance of $^7$Li, e.g. by ‘solving the $^7$Li problem’. To achieve this, energy injection should occur in the temperature interval $60 \text{ keV} \gtrsim T \gtrsim 30 \text{ keV}$, i.e. during or just after $^7$Be synthesis. The essence of this mechanism consists in the injection of $\gtrsim 10^{-5}$ neutrons per baryon that will enhance $^7\text{Be} \rightarrow ^7\text{Li}$ interconversion followed by the $p$-destruction of $^7$Li via the thermal reaction sequence [48].

$$n + ^7\text{Be} \rightarrow p + ^7\text{Li}; \quad p + ^7\text{Li} \rightarrow ^4\text{He} + ^4\text{He}. \quad (5)$$

Note that this is the same mechanism that depletes $^7$Be in SBBN, but with elevated neutron concentration due to hadronic energy injection. This mechanism of depleting $^7$Be is tightly constrained by deuterium abundance, as extra neutrons could easily overproduce D.

Figure 1 summarizes constraints on abundance versus lifetime of relic decaying particles. It is convenient to measure the abundance in terms of $\Omega_X h^2$, the present-day fraction of total energy density if these particles were to remain stable. This quantity relates to the frequently used $\zeta = n_X M_X / s$ via $\zeta = 3.6639 \times 10^{-9} \text{ GeV} \Omega_X h^2$, where $n_X$ is particle number density, $M_X$ is particle mass and $s$ is entropy. It is seen that constraints become increasingly more stringent when the lifetime $\tau_X$ increases, implying also that, under generic circumstances, the production
of dark matter $X$ (with $\Omega_X h^2 \sim 0.1$) by the decay of a parent particle $Y \rightarrow X + \cdots$ at $\tau_X \gg 10^3$ s is extremely problematic, if at all possible.

5. Residual dark matter annihilation during BBN

Many dark matter candidates $X$ may be ultimately visible in our galaxy due to the cosmic rays they inject induced by residual $XX$ self-annihilations. This is, for example, the case for supersymmetric neutralinos, provided their annihilation products can be distinguished from astrophysical backgrounds. Residual annihilation events in the early universe may also be of importance as they may lead to cosmologically significant $^6$Li abundances [53] induced by the non-thermal nuclear reaction discussed in section 4. Given an $X$ annihilation rate $\langle \sigma v \rangle$ and $X$-density $n_X$, one may determine the approximate fraction $f_X$ of particles that annihilate in the early universe at temperature $T$:

$$f_X \approx \frac{1}{n_X} \frac{dn_X}{dt} \Delta t_H \approx \frac{\langle \sigma v \rangle n_X}{s} s^{-1},$$

where $\Delta t_H \approx H^{-1} = (90 M_{pl}^2/\pi^2 g T^4)^{1/2}$ is the characteristic Hubble time at $T$, $g$ is the appropriate particle statistical weight, $s = 4\pi^2/90 g T^3$ is radiation entropy and $\langle \cdots \rangle$ denotes a thermal average, which can be taken once the velocity dependence of $\sigma v$ is specified. Many scenarios for the production of dark matter envision a stable self-annihilating particle, typically a weakly interacting massive particle (WIMP), whose final asymptotic abundance is given by its annihilation rate. Straightforward considerations of thermal WIMP freezeout require the annihilation rate at $T_i^{th} \simeq 0.05 m_X$ to be $\langle \sigma v \rangle_i^{th} \approx 1$ pb nuc cm$^{-1}$ s$^{-1}$ if the $X$ particle is to be the dominant component of dark matter, $\Omega_X h^2 \approx 0.1$. Less straightforward but still plausible scenarios that include the non-thermal production of dark matter, e.g. via evaporation of Q-balls and/or decay of relic particles $Y \rightarrow X + \cdots$ with subsequent $X$ self-annihilation, may require $\langle \sigma v \rangle_i$ well in excess of $\langle \sigma v \rangle_i^{th}$ for $\Omega_X h^2 \approx 0.1$. We concentrate on the WIMP example, and parameterize the velocity dependence of $\langle \sigma v \rangle$ by $\langle \sigma v \rangle = (\sigma v)_0 S(v)$, with the chosen normalization $\langle \sigma v \rangle_0 = \langle \sigma v \rangle_i$. We are interested in finding the fraction of annihilating WIMP particles at $T < 10$ keV, a temperature scale below which $^6$Li is no longer susceptible to nuclear burning (destruction). Exploiting (6) at the freezeout temperature $T_i$, where $f_X(T_i) \approx 1$, as well as at an arbitrary other $T < T_i$, we obtain

