

MeV-mass dark matter and primordial nucleosynthesis

Pasquale D. Serpico and Georg G. Raffelt
Max-Planck-Institut für Physik (Werner-Heisenberg-Institut), Föhringer Ring 6, 80805 München, Germany

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The annihilation of new dark matter particles with masses \( m_X \) in the MeV range may account for the galactic positrons that are required to explain the 511 keV \( \gamma \)-ray flux from the galactic bulge. We study the impact of MeV-mass thermal relic particles on the primordial synthesis of \(^4\)H, \(^3\)He, and \(^7\)Li. If the new particles are in thermal equilibrium with neutrinos during the nucleosynthesis epoch they increase the helium mass fraction for \( m_X \lesssim 10 \) MeV and are thus disfavored. If they couple primarily to the electromagnetic plasma they can have the opposite effect of lowering both helium and deuterium. For \( m_X = 4-10 \) MeV they can even improve the overall agreement between the predicted and observed \(^4\)H and \(^3\)He abundances.

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I. INTRODUCTION

In a recent series of papers the intriguing possibility was explored that the cosmic dark matter consists of new elementary particles with masses in the MeV range \[ \text{[1,2,3,4,5]} \]. While weakly interacting massive particles (WIMPs) as dark matter candidates, notably supersymmetric particles, are usually thought to have masses exceeding tens of GeV, it is not difficult to come up with viable MeV-mass candidates such as sterile neutrinos \[ \text{[6]} \] or axinos \[ \text{[7]} \]. The remarkable aspect of the MeV-mass dark matter hypothesis is that the cosmic dark matter consists of these dark matter particles in the galactic bulge can be in thermal equilibrium with the electromagnetic plasma they can have the opposite effect of lowering both helium and deuterium. For \( m_X = 4-10 \) MeV they can even improve the overall agreement between the predicted and observed \(^4\)H and \(^3\)He abundances.

Quite on the contrary, it was argued that the annihilation rate much larger than given by ordinary weak interactions. Contrary to naive intuition, such particles are not excluded by any obvious laboratory measurement or astrophysical argument \[ \text{[1,2,3,4,5]}. \]

In addition to our BBN study presented in Sec. II we briefly consider two other possible consequences of MeV-mass dark matter particles. In Sec. III we note that the same mechanism proposed to explain the \( \gamma \)-ray signature from the galactic bulge should also produce a diffuse background of low-energy cosmic-ray positrons in the solar neighborhood. This argument strengthens the conclusion reached independently in Refs. \[ \text{[1,2,3,4,5]}. \]

II. MEV-MASS PARTICLES AND BIG-BANG NUCLEOSYNTHESIS

A. Impact of New Particles

Thermal relic particles freeze out at a cosmic temperature \( T_F \) that for weakly interacting particles is about 1 MeV. If the particle mass is somewhat larger than \( T_F \) the number density is suppressed by annihilations so that the relic density is maximal for \( m_X \) of order \( T_F \), i.e. for neutrinos the relic mass density would be maximal for MeV-range masses and overclose the universe by about five orders of magnitude \[ \text{[14]}. \]

Therefore, MeV-mass thermal relics must interact far more strongly than neutrinos so that they annihilate more efficiently to reduce their density to a level compatible with the dark-matter abun-
dance. We briefly summarize the relationship between the annihilation cross section and the relic density in Appendix A. The result reveals that MeV-mass dark matter particles must have been in thermal equilibrium throughout most of the primordial nucleosynthesis epoch so that the exact final dark-matter abundance is not relevant for our study. We will use the equilibrium assumption in our numerical implementation that was performed by a modified version of the new BBNCode recently developed in Naples and documented in Ref. [15].

The impact of our new particles on BBN differs qualitatively and quantitatively depending on the dominant annihilation and scattering channels. Before turning to specific cases we discuss how the BBN equations and inputs are affected. To this end we first note that the new particles enter BBN through the second Friedmann equation $H^2 = (8\pi/3) G_N \rho_{tot}$ by their contribution to the total mass-energy density $\rho_{tot}$. We consider the $X$-particles to be in perfect thermal equilibrium and assume that the number densities of particles and antiparticles are the same for those cases where $X \neq \bar{X}$. The number density, pressure, and energy density contributed by the new particles is thus taken to be

$$n_X = \frac{g_X}{2\pi^2} T_X^3 \int_x^\infty \frac{dy (y^2 - x^2)^{1/2}}{e^{y^2} + 1},$$

$$P_X = \frac{g_X}{6\pi^2} T_X^4 \int_x^\infty \frac{(y^2 - x^2)^{3/2}}{e^{y^2} + 1},$$

$$\rho_X = \frac{g_X}{2\pi^2} T_X^4 \int_x^\infty \frac{y^2 (y^2 - x^2)^{1/2}}{e^{y^2} + 1},$$

(1)

where $x = m_X/T_X$ and the sign $+$ ($-$) refers to fermion (boson) statistics. $T_X$ is identical with the ambient neutrino temperature $T_\nu$ for $X$-particles that dominantly couple to neutrinos, while $T_h$ for the electromagnetically coupled case. Since $y > x > 0$ and

$$(e^y + 1)^{-1} = e^{-y} \sum_{n=0}^\infty (\mp 1)^n e^{-ny}$$

(2)

we may expand the thermal integrals as

$$\int_x^\infty \frac{dy (y^2 - x^2)^{1/2}}{e^{y^2} + 1} = x^2 \sum_{n=0}^\infty (\mp 1)^n K_2(1 + n)x \frac{1}{1 + n},$$

$$\int_x^\infty \frac{(y^2 - x^2)^{3/2}}{e^{y^2} + 1} = 3x^2 \sum_{n=0}^\infty (\mp 1)^n K_2(1 + n)x \frac{1}{(1 + n)^2},$$

$$\int_x^\infty \frac{y^2 (y^2 - x^2)^{1/2}}{e^{y^2} + 1} = \int_x^\infty \frac{dy (y^2 - x^2)^{1/2}}{e^{y^2} + 1} + x^2 \int_x^\infty \frac{(y^2 - x^2)^{1/2}}{e^{y^2} + 1}$$

$$= \sum_{n=0}^\infty (\mp 1)^n \left( 3x^2 K_2(1 + n)x \frac{1}{(1 + n)^2} + x^3 K_1(1 + n)x \frac{1}{1 + n} \right),$$

(3)

where $K_i(x)$ is the special Bessel function of order $i$. In the numerical code the series are truncated at $n = 4$. The error of the physical quantities is always \leq 0.18%.

A more subtle effect is caused by the conservation of entropy that implies a modified relation between the neutrino temperature $T_\nu$ and that of the electromagnetic plasma $T \equiv T_\gamma$. The cosmic entropy density is

$$s = \frac{2\pi^2}{45} g_{*s} T^3 = \frac{2\pi^2}{45} T^3 \sum_i g_{*s}^i,$$

(4)

where

$$g_{*s}^i(T) = \frac{2\pi^2}{45 T^3} \rho_i + P_i T_i$$

$$= \left( \frac{T}{T_\nu} \right)^4 \left( \sum_{i=0}^{\text{baryons}} g_i + \frac{7}{8} \sum_{i=\text{fermions}} g_i \right).$$

(5)

The second equality applies for relativistic species, i.e. in the absence of warm species where $m \sim T$. Note that $g_{*s}^i$ is always considered as a function of the temperature $T$ of the electromagnetic plasma, not of $T_i$ (a different notation is used in Ref. [13]). After neutrino decoupling at $T_D \approx 2.3$ MeV [16], the entropy in a comoving volume is separately conserved for the “neutrino plasma” and the electromagnetic one so that

$$\left[ \frac{g_{*s}^\nu(T_D)}{g_{*s}^\nu(T)} \right] \sum_i g_{*s}^i(T_D) = 1. \quad \left[ \frac{g_{*s}^\nu(T_D)}{g_{*s}^\nu(T)} \right] \sum_i g_{*s}^i(T_D)$$

(6)

For new particles being only electromagnetically coupled this equation simplifies to

$$\frac{T}{T_\nu} = \left[ \frac{g_{*s}^\nu(T_D)}{g_{*s}^\nu(T)} \right] \left[ \frac{g_{*s}^\nu(T_D)}{g_{*s}^\nu(T)} \right],$$

(7)

thus giving a higher temperature ratio relative to the standard case. This simulates the effect of $N_{\text{eff}} < 3$ thermally excited neutrino species at BBN and thus a reduced primordial helium abundance. The opposite is true for $\nu$-coupled new particles.

The modified $T_\nu(T)$ relation also affects the $n \leftrightarrow p$ weak rates that depend on both $T$ and $T_\nu$ through the phase-space dependence of the initial and final states of the processes. We implemented the modification of these rates in a perturbative way by introducing the small parameter

$$\delta(T) = \frac{T_\nu^0(T) - T_\nu(T)}{T_\nu(T)},$$

(8)

where $T_\nu^0(T)$ is the standard dependence on the electromagnetic temperature $T$. Typically $\delta$ assumes values of order 0.01 and is always smaller than about 0.1. The neutrino temperature enters the weak rates through Fermi factors of the kind

$$\frac{1}{1 + \exp(az_\nu)}, \quad \frac{1}{1 + \exp(az_\nu)}.$$
where \( z_{\nu} = m_{\nu}/T_{\nu} \). Therefore, the additional terms for the rates can be obtained by integrating the factors

\[
\frac{1}{1 + \exp[a z_{\nu}^2 (1 + \delta)]} - \frac{1}{1 + \exp(a z_{\nu}^2)} = \sum_{n=1}^{\infty} f_n(a z_{\nu}^2) \delta^n,
\]

where in our numerical treatment the series was truncated to the third term. The corrections and the standard rates in the limit \( \delta \to 0 \) one recovers the \( n \to p \) rates including finite mass, QED radiative and thermodynamic corrections, while disregarding modifications of these subleading effects in the \( \delta \neq 0 \) corrections. We finish with six functions \( \epsilon_i(T) \)

\[
\frac{\Delta \Gamma_{n \to p}}{\Gamma_{n \to p}} = \epsilon_n(T) \delta + \epsilon_{n2}(T) \delta^2 + \epsilon_{n3}(T) \delta^3,
\]

\[
\frac{\Delta \Gamma_{p \to n}}{\Gamma_{p \to n}} = \epsilon_p(T) \delta + \epsilon_{p2}(T) \delta^2 + \epsilon_{p3}(T) \delta^3,
\]

fitting the change in the weak rates with an accuracy better than 5%.