$$f_X(T) \approx \left( \frac{g(T)}{g(T_i)} \right)^{1/2} \left( \frac{T}{T_i} \right) \langle S(v) \rangle_T \langle S(v) \rangle_{T_i}^{-1}$$

for $X \ll 1$. Several generic options are possible for the temperature scaling of the $S$-ratio in equation (7). If the s-wave annihilation is mediated by short-distance physics and occurs away from sharp narrow resonances, $\langle S(v) \rangle_T = \langle S(v) \rangle_{T_i} = 1$, where the second equality is due to our chosen normalization. Using this conservative assumption and equation (7), for a WIMP of mass $M_X = 100$ GeV, so that $g(T)/g(T_i) \simeq 0.1$, one finds that only a small fraction, $f_X \approx 6 \times 10^{-7}$, of $X$ particles has a chance to annihilate at $T \simeq 10$ keV and below. Nevertheless, even this tiny fraction is still sufficient to produce a $^6$Li abundance of $^6$Li/H $\approx 1.6 \times 10^{-12}$.

The existence of attractive Coulomb-like force of some strength $\alpha'$ in the WIMP sector may lead to a significant enhancement of annihilation at low temperatures/velocities [54], possibly leading to a much higher yield of $^6$Li. In this case, the Sommerfeld-like scaling $\sigma v \sim (\pi \alpha'/v)(1 - \exp(-\pi \alpha'/v))^{-1}$ enhances the annihilation at small $v \simeq \pi \alpha'$, i.e. $\langle S(v) \rangle_T \simeq \pi \alpha'/v$. This leads to a $\sim T^{-1/2}$ scaling of $\langle S(v) \rangle_T$ in (7) when $X$ particles are still in thermal

New Journal of Physics 11 (2009) 105028 (http://www.njp.org/)
Figure 2. Upper bound on the annihilation cross section of particle dark matter of mass $M_X$ from BBN. Here the upper line assumes annihilation into only electromagnetically interacting particles, whereas the two lower lines assume annihilation into a light quark–anti-quark pair. Adopted limits on the light element abundances are as indicated in the figure. From [52].

equilibrium with the plasma. After they have dropped out of thermal equilibrium, $\langle S(v) \rangle_T$ falls even more rapidly as $\sim T^{-1}$, with the net effect that weak mass scale $X$ particles usually have much smaller velocities at the end of BBN than in the Milky Way. Similarly, the presence of narrow resonances just above the $XX$ annihilation threshold may drastically boost the annihilation at low energies. Both mechanisms of enhancing the annihilation have been widely discussed (see e.g. [55]) in an attempt to link some cosmic ray anomalies to dark matter annihilation, as for example an elevated positron fraction $e^+/(e^- + e^+)$ observed by the PAMELA instrument [56].

Keeping the annihilation rate as a free parameter, figure 2 shows the upper limit on the effective annihilation cross section imposed by BBN. Here electromagnetically (upper line) and hadronically (lower lines) annihilating particle DM has been considered. The former is mostly constrained by overproduction of $^3\text{He}/\text{D}$ at $T \approx 0.1$ keV, while the latter by $^6\text{Li}/^7\text{Li}$ overproduction at $T \approx 10$ keV, such that the effective annihilation cross section refers to $\langle \sigma v \rangle$ at those temperatures. Due to the possibility of $^6\text{Li}$ destruction, a fairly conservative $^6\text{Li}/^7\text{Li} < 0.66$ constraint has also been considered. It is seen that much $^6\text{Li}$ may be produced by hadronic annihilations. Figure 3 shows darkmatter parameters that lead to the production of a $^6\text{Li}/^7\text{Li}$ ratio as claimed to be observed in star HD84937 at $1\sigma$ and $2\sigma$ (dark blue and light blue), respectively. Here a completely hadronic $XX \rightarrow q\bar{q}$ annihilation has been employed, an assumption that could be further confronted with constraints on antiproton fluxes in our galaxy. Electromagnetic annihilations, as often implied in recent dark matter interpretations of PAMELA [56], may also lead to $^6\text{Li}$, but only annihilations below $T \lesssim 0.3$ keV may do so efficiently. Note that the figure already implicitly assumes a factor of $\sim 3$–4 stellar destruction of $^7\text{Li}$ (and $^6\text{Li}$) to solve the lithium problem. It is therefore found that weak-mass-scale dark matter particles, if fairly light, and if annihilating into hadronically interacting particles, may account for all of the observed $^6\text{Li}$ in HD84937. The figure also shows by the orange-red-green-yellow areas the
Figure 3. Dark matter annihilation rate versus dark matter mass. The blue band shows parameters where $^6$Li due to residual dark matter annihilation may account for the $^6$Li abundance as inferred in HD84937 ($^6$Li/$^7$Li $\approx$ 0.014–0.09 at $2 - \sigma$), whereas the orange–red–green–yellow region shows where $^7$Li is efficiently destroyed, i.e. $^7$Li/$^7$H < 1.5, 2, 3 and $4 \times 10^{-10}$, respectively. Above the lower (upper) dashed line D/H exceeds $4 \times 10^{-5}$ ($5.3 \times 10^{-5}$), such that parameter space above the upper dashed line is ruled out by D overproduction. Scenarios between this line and the upper edge of the blue band are problematic since they severely overproduce $^6$Li. Dark matter annihilation into light quarks has been assumed. From [52].