In principle, one has a single covariant energy conservation equation for all components of the primordial plasma. For the sake of simplicity, however, in the previous considerations the “two fluids entropy conservation” was used to obtain the \( T_o(T) \) relation. We can now derive the evolution of the thermodynamical quantities by applying the covariant energy conservation law to one of the two plasmas, e.g. the electromagnetic one, so that the first Friedmann equation is

\[
\frac{dT}{dt} = -3H (\rho_{em} + P_{em}) \left( \frac{d\rho_{em}}{dT} \right)^{-1},
\]

where \( H \) depends on \( \rho_{tot} \) through the second Friedmann equation. If the additional species \( X \) couples to the electromagnetic fluid, the \( T \)-relation is further affected by a modified \( (\rho_{em} + P_{em}) \) factor, at least until the scattering freeze-out is reached. This has been roughly estimated to happen at \( T \sim 35 \) keV (Appendix B). In the numerical code, the \( X \)-particles were considered to decouple instantaneously from the electromagnetic component of the plasma for \( T \leq 35 \) keV. Relaxing this assumption our results remain essentially unchanged because for such low temperatures nucleosynthesis has almost completely stopped and the \( X \)-particles have a negligible impact, at least for the interesting mass range.

Another input parameter for the BBN calculation is the radiation density contributed by ordinary neutrinos which we fix to \( N_{\text{eff}} = 3 \).

Finally, we need the cosmic baryon density. The best-fit value from the temperature fluctuations of the cosmic microwave radiation as measured by the WMAP satellite is \( \Omega_{B}h^2 = 0.024 \pm 0.001 \) [17], including large-scale structure data from the 2dF galaxy redshift survey shifts this result to \( \Omega_{B}h^2 = 0.023 \pm 0.001 \), and including Lyman-\( \alpha \)

data further shifts it to \( 0.0226 \pm 0.0008 \) [17]. In our study we always use a fixed value of

\[
\Omega_{B}h^2 = 0.023.
\]  

We have checked that for 2\( \sigma \) variations of \( \Omega_{B}h^2 \) our conclusions remain essentially unchanged. Note that in our BBN code the final value of \( \eta \equiv n_{B}/n_{\gamma} \) or of \( \Omega_{B}h^2 \) is used to work out the one that enters the initial condition of the problem. We have

\[
\eta = \frac{n_B}{n_{\gamma}} = \frac{n_B^{12} n_B^{13} n_B^{14} n_B^{15}}{n_{\gamma} (a T_f)^3},
\]

so that entropy conservation implies the well-known factor 11/4 in the standard case. In general it is a factor depending on the \( X \)-particle properties and was numerically evaluated.

B. Neutrino-Coupled Particles

As a first generic type of \( X \)-particles we consider particles that annihilate predominantly into neutrinos \( XX \leftrightarrow \nu \bar{\nu} \). We explicitly study three cases, that of Majorana fermions with a total of \( g_X = 2 \) inner degrees of freedom (case F2), self-conjugate scalar bosons with \( g_X = 1 \) (case B1), and scalar bosons with a particle and anti-particle degree of freedom \( g_X = 2 \), case B2. With the ingredients discussed in the previous section we calculated the abundances for the light elements \( ^2\text{H}, ^3\text{He}, \) and \( ^7\text{Li} \) shown in Fig. 4 as a function of the new particle mass \( m_{X} \). For \( m_{X} \gtrsim 20 \) MeV we recover the standard BBN predictions. For very small masses \( m_{X} \to 0 \) these particles freeze out relativistically and their effect on BBN is exactly that of \( \Delta N_{\text{eff}} = 4/7, 1 \), or 8/7 additional relativistic neutrinos for the three cases B1, F2 and B2, as summarized in Table II.

The BBN effect of the new particles is dominated by their contribution to the primordial energy density and thus similar to an additional neutrino species. However, it is worthwhile to note the extra effect for intermediate \( m_{X} \) relative to the asymptotic case \( m_{X} \to 0 \). It is caused by the modified \( T_o/T \) ratio previously described.

Our results essentially agree, both qualitatively and quantitatively, with those of Ref. [13], except for lithium. The difference is explained by our value of \( \Omega_{B}h^2 \) where the dependence of \( ^7\text{Li} \) on \( N_{\text{eff}} \) is opposite from the situation in Ref. [13]. In our case this nuclide is essentially produced through the channel \( ^4\text{He}(^3\text{He}, \gamma)^7\text{Be} (e^-, \nu_e)^7\text{Li} \)

| Case     | \( g_X \) | \( \Delta N_{\text{eff}} \) | \( \eta / \eta_i \) |
|----------|-----------|-----------------|-----------------|
| B1       | 1         | 4/7             | 3.25            |
| B2       | 2         | 8/7             | 3.75            |
| F2       | 2         | 1               | 3.65            |
FIG. 1: Calculated light-element abundances for $^4$He (top), $^2$H (middle), and $^7$Li (bottom) for neutrino-coupled new particles. The indicated cases B1, B2, and F2 are described in Table I. The horizontal lines indicate the effect of $N_{\text{eff}} = 3$, 3.5, and 4 relativistic neutrinos.

while at the lower value of $\Omega_B h^2$ used in Ref. [13] the direct channel $^4$He($^3$H, $\gamma$)$^7$Li dominates.

Our theoretical predictions can be compared with the measured primordial abundances summarized in Table II. For helium, the standard BBN prediction significantly exceeds the most recent measured value [19], and this discrepancy is even worse for other helium determinations that are lower (for a review see, e.g. Ref. [18]). Therefore, a meaningful comparison of our non-standard BBN prediction for helium with observations is difficult. We show these results primarily for the purpose of illustration.

A clear interpretation of the Spite plateau in the lithium data from metal-poor halo stars is still lacking so that it is not clear how to compare the lithium observations (average value given in Table II) with the theoretical prediction that is roughly a factor 2–3 larger. Both standard and non-standard physics explanations of this discrepancy have been invoked, e.g. Ref. [21] and references. Therefore, while different observations of the primordial $^4$He abundance disagree on its exact value, a tension with the BBN prediction always exists. Our new particles exacerbate this discrepancy for $m_X < \sim 10$ MeV and are thus disfavored or even excluded.

The deuterium abundance extracted from the QSO systems agrees perfectly with the BBN prediction, although it could be affected by possibly underestimated systematic errors. In any event, the new particles change the prediction only within the 1$\sigma$ observational range so that deuterium adds little new information on the viability of the new particles.

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C. Electromagnetic Couplings

We next turn to new particles that do not interact with neutrinos but remain in perfect thermal equilibrium with the electromagnetic plasma throughout the BBN epoch by virtue of processes such as $X \bar{X} \leftrightarrow e^+e^−$ and $X + e \leftrightarrow X + e$. The BBN impact of these particles is more subtle and can be opposite to the neutrino-coupled case. Moreover, this is the case relevant as a source for galactic positrons [1, 2, 3, 4, 5]. For particles of this sort, there is a tension between the required primordial annihilation cross section to achieve the correct dark-matter density and the one in the galaxy to avoid overproducing positrons. The solution preferred in Refs. [1, 2, 3, 4, 5] is that of a predominant p-wave annihilation channel that will suppress the annihilation rate in the galaxy relative to that in the early universe. In particular, a specific model for a new boson was constructed where particles and anti-particles are not identical, i.e. our case B2. Of course, the detailed form of the cross section is not relevant for our work because we assume that the new particles are in perfect equilibrium with the electromagnetic plasma.

The calculated light-element abundances as functions of $m_X$ are shown in Fig. 2. For $m_X \lesssim 2$ MeV the abundances are all shifted away from the observed values. However, for $m_X \gtrsim 2$ MeV the “entropy effect” works in the direction of lowering $Y_p$, by up to $\Delta Y_p \approx -0.002$ for the B2 case, without significantly affecting the $^2$H and $^7$Li predictions. Therefore, the concordance with the $^4$He observations is improved. This point is illustrated more directly by Fig. 3 where we compare the new $^4$He predictions for the cases B1 and B2 with the standard $N_{\text{eff}} = 3$
FIG. 2: Calculated light-element abundances for $^4$He (top), $^2$H (middle), and $^7$Li (bottom) for electromagnetically coupled new particles. The indicated cases B1, B2, and F2 are described in Table I. The horizontal dotted lines indicate the standard predictions for the indicated values of the effective number of relativistic neutrino species, $N_{\text{eff}}$. The $1\sigma$ error band is shown for which the $1\sigma$ error band is shown. $Y_p$ would benefit from this effect up to $m_X \lesssim 15$ MeV, even if the values $m_X \lesssim 10$ MeV are preferred. Our results again agree qualitatively with those of Ref. \cite{13}.

For low masses, the value of $\eta_i$ needed to match the WMAP finding for $\eta_i$ is significantly increased relative to the standard factor 2.75, as shown in the last column of Table I. This explains physically the huge decrease (increase) in the $^2$H ($^7$Li) yield that does not strongly depend on $N_{\text{eff}}$. Quite on the contrary, as $Y_p$ depends only logarithmically on $\eta_i$ its change is essentially dominated by the addition of extra degrees of freedom to the electromagnetic plasma. The $X$-particles are now hotter than in the neutrino case and thus provoke a bigger effect. For $m_X \gtrsim 20$ MeV one recovers the standard predictions.

D. Particles Coupled both to the Electromagnetic Plasma and to Neutrinos

It may also be that the new particles interact strongly enough with both charged leptons and neutrinos to keep the two fluids in equilibrium beyond the usual decoupling epoch, a situation that was not previously treated. In this case $T_\gamma = T_\nu$ is maintained longer and perhaps throughout the nucleosynthesis epoch, depending on the new particle properties. This effect would be present even for very high $m_X$ if one introduces an additional direct coupling of the neutrinos to the charged leptons through an extra gauge boson as in the preferred $U$-boson model discussed in Refs. \cite{1, 2, 3, 4, 5}. Of course, in this case laboratory data, e.g. from $\nu$-$e$ scattering \cite{22, 23}, provide strong limits so that this situation may be rather unphysical.

In Fig. 4 we show the light-element abundances as a function of the assumed neutrino decoupling temperature $T_D$. If it is lowered to values $T_D \lesssim m_e$, the resulting modification of the light-element abundances is quite extreme and strictly excluded. For $^2$H and $^7$Li the “eta-effect” dominates, while for $^4$He the change in $T_\nu/T$ and that of

We have always neglected the role of dark matter residual annihilations during the freeze-out epoch. A detailed treatment of such effects is beyond the scope of our paper. However, an approximate study of this phenomenon (Appendix \cite{14}) indicates that this late-time entropy generation effect is sub-dominant. Its effect goes in the direction of a further marginal reduction of $Y_p$ and a small increase in the deuterium yield.

We have also neglected the possible photo-dissociation of $^2$H and $^7$Li induced by late dark matter annihilations. As shown in Appendix \cite{14} this effect is marginal unless, perhaps, if the dark matter particles couple directly with photons.

FIG. 3: $^4$He abundance as in Fig. 2 here on a linear scale for $m_X$. The horizontal dotted line indicates the standard prediction for $N_{\text{eff}} = 3$. The gray band is the $1\sigma$ observational range for $Y_p$ according to Ref. \cite{19}.
\[ \tau(E) + d[E / n(E)] dE = 0, \] (15)

\[ n(E) = b(E) \int_{E}^{\infty} ds q(s) \exp \left( - \int_{E}^{s} \frac{dy}{\tau(y) b(y)} \right). \] (16)

The expected positron flux at the top of the atmosphere, neglecting solar modulation effects, is 

\[ j = n(E) / 4\pi, \]

where we have assumed relativistic velocities. In the relevant energy range and for an essentially neutral environment, ionization and bremsstrahlung are the dominant energy-loss mechanisms so that we have

\[ b(E) \approx 36.18 m_e \sigma_T n_H \left[ 1 + 0.146 \ln \left( \frac{E}{m_e} \right) + \frac{E}{709 m_e} \right] = 3.69 \times 10^{-13} \text{ MeV s}^{-1} n_H \]. (17)

Here, \( \sigma_T \) is the Thomson cross section and \( n_H \) the hydrogen density that dominates the interstellar medium. For the relevant energies, \( b(E) \) is a slowly varying function.