dark matter parameters that would lead to significant $^7$Li destruction due to residual dark matter annihilations, with the nuclear destruction mechanism described in the previous section. Since those regions are much above the $^6$Li band, the possibility of a factor of 2 or more depletion in $^7$Li is severely constrained by $^6$Li overproduction. We note in passing that our constraints are significantly more conservative than those found in [57]. It is intriguing to realize that primordial $^6$Li production by residual (hadronic) dark matter annihilations dominates standard BBN $^6$Li production for annihilation rates as small as $\langle \sigma v \rangle \sim 10^{-27}$ cm$^3$ s$^{-1}$, well below $\langle \sigma v \rangle_{th}$. It is thus possible that in the most primeval gas clouds and in the oldest stars the bulk of $^6$Li is due to dark matter annihilations. Unfortunately, $^6$Li abundances as low as $^6$Li/$^7$H $\lesssim 10^{-12}$ are difficult to observe.

6. Catalyzed BBN (CBBN)

The idea of particle physics catalysis of nuclear reactions goes back to the 1950s, and muon-catalyzed fusion has been a subject of active theoretical and experimental research in nuclear physics. In recent years, there has been a significant interest towards a possibility of nuclear catalysis by hypothetical negatively charged particles that live long enough to participate in nuclear reactions at the BBN time [38, 51], [58]–[68] (see also [69]–[71] for earlier work on the subject). An essence of the idea is very simple: a negatively charged massive particle that we call
Table 1. Properties of bound states: Bohr radius $a_0 = 1/(Za m N)$, binding energies $E_b$ calculated for realistic charge radii and ‘photodissociation decoupling’ temperatures $T_0$.

| Bound state | $a_0$(fm) | $|E_b|$(keV) | $T_0$(keV) |
|-------------|-----------|-------------|-----------|
| $pX^-$      | 29        | 25          | 0.6       |
| $^4\text{He}X^-$ | 3.63   | 346         | 8.2       |
| $^7\text{Be}X^-$ | 1.03   | 1350        | 32        |
| $^8\text{Be}X^-$ | 0.91   | 34          | 1430      |

$X^-$ gets into a bound state with a nucleus of mass $m_N$ and charge $Z$, forming a large compound nucleus with charge $Z = 1$, mass $M_X + m_N$ and binding energy in the $O(0.1 – 1)$ MeV range. Once the bound state is formed, the Coulomb barrier is reduced signalling a higher ‘reactivity’ of the compound nucleus with other nuclei. But what proves to be the most important effects of catalysis are new reaction channels that may open up and avoid SBBN-suppressed production mechanisms [58], e.g. equation (3), thus clearing the path to synthesis of elements such as $^6\text{Li}$ and $^9\text{Be}$. Although in this section we discuss the catalysis by negatively charged heavy relics, this is not the only option for CBBN, as for example, strongly interacting relics may also participate and catalyze certain nuclear reactions.

Although the connection between dark matter and CBBN is not immediate—after all the dark matter may not be charged—it is possible that dark matter particles do have a relatively long-lived charged counterpart. One example of this kind is supersymmetry with the lightest supersymmetric particle (LSP), the gravitino, and the next-to-LSP (NLSP), a charged slepton, to be examined in the next section. In that case the decay of the NLSP is tremendously delayed by the smallness of the gravitino–lepton–slepton coupling $\sim G^2_1/2$. Another example in the same vein is the nearly degenerate stau–neutralino system, in which case the longevity of the charged stau against the decay to the dark matter neutralino is ensured as long as the mass splitting of the stau–neutralino system is below 100 MeV. Both the gravitino and neutralino in these two examples represent viable dark matter candidates. A very important aspect of CBBN is that the abundance of charged particles before they start decaying is given by their annihilation rate at freezeout. In most of the models their abundance is easily calculable, and if no special mechanisms are introduced to boost the annihilation rate, the abundance of charged particles per nucleon is not small, and in the typical ballpark of $Y_X \sim (0.001 – 0.1) \times m_X$/TeV.