The latter can be written in the form

\[ \tau(E) = \frac{1.33 \times 10^{14} \text{ s}}{f(E) n_H} \], (18)

where \( f(E) = 0.02 - 1 \) in the energy range of interest.

Finally we need the positron injection spectrum from dark-matter annihilation,

\[ q(E) = n_X^\delta(\sigma_a v)(1 - E / m_X), \] (19)

III. LOW-ENERGY COSMIC RAY POSITRONS

The proposed MeV-mass dark matter particles discussed in Refs. 1, 2, 3, 4, 5 are supposed to annihilate in the galactic bulge and produce a flux of low-energy cosmic-ray positrons that can explain the observed 511 keV \( \gamma \)-ray signature. For such \( X \)-particles we also expect a flux of low-energy positrons at Earth from local dark-matter annihilation. Such a signature was already proposed for the detection of the traditional GeV-TeV mass range of dark matter particle candidates 24, 25. For the case of annihilating MeV-mass dark matter particles we expect a much larger positron flux and therefore we study if additional constraints arise from the low-energy cosmic-ray positrons in the solar neighborhood.

To this end we derive the expected positron flux from \( X \)-particle annihilation. For a stationary situation the continuity equation for the cosmic-ray positrons is

\[ q(E) - \frac{n(E)}{\tau(E)} + \frac{d[n(E) b(E)]}{dE} = 0, \] (15)

where \( n(E) \) is the differential positron density, \( q(E) \) the injection spectrum, \( \tau(E) \) an effective containment time, and \( b(E) \equiv -dE/dt \) the energy-loss function. This equation is solved by

\[ n(E) = b(E) \int_{E}^{\infty} ds q(s) \exp \left( - \int_{E}^{s} \frac{dy}{\tau(y) b(y)} \right). \] (16)

\[ b(E) \approx 36.18 m_e \sigma_T n_H \left[ 1 + 0.146 \ln \left( \frac{E}{m_e} \right) + \frac{E}{709 m_e} \right] \]

Here, \( \sigma_T \) is the Thomson cross section and \( n_H \) the hydrogen density that dominates the interstellar medium. For the relevant energies, \( b(E) \) is a slowly varying function.

It is easy to estimate that low-energy positrons do not travel far before annihilating so that the containment time \( \tau(E) \) is identical with the annihilation time scale. The latter can be written in the form

\[ \tau(E) = \frac{1.33 \times 10^{14} \text{ s}}{f(E) n_H}, \] (18)

where \( f(E) = 0.02 - 1 \) in the energy range of interest.

Finally we need the positron injection spectrum from dark-matter annihilation,

\[ q(E) = n_X^\delta(\sigma_a v)(1 - E / m_X), \] (19)
where we have assumed that the $X$-particles are different from their antiparticles and that $n_X = n_{\bar{X}}$.

Equation (10) cannot be expressed in closed form, but for our purposes an analytical approximation is accurate enough. The factor $b \tau$ does not depend on $n_H$ and is found to be

$$b \tau = 50–5000 \text{ MeV},$$

(20)

where the range reflects the monotonic energy dependence. Ignoring this energy-dependence allows us to write the solution of Eq. (10) for $E \leq m_X$ as

$$n(E) \approx \frac{n_X^2 \langle \sigma_a v \rangle}{b(E)} \exp \left( -\frac{m_X - E}{b \tau} \right).$$

(21)

For the interesting range of $m_X$ and $E$ the exponential factor is always close to 1. Physically this represents the fact that positrons produced at energy $E = m_X$ will be lost from this “energy bin” primarily by down-scattering, not by annihilation.

In order to predict the positron flux we use $n_H = 1 \text{ cm}^{-3}$. For the local dark matter density we use the canonical value $\rho_{DM} = 300 \text{ MeV cm}^{-3} = m_X (n_X + n_{\bar{X}})$. For the annihilation cross section we first consider an $s$-wave model with $\langle \sigma_a v \rangle = \sigma_0$ where $\sigma_0$ is fixed by the early-universe freeze-out calculation (Appendix A). We compare the flux prediction with the best 95% CL upper limits in the 20–90 MeV range that are given in Fig. 4 of Ref. 28. At 20 MeV, the flux limit is approximately $1.2 \times 10^{-5} \text{ cm}^{-2} \text{s}^{-1} \text{ sr}^{-1} \text{ MeV}^{-1}$, for 30 MeV it is $8.5 \times 10^{-6}$ in these units, for 50 MeV it is $2.5 \times 10^{-6}$, and for 90 MeV it is $9 \times 10^{-7}$.

We find flux predictions exceeding these limits by factors $\gtrsim 500$. This implies that the annihilation cross section required in the early universe for reducing the particle density to the dark-matter level produces an unacceptably large local positron flux and is thus excluded for $m_X > 20 \text{ MeV}$. For smaller $m_X$ we have no constraint. For a $p$-wave channel we do not obtain any limits because $\langle \sigma_a v^2 \rangle$ is strongly suppressed in the galaxy because of the small dark-matter velocities $v$ of order $10^{-3}$.