6.1. Properties of the bound states

For light nuclei participating in BBN, we can assume that the reduced mass of the nucleus–$X^-$ system is well approximated by the nuclear mass, so that the binding energy is given by $Z^2 a^2 m_A/2$ when the Bohr orbit is larger than the nuclear radius. It turns out that this is a poor approximation for all nuclei heavier than $A = 4$, and the effect of the finite nuclear charge radius has to be taken into account. In table 1, we give the binding energies as well as the recombination temperature, defined as the temperature at which the photodissociation rate of bound states becomes smaller than the Hubble expansion rate. Below these temperatures bound states are practically stable, and the most important benchmark temperatures for CBBN are $T \sim 30, 8$ and $0.5$ keV, when ($^7\text{Be}X^-$), ($^4\text{He}X^-$) and ($pX^-$) can be formed without efficient suppression by the photodissociation processes. It is important to emphasize that these
properties of the bound states are generic to any CBBN realization: i.e. they are completely determined by the charge of $X^-$ and electromagnetic properties of nuclei, and thus are applicable to supersymmetry (SUSY) or non-SUSY models alike. It is also important to note that the $(^8\text{Be}X^-)$ compound nucleus is stable, which may open the path to synthesis of $A > 8$ elements in CBBN.

6.2. Catalysis at 30 keV: suppression of $^7\text{Be}$

When the universe cools to temperatures of 30 keV, the abundances of deuterium, $^3\text{He}$, $^4\text{He}$, $^7\text{Be}$ and $^7\text{Li}$ are already close to their freezeout values, although several nuclear processes remain faster than the Hubble rate. At these temperatures, a negatively charged relic can get into bound states with $^7\text{Be}$ and form a $(^7\text{Be}X^-)$ composite object. Once this object is formed, some new destruction mechanisms for $^7\text{Be}$ appear. For models with weak currents connecting nearly mass-degenerate $X^-X^0$ states, there is a very fast internal conversion followed by the $p$-destruction of $^7\text{Li}$:

$$(^7\text{Be}X^-) \rightarrow ^7\text{Li} + X^0; \quad ^7\text{Li} + p \rightarrow 2\alpha.$$ (8)

When $X^- \rightarrow X^0$ is energetically disallowed, the destruction of $^7\text{Be}$ can be achieved via the following chain:

$$(^7\text{Be}X^-) + p \rightarrow (^8\text{Be}X^-) + \gamma; \quad (^8\text{Be}X^-) \rightarrow (^8\text{Be}X^-) + e^+\nu,$$ (9)

which is much enhanced by the atomic resonances in the $(^7\text{Be}X^-)$ system [62].

The rates for both mechanisms may be faster than the Hubble rate, possibly leading to a sizable suppression of $^7\text{Be}$ abundance if $(^7\text{Be}X^-)$ bound states are efficiently forming. In other words, $(^7\text{Be}X^-)$ serves as a bottleneck for the CBBN depletion of $^7\text{Be}$. The recombination rate per $^7\text{Be}$ nucleus leading to $(^7\text{Be}X^-)$ is given by the product of recombination cross section and the concentration of $X^-$ particles. It can be easily shown that for $Y_X < 0.01$ the recombination rate is too slow to lead to a significant depletion of $^7\text{Be}$. Detailed calculations of recombination rate and numerical analyzes of the CBBN at 30 keV [62, 67] indicate that the suppression of $^7\text{Be}$ by a factor of 2 is possible for $Y_X \geq 0.1$ if only mechanism (9) is operative, and for $Y_X \geq 0.02$ if internal conversion (8) is allowed.

6.3. Catalysis at 8 keV: enhancement of $^6\text{Li}$ and $^9\text{Be}$

As the universe continues to cool below 10 keV, an efficient formation of $(^4\text{He}X^-)$ bound states becomes possible. With the reasonable assumption of $Y_X < Y_{^4\text{He}}$, the rate of formation of bound states per $X^-$ particle is given by the recombination cross section and the concentration of the helium nuclei. Numerical analysis of recombination reveals that at $T \simeq 5$ keV, about 50% of available $X^-$ particles will be in bound states with $^4\text{He}$ [58].

As soon as $(^4\text{He}X^-)$ is formed, new reaction channels open up. In particular, a photonless thermal production of $^6\text{Li}$ becomes possible:

$$(^4\text{He}X^-) + D \rightarrow ^6\text{Li} + X^-, \quad Q \simeq 1.13 \text{ MeV},$$ (10)

which exceeds the SBBN production rate by $\sim$six orders of magnitude. The production of $^9\text{Be}$ may also be catalyzed, possibly by many orders of magnitude relative to the SBBN values, with the following thermal nuclear chain [66]:

$$(^4\text{He}X^-) + ^4\text{He} \rightarrow (^8\text{Be}X^-) + \gamma; \quad (^8\text{Be}X^-) + n \rightarrow ^9\text{Be} + X^-.$$ (11)
Both reactions at these energies are dominated by the resonant contributions, although the efficiency of the second process in (11) is not fully understood.

Current estimates/calculations of the CBBN rates are used to determine the generic constraints on lifetimes/abundances of charged particles. The essence of these limits is displayed in figure 4, which shows that for typical $X^-$ abundances the lifetime of the charged particles would have to be limited by a few thousand seconds! This is the main conclusion to be learned from CBBN. Note that while, to lowest order, non-thermal BBN is sensitive to the energy density of decaying particles, CBBN processes are controlled by the number density of $X^-$, which underlines the complementary character of these constraints. In some models, where both catalysis and cascade nucleosynthesis occur, catalysis dominates the cascade production of $^6$Li for all particles with hadronic branching ratio $B_h \lesssim 10^{-2}$ [65], whereas $^7$Li destruction is usually dominated by cascade effects unless $B_h \lesssim 10^{-4}$. $^9$Be production, on the other hand, is conceivable only through catalysis.

6.4. Catalysis below 1 keV and nuclear uncertainties

Finally, we comment on the possibility of ($pX^-$) catalysis of nuclear reactions, discussed in [64, 70]. Although it is conceivable that the absence of the Coulomb barrier for this compound nucleus may lead to significant changes of SBBN/CBBN predictions, in practice it turns out that in most cases ($pX^-$)-related mechanisms are of secondary importance. The large radius and shallow binding of this system leads to a fast charge-exchange reaction on helium, $(pX^-) + ^4\text{He} \rightarrow (^4\text{He}X^-) + p$, that reduces the abundance of ($pX^-$) below $10^{-6}$ relative to hydrogen, as long as $Y_{X^-} \lesssim Y_{^4\text{He}}$, making further reactions inconsequential for any observable element [38]. In the less likely case, $Y_{X^-} \gtrsim Y_{^4\text{He}}$, significant late-time processing due to ($pX^-$) bound states may still occur. Such late-time BBN, nevertheless, typically leads to observationally unacceptable final BBN yields.

Unlike in the SBBN case and even in cascade nucleosynthesis that utilizes mostly measured nuclear reaction rates, CBBN rates cannot be measured in the laboratory, and significant nuclear theory input for the calculation of the reaction rates is required. However, since $X^-$ participates

---

**Figure 4.** The left panel shows CBBN constraints on the abundance versus lifetime of $X^-$. The red cross corresponds to a point in parameter space, for which the temporal development of $^6\text{Li}$ and $^9\text{Be}$ is shown in the right panel, following [38].
Figure 5. Parameter space in the relic decaying particle abundance times hadronic branching ratio $B_h$, i.e. $\Omega_X h^2 B_h$, and lifetime $\tau_X$ plane, where $^7\text{Li}$ is significantly reduced (red and blue) and $^6\text{Li}$ is efficiently produced (green and blue). See text for further details. From [72].

only in electromagnetic interactions, such calculations are feasible, and dedicated nuclear theory studies [61] in this direction have already commenced. The reaction rates for some CBBN processes, such as (9) and (10), are already known within a factor of 2 accuracy, and detailed calculations for $^9\text{Be}$ synthesis are underway [68].

7. Dark Matter production during BBN: NLSP $\rightarrow$ LSP example

Dark matter particles may be produced by the decay of relic parent particles $X$ during BBN. Examples, well studied by different groups, include the production of gravitino-LSP dark matter by NLSP decays (often charged sleptons or neutralinos) or the production of neutralino dark matter by heavier gravitinos. Other conceivable possibilities include the production of superweakly interacting Kaluza–Klein dark matter, and more generally cascade decays to any superweakly interacting dark matter candidates. In case of charged NLSP decays, both nonthermal and CBBN processes must be accounted for. In the framework of gravitino LSP/stau NLSP, the lifetime of the charged slepton in the limit of $m_\tilde{G} \ll M_{\text{NLSP}}$ is given by

$$\tau_{\text{NLSP}} \approx 2.4 \times 10^4 \text{ s} \times \left(\frac{M_{\text{NLSP}}}{300 \text{ GeV}}\right)^{-5} \left(\frac{m_\tilde{G}}{10 \text{ GeV}}\right)^2,$$

where $M_{\text{NLSP}}$ and $m_\tilde{G}$ denote NLSP and gravitino mass, respectively.