Therefore, we confirm the conclusion of Refs. 1, 2, 4, 5 that an $s$-wave annihilation channel $X \bar{X} \to e^+ e^-$ is not acceptable for thermal relics. However, our argument does not depend on the uncertain dark-matter profile of the galactic bulge.

**IV. ENERGY TRANSFER IN STARS**

New MeV-mass particles could play an important role in stars. For example, they would be thermally produced in the collapsed core of a supernova and could contribute to the energy loss or the transfer of energy in these systems 29. However, the $X$-particles discussed here have stronger-than-weak interactions, implying that in a supernova these effects would be small compared to those of ordinary neutrinos. Therefore, even though the new particles would be thermally excited in a supernova core, there are no obvious observational consequences.

In ordinary stars, and especially in our Sun, dark-matter particles will be trapped and contribute to the transfer of energy in potentially observable ways. 20. In the following we investigate if MeV-mass particles could be relevant in this context. The result is that for MeV-range masses the evaporation time is very short so that the steady-state abundance of $X$-particles in the Sun is too small to be important.

A simple estimate of the energy conduction by the new particles can be worked out in a one-zone model of the Sun 20. One assumes that the dark-matter stationary distribution in the Sun is globally Maxwellian at a uniform temperature $T_X$, with its maximum density found at a scale radius $r_X$. Assuming that $T_X$ is identical with the temperature at the solar center, and taking $r_X$ to be of order the solar radius, the luminosity carried by the new particles is of order

$$L_X \sim \frac{10^{-22} N_X}{\langle \sigma_s / \text{pb} \rangle \sqrt{m_X / \text{MeV}}},$$

(22)

where $N_X$ is the total number of $X$-particles trapped in the Sun and $\sigma_s$ is the scattering cross section on electrons, taken to be comparable to the annihilation cross section as described in Appendix A. For particles even more weakly interacting, i.e. for $\sigma_s \ll \text{pb}$, one enters the Knudsen regime where the effect of energy transfer is much smaller 51, 52.

The steady-state number of dark-matter particles collected by the Sun arises from an equilibrium between capture and evaporation, i.e. $N_X = A/P_e$ with $A$ the number of particles captured per unit time and $P_e$ the escape probability per unit time. We estimate the capture rate to be 33, 34

$$A \sim 4.34 \frac{n_X G_N M_\odot R_\odot}{v_{\text{gal}}},$$

(23)

where $v_{\text{gal}} \approx 300 \text{ km s}^{-1}$ is the mean square velocity of the galactic dark matter near the orbit of the Sun and $n_X$ its number density.

The escape probability $P_e$ is physically the ratio of the fraction of particles in the “escape-tail” of their distribution to the typical time needed to repopulate it. A simple estimate is 30

$$P_e \sim \left( \frac{v_f \delta v}{G_N M_\odot} \right)^{3/2} \sqrt{\frac{m_X v_f^2}{2T_X}} \exp \left( -\frac{m_X v_f^2}{2T_X} \right),$$

(24)

where $v_f$ is the escape velocity from a typical position in the Sun. Near the solar center one has $v_f \sim (15.8 T_X/m_p)^{1/2}$ 34. Further, $\delta v \sim v_s m_s/m_X$ where $m_s$ and $v_s$ refer to the scattered particles, electrons in our case, for which we assume a thermal velocity distribution so that $v_s = (3 T_s/m_s)^{1/2}$ and we assume $T_s = T_X$.

Based on these simple estimates and using a simplified solar model we find $N_X \sim 10^{38} (\text{pb}/\sigma_s)$. Comparing this
result with Eq. [22] we conclude that the effect of the new particles is too small to be significant for the Sun. Of course, our estimate is rather crude considering that the escape lifetime of the dark matter particles is comparable to their orbital period.

V. SUMMARY AND CONCLUSIONS

We have analyzed some astrophysical and cosmological consequences of the intriguing possibility that the cosmic dark matter consists of MeV-mass particles [1, 2, 3, 4, 5]. These particles are assumed to be thermal relics and thus have interaction cross sections that are larger than weak.

Such particles do not have any apparent consequences for stellar evolution. In supernovae, their interaction is too strong so that neutrinos continue to play the leading role for energy loss and energy transfer. In the Sun, the particle mass is so small that evaporation prevents the trapped steady-state population from growing large enough for a significant contribution to energy conduc-
tion.

We have derived new constraints coming from the low-energy positron component of cosmic rays. An s-wave annihilation cross section for $XX \rightarrow e^+e^-$ as large as implied by the early-universe freeze-out calculation causes an excessive positron flux for $m_X = 20-90$ MeV where experimental upper limits are available. Therefore, only p-wave annihilation is compatible with these constraints. These conclusions agree with those reached in Refs. [1, 2, 3, 4, 5], but in our case they do not depend on the assumed dark-matter profile of the galactic bulge.

The main subject of our work, however, was the impact of MeV-mass particles on big-bang nucleosynthesis (BBN). Significant modifications arise only for $m_X \lesssim 20$ MeV. The effects found are largely independent of the energy dependence of the annihilation cross section and even of its exact value, provided that it is high enough to take the new particles to equilibrium with the neutrinos or the electromagnetic plasma.

For the neutrino-coupled case, the impact of the new particles is comparable to that of additional neutrino species. Notably, the primordial helium abundance is increased, exacerbating the discrepancy between predictions and observations. Therefore, such $X$-particles are disfavored by BBN for masses up to about 10 MeV.