It becomes exceedingly more difficult with increasing $\tau_X$ to obtain observational consistency with inferred primordial abundances (cf figure 1). BBN therefore plays an important role in constraining such scenarios (cf sections 4 and (6)). However, BBN may not only constrain but also favor particular scenarios, if current discrepancies with $^6\text{Li}$ and $^7\text{Li}$ abundances are to be taken seriously. Both trends, the reduction of $^7\text{Li}$ and the production of $^6\text{Li}$ via mechanisms described in sections 4 and (6), are seen in figure 5. There the red area shows decaying particle parameter space resulting in more than a factor of 2 suppressed
$^7\text{Li}$ abundance relative to the SBBN prediction, and the green area shows regions where significant $^6\text{Li}$ production ($0.015 \lesssim ^6\text{Li}/^7\text{Li} \lesssim 0.3$) occurs. In the overlap of these areas, the blue area, both effects may be achieved simultaneously [48, 51, 73]. Figure 5 also shows the prediction of supersymmetric scenarios with the gravitino LSP, for some representative values of other supersymmetric mass parameters. In particular, the gray dots show predictions of stau NLSPs with gravitino LSPs of mass $m_{\tilde{G}} = 50 \text{GeV}$ within the so-called constrained minimal supersymmetric SM (CMSSM), whereas the blue dots show the case of neutralino NLSPs decaying into $m_{\tilde{G}} = 100 \text{MeV}$ gravitino LSPs within the gauge-mediated supersymmetry breaking scenario. It is seen that both scenarios naturally cross the region of $^7\text{Li}$ destruction. The assumption underlying these models is a thermal freezeout abundance of the NLSP. Since this typically leads to NLSP abundances, $10^{-3} \lesssim \Omega_{\text{NLSP}} \lesssim 10^3$, and taking into account that gravitino energy density due to NLSP decays is $\Omega_{\tilde{G}} = \Omega_{\text{NLSP}} (m_{\tilde{G}}/M_{\text{NLSP}})$, the resulting $\Omega_{\tilde{G}}$ produced in such scenarios may come close to the observed dark matter density. This is particularly the case for heavy gravitinos $m_{\tilde{G}} \sim 100 \text{GeV}$ in the CMSSM, for which more detailed results are shown in figure 6. It is intriguing, and perhaps purely coincidental, that when resolving the tension between observed and predicted $^7\text{Li}$ abundance by staus decaying into gravitinos, the resulting gravitino abundance may account for all the dark matter. For stau decay times $\tau \approx 10^3 \text{s}$, it is furthermore possible to synthesize a primordial $^6\text{Li}$ abundance as claimed to be observed in low-metallicity stars. Moreover, although less certain, the same parameter space could also lead to an important $^9\text{Be}$ abundance due to catalytic effects (cf section 6), as indicated by the cross-hatched region in figure 6. Finally, since produced by decays, gravitino dark matter in such scenarios is significantly warm, with free-streaming velocities of the same order as those of an $m \approx 3 \text{ keV}$ early freezing-out relic particle, which has important implications for the small-scale structures in the present-day universe. It is therefore not impossible that at some time in the future, anomalies in the primordial light elements may be understood as signs
of dark matter. Nevertheless, independent verification by particle accelerators, such as the LHC, is required. Unfortunately, scenarios as presented in figure 6 require staus of mass $m_\tilde{\tau} \gtrsim 1$ TeV, too heavy to be produced at the LHC.

8. Conclusions

Even though the concept of BBN is more than 50 years old, it continues to be relevant due to the constant progress in nuclear physics and astrophysics, and the refined quality of calculations. In this paper, we have reviewed the status of BBN, and have shown how New Physics can modify the synthesis of light element abundances. All three generic ways, extra degrees of freedom modifying the Hubble expansion during BBN, energy injection due to annihilation or decay of heavy particles and particle catalysis of BBN reactions, are highly relevant to the physics associated with particle dark matter or with particles intimately tied to dark matter. We have illustrated how the existing overall concordance between predicted elemental abundances and observations leads to some very nontrivial constraints on the properties of DM particles and their companions (e.g. stau–gravitino system).

Perhaps even more intriguing is the current discrepancy between the observed and standard BBN predicted abundance of $^7$Li at the level of 2–3. This discrepancy has hirmed up since the last unknown SBBN parameter, the baryon-to-photon ratio, has been determined with better than 5% accuracy by recent high-precision CMB experiments. At this moment it is premature to talk on how the $^7$Li problem is resolved, but it is nonetheless intriguing that certain models with unstable particles are capable of alleviating this discrepancy. The continuing improvement of the observational determination of primordial light element abundances as well as future breakthroughs in electroweak-scale particle physics would help in solving this important problem.

References

[1] Dunkley J et al (WMAP) 2008 arXiv:0803.0586
[2] Malaney R A and Mathews G J 1993 Phys. Rep. 229 145–219
[3] Sarkar S 1996 Rep. Prog. Phys. 59 1493–610 arXiv:hep-ph/9602260
[4] Iocco F, Mangano G, Miele G, Pisanti O and Serpico P D 2008 arXiv:0809.0631
[5] Wagener R V, Fowler W A and Hoyle F 1967 Astrophys. J. 148 3–49
[6] Mukhanov V F 2004 Int. J. Theor. Phys. 43 669–93 arXiv:astro-ph/0303073
[7] Cyburt R H, Fields B D and Olive K A 2008 arXiv:0808.2818
[8] Nara Singh B S, Hass M, Nir-El Y and Haquin G 2004 Phys. Rev. Lett. 93 262503 arXiv:nucl-ex/0407017
[9] Gyurky G et al 2007 Phys. Rev. C 75 035805 arXiv:nucl-ex/0702003
[10] Brown T A D et al 2007 Phys. Rev. C 76 055801 arXiv:0710.1279
[11] Peimbert M, Luridiana V and Peimbert A 2007 (arXiv:astro-ph/0701580)
[12] Izotov Y I, Thuan T X and Stasinska G 2007 Astrophys. J. 662 15–38 (arXiv:astro-ph/0702072)
[13] Porter R L, Ferland G J and MacAdam K B 2007 Astrophys. J. 657 327–37 (arXiv:astro-ph/0611579)
[14] Olive K A and Skillman E D 2004 Astrophys. J. 617 29 (arXiv:astro-ph/0405588)
[15] Fukugita M and Kawasaki M 2006 Astrophys. J. 646 691
[16] Burles S and Tytler D 1998 Astrophys. J. 499 699 (arXiv:astro-ph/9712108)
[17] Levshakov S A, Dessauges-Zavadsky M, D’Odorico S and Molaro P 2001 arXiv:astro-ph/0105529
[18] Crighton N H M, Webb J K, Ortiz-Gill A and Fernandez-Soto A 2004 Mon. Not. R. Astron. Soc. 355 1042 (arXiv:astro-ph/0403512)