For the electromagnetically coupled case, the BBN concordance would be severely disturbed for $m_X \lesssim 2$ MeV. However, there is a region $m_X = 4–10$ MeV where the primordial helium mass fraction $Y_p$ is actually reduced relative to the standard case while the predicted $^2\text{H}$ remains compatible with observations. This non-trivial phenomenon is a consequence of the “entropy effect” discussed in the paper. While this effect was already found in Ref. [13], it now assumes greater importance because it slightly improves the discrepancy between the BBN predictions for $Y_p$ and its observed value.

In summary, the MeV-mass dark matter particles proposed in Refs. [1, 2, 3, 4, 5] as a source for positrons in the galactic bulge are not incompatible with BBN, provided the particle mass exceeds a few MeV. On the contrary, in the mass range $m_X = 4–10$ MeV these particles slightly improve the concordance between BBN calculations and the observed helium abundance. It is quite fascinating that such exotic particles, far from being excluded, seem to have several beneficial consequences in astrophysics and cosmology.

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APPENDIX A: RELIC DENSITY AND ANNIHILATION CROSS SECTION

Once the mass $m_X$ and the spin multiplicity $g_X$ of a thermal relic are fixed, the requirement that it is the dark matter uniquely determines the annihilation freeze-out temperature $T_F$ and the annihilation reaction rate at this temperature $\left< \sigma v \right>_T$. The relic density is

$$\Omega_X h^2 = \frac{8\pi^2 m_X n_X}{3 m_{Pl}^2 H^2} = \frac{8\pi m_X s}{3(m_{Pl}H/h)^2} \left( \frac{n_X}{s} \right),$$

where $H/h = 100$ km s$^{-1}$ Mpc$^{-1}$, $m_{Pl} = 1.22 \times 10^{22}$ MeV is the Planck mass, and $s$ is the entropy density. Note that for particles that are not self-conjugate $\Omega_X$ only includes the particles so that $\Omega_{\text{DM}} = \Omega_X + \Omega_{\bar{X}}$.

Neglecting entropy-producing phenomena, the entropy per comoving volume remains constant so that

$$\left( \frac{n_X}{s} \right)_0 = \left( \frac{n_X}{s} \right)_T$$

for $T \leq T_F$. If one assumes a functional form for $\left< \sigma v \right>_T$, an order-of-magnitude estimate of $T_F$ is given by the condition $n_X\left< \sigma v \right>_T = H(T_F)$. A more accurate formula is obtained by a semi-analytical treatment of the Boltzmann equation [14]:

$$x_F = x_0 - (n + \frac{1}{2}) \log x_0,$$

where

$$x_0 = \log \left[ 0.038 (n + 1) \frac{g_X}{\sqrt{g_{_{\text{eff}}}}} m_{Pl} m_X \sigma_n \right].$$

Here, the parameterization

$$\left< \sigma v \right>_T = \sigma_n x^{-\eta}$$

is obtained by a semi-analytical treatment of the Boltzmann equation [14].
with \( x = m_X/T \) was assumed. The most interesting cases are obtained for \( n = 0 \) (s-wave annihilation) and \( n = 1 \) (p-wave). The previous equations yield

\[
\left( \frac{n_X}{s} \right)_0 = \frac{3.79 \sqrt{g_F} (n+1) \sigma^{n+1}_F}{g_{sF} \eta^{(n+1)/2} \sqrt{g_F}},
\]

\[
\Omega_X h^2 = 0.032 \frac{(n+1) \sigma^{n+1}_F \eta^{(n+1)/2} \sqrt{g_F}}{g_{sF} (\sigma_0 / pb)}, \tag{A6}
\]

with an accuracy better than 5\%. Since \( \Omega_{DM} h^2 \approx 0.11 \) \cite{17}, and \( x_F = 15-20 \), it is easily seen that, for \( X \neq \bar{X} \), one has \( \sigma_0 \sim 5 \) pb and \( \sigma_1 \sim 10^2 \) pb.

**APPENDIX B: KINETIC FREEZE-OUT**

We estimate the kinetic freeze-out temperature for \( X \)-particles coupled to the electromagnetic plasma. We make the rough approximation that the annihilation and scattering cross sections are comparable, \( \sigma_a \sim \sigma_s \), so that the ratio between the annihilation and scattering rates is essentially given by \( \Gamma_a/\Gamma_s \sim n_X/n_s \). For our case the scattering targets are electrons and positrons so that

\[
n_s = n_{e^+} + n_{e^-} = 4 \left( \frac{m_e^2}{2 \pi \bar{z}} \right)^{3/2} e^{-\bar{z}} \cosh \xi_e, \tag{B1}
\]

where \( \xi_e \equiv m_e/T \) is the electron degeneracy parameter and \( \bar{z} \equiv m_e/T \). Therefore, as long as the \( X \)-particles are in kinetic equilibrium we have

\[
\frac{\Gamma_a}{\Gamma_s} \approx \left( \frac{m_X}{m_e} \right)^{3/2} \exp \left( -\frac{m_X - m_e}{T} \right) \frac{1}{2 \cosh \xi_e}, \tag{B2}
\]

This result applies at \( T_F \) only if \( \left( \Gamma_a/\Gamma_s \right) T_F \ll 1 \). Assuming \( m_X = 1 \) MeV, one finds from Eq. \( \text{A3} \) that \( T_F \approx 0.07 \) MeV and from a standard BBN code that \( \xi_e(T_F) \approx 0.32 \times 10^{-7} \) so that \( (\Gamma_a/\Gamma_s)T_F \approx 4 \times 10^{-4} \) which is indeed \( \ll 1 \).