New Journal of Physics 11 (2009) 105028 (http://www.njp.org/)
[19] O’Meara J M et al 2006 Astrophys. J. 649 L61–6 (arXiv:astro-ph/0608302)
[20] Pettini M, Zych B J, Murphy M T, Lewis A and Steidel C C 2008 arXiv:0805.0594
[21] Fields B and Sarkar S 2006 arXiv:astro-ph/0601514
[22] Steigman G 2007 Annu. Rev. Nucl. Part. Sci. 57 463–91 (arXiv:0712.1100)
[23] Geiss J and Gloeckler G 2007 Space Sci. Rev. 130 5
[24] Sigl G, Jedamzik K, Schramm D N and Berezinsky V S 1995 Phys. Rev. D 52 6682–93 (arXiv:astro-ph/9509309)
[25] Ryan S G, Norris J E and Beers T C 1999 Astrophys. J. 523 654–77 (arXiv:astro-ph/9903059)
[26] Hosford A, Ryan S G, Perez A E G, Norris J E and Olive K A 2008 arXiv:0811.2506
[27] Asplund M, Lambert D L, Nissen P E, Primas F and Smith V V 2006 Astrophys. J. 644 229–59 (arXiv:astro-ph/0510636)
[28] Bonifacio P et al 2002 arXiv:astro-ph/0204332
[29] Richard O, Michaud G and Richer J 2005 Astrophys. J. 619 538–48 (arXiv:astro-ph/0409672)
[30] Korn A et al 2006 Nature 442 657–9 (arXiv:astro-ph/0608201)
[31] Cayrel R, Spite M, Spite F, Vangioni-Flam E, Casse M and Audouze J 1999 Astron. Astrophys. 343 923 (arXiv:9901205)
[32] Prantzos N 2005 arXiv:astro-ph/0510122
[33] Tatischeff V and Thibaud J P 2007 Astron. Astrophys. 469 265 (arXiv:astro-ph/0610756)
[34] Cayrel R et al 2007 arXiv:0708.3819
[35] Primas F, Asplund M, Nissen P E and Hill V 2000 arXiv:astro-ph/0009482
[36] Boesgaard A M and Novicki M C 2006 Astrophys. J. 641 1122–30 arXiv:astro-ph/0512317
[37] Fields B, Olive K A and Vangioni-Flam E 2005 Astrophys. J. 623 1083–91 (arXiv:astro-ph/0411728)
[38] Pospelov M, Pradler J and Steffen F D 2008 J. Cosmol. Astropart. Phys. JCAP11(2008)020 (arXiv:0807.4287)
[39] Jedamzik K 2006 Phys. Rev. D 74 103509 (arXiv:hep-ph/0604251)
[40] Ellis J R, Nanopoulos D V and Sarkar S 1985 Nucl. Phys. B 259 175
[41] Levitan Y L, Sobol I M, Khlopov M Y and Chechetkin V M 1988 Sov. J. Nucl. Phys. 47 109–15
[42] Dimopoulos S, Esmailzadeh R, Hall L J and Starkman G D 1988 Astrophys. J. 330 545
[43] Reno M H and Seckel D 1988 Phys. Rev. D 37 3441
[44] Dimopoulos S, Esmailzadeh R, Hall L J and Starkman G D 1989 Nucl. Phys. B 311 699
[45] Ellis J R, Gelmini G B, Lopez J L, Nanopoulos D V and Sarkar S 1992 Nucl. Phys. B 373 399–437
[46] Khlopov M Y, Levitan Y L, Sedelnikov E V and Sobol I M 1994 Phys. At. Nucl. 57 1393–7
[47] Kawasaki M and Moroi T 1995 Astrophys. J. 452 506 (arXiv:astro-ph/9412055)
[48] Jedamzik K 2004 Phys. Rev. D 70 063524 (arXiv:astro-ph/0402344)
[49] Kawasaki M, Kohri K and Moroi T 2005 Phys. Lett. B 625 7–12 (arXiv:astro-ph/0402490)
[50] Kawasaki M, Kohri K and Moroi T 2005 Phys. Rev. D 71 083502 (arXiv:astro-ph/0408462)
[51] Cyburt R H, Ellis J R, Fields B D, Olive K A and Spanos V C 2006 J. Cosmol. Astropart. Phys. JCAP11(2006)014 (arXiv:astro-ph/0608562)
[52] Jedamzik K Pospelov M and Ritz A in preparation
[53] Jedamzik K 2004 Phys. Rev. D 70 083510 (arXiv:astro-ph/0405583)
[54] Hisano J, Matsumoto S and Nojiri M M 2004 Phys. Rev. Lett. 92 031303 (arXiv:hep-ph/0307216)
[55] Pospelov M and Ritz A 2009 Phys. Lett. B 671 391–7 (arXiv:0810.1502)
[56] Adriani O et al 2008 arXiv:0810.4995
[57] Hisano J, Kawasaki M, Kohri K and Nakayama K 2009 Phys. Rev. D 79 063514 (arXiv:0810.1892)
[58] Pospelov M 2007 Phys. Rev. Lett. 98 231301 (arXiv:hep-ph/0605215)
[59] Kohri K and Takayama F 2007 Phys. Rev. D 76 063507 (arXiv:hep-ph/0605243)
[60] Kaplinghat M and Rajaraman A 2006 Phys. Rev. D 74 103004 (arXiv:astro-ph/0606209)
[61] Hamaguchi K, Hatsuda T, Kamimura M, Kino Y and Yanagida T T 2007 Phys. Lett. B 650 268–74 (arXiv:hep-ph/0702274)

New Journal of Physics 11 (2009) 105028 (http://www.njp.org/)
[62] Bird C, Koopmans K and Pospelov M 2008 *Phys. Rev.* D **78** 083010 (arXiv:hep-ph/0703096)
[63] Jittoh T *et al.* 2007 *Phys. Rev.* D **76** 125023 (arXiv:0704.2914)
[64] Jedamzik K 2008 *Phys. Rev.* D **77** 063524 (arXiv:0707.2070)
[65] Jedamzik K 2008 *J. Cosmol. Astropart. Phys.* JCAP03(2008)008 (arXiv:0710.5153)
[66] Pospelov M 2007 arXiv:0712.0647
[67] Kusakabe M, Kajino T, Boyd R N, Yoshida T and Mathews G J 2007 arXiv:0711.3858
[68] Kamimura M, Kino Y and Hiyama E 2008 arXiv:0809.4772
[69] De Rujula A, Glashow S L and Sarid U 1990 *Nucl. Phys.* B **333** 173
[70] Dimopoulos S, Eichler D, Esmailzadeh R and Starkman G D 1990 *Phys. Rev.* D **41** 2388
[71] Rafelski J, Sawicki M, Gajda M and Harley D 1991 *Phys. Rev.* A **44** 4345
[72] Bailly S, Jedamzik K and Moultaka G 2008 arXiv:0812.0788
[73] Cumberbatch D *et al.* 2007 *Phys. Rev.* D **76** 123005 (arXiv:0708.0095)