Once a value for the annihilation effective cross section \( \sigma_n \) has been determined, one can easily deduce a kinetic freeze-out temperature \( T_K \) for the decoupling of the dark matter particles from the electromagnetic plasma, assuming that \( \sigma_a \sim \sigma_s \). The condition \( \Gamma_s(T_K) = H(T_K) \) implies

\[
\langle \sigma_a v \rangle_{T_K} = \frac{4.44 T_K^2}{m_{p1} n_s(T_K) \zeta(3)} \sqrt{g_F(T_K)}/10.75, \tag{B3}
\]

With \( \cosh \xi_e \sim 1 \) and \( g_s \sim 3.36 \) one obtains

\[
z_K - (0.5 - \ell) \log z_K \approx 14.1 + \log(\sigma_s / \ell / \text{pb}), \tag{B4}
\]

where \( \langle \sigma_a v \rangle \equiv \sigma_s / \ell \). For typical values of \( \sigma_s, n \) this gives \( T_K \approx 35 \) keV.

**APPENDIX C: ENTROPY GENERATION DURING DECOUPLING**

We sketch an argument that indicates that the entropy production associated with the dark-matter freeze-out is a sub-leading effect. To this end we assume that every \( e^+ e^- \) annihilation product is instantaneously thermalized and converted into photons. In this case the evolution equation for the \( X \) number density is \( \text{[14]} \)

\[
\dot{n}_X + 3H n_X = -C[n_X], \tag{C1}
\]

where

\[
C[n_X] = \langle \sigma_a v \rangle [n_X^2 - n_{X,eq}^2]. \tag{C2}
\]

The same equation applies to \( n_X = n_{\bar{X}} \) because we always assume equal distributions for both \( X \) and \( \bar{X} \). The subsequent \( e^+ e^- \) annihilations imply

\[
\dot{n}_\gamma + 3H n_\gamma = +2C[n_X]. \tag{C3}
\]

The equilibrium energy densities are

\[
\rho_\gamma = \left( \frac{\pi^2}{15} \right) \left( \frac{\pi^2}{2 \zeta(3)} \right) n_\gamma^{4/3},
\rho_X = n_X \left( m_X + \frac{3}{2} T \right), \tag{C4}
\]

where the non-relativistic regime was used for the dark-matter particles. The energy injection by late annihilations is estimated as

\[
\dot{\rho}_\gamma = \frac{8}{3} \left( \frac{\pi^4 T}{30 \zeta(3)} \right) C[n_X],
\dot{\rho}_X = -m_X \left( 1 + \frac{3T}{2m_X} \right) C[n_X]. \tag{C5}
\]

The direct change of the total energy density due to this conversion of dark-matter particles into photons is a very small effect, as one can see by evaluating the ratios \( \dot{\rho}_\gamma / \rho_\gamma \) and \( \dot{n}_\gamma / n_\gamma \).

The main non-negligible consequence is on the \( X \)-particles. The energy injection by late annihilations is estimated as

\[
\frac{d \rho_{\text{em}}}{dt} = -3H (\rho + P)_{\text{em}}
- \left[ m_X - \left( \frac{4\pi^4}{45\zeta(3)} - 3 \right) T \right] C[n_X]. \tag{C6}
\]

The problem is then restricted to the solution of the equation for \( n_X \) in order to calculate the relevant quantity \( C[n_X] \). This was done by standard techniques after some standard substitutions (see e.g. Ref. \[14\]). We found a negligible effect on \( Y_{\text{pr}} \) and a change of the \( ^3\text{H} \) and \( ^7\text{Li} \) yields of order 0.1 in the units of Fig. \[2\]. Therefore, we are indeed dealing with a sub-leading effect.
We show that the dissociation of fragile nuclides, principally deuterium, by late dark-matter annihilation is completely negligible. To this end we first follow standard works on this subject (e.g. Ref. [35]) and note that the “first-generation” of nonthermal photons is rapidly degraded to have a spectrum with a high energy cut $E_C \approx m_X^2/22T$ because of the highly efficient pair creation and subsequent reactions on the background medium. This means that in order for some photons to survive and to be able to dissociate nuclei the universe has to drop at least to values $T \lesssim m_X^2/22B_D \approx 5 \text{ keV}$, where $B_D \approx 2.2 \text{ MeV}$ is the binding energy of deuterium. This implies that neglecting this phenomenon during the BBN epoch is certainly a good approximation. Note also that in our preferred models, MeV-mass $X$ particles would not annihilate directly into photons, so the “first-generation” photon spectrum is already degraded in energy because it is produced by secondary effects of the primary electrons and positrons.

Therefore, by simply looking at the expected energy spectrum, the most dangerous process would be the electro-disintegration of deuterium, for which a further 0.5 MeV penalty in the energy should be considered as the electron rest mass would not be released. Without invoking a detailed analysis we conclude that our results up to $m_X \lesssim 2.7 \text{ MeV}$ remain unchanged.

For higher values of $m_X$ considered in our BBN analysis, say 3–20 MeV, we reach the same conclusion by following the treatment of the late-time annihilations of a relic particle described in Ref. [34]. Assuming that electro-dissociation of deuterium is the only relevant phenomenon, the Deplination factor can be written as

$$\exp \left[ - \int_{t_i}^{t_f} dt \Gamma_{eD}(t) \right], \quad (D1)$$

where $t_i$ is the time at the onset of dissociation, $t_f$ is some late time where the D abundance is observed, and

$$\Gamma_{eD} = s \int_{B_D}^{E_C(T)} dE f_c(E, T) \sigma_{eD}(E). \quad (D2)$$

Here, $\sigma_{eD}(E) \approx 13.3 \mu b \left( E - B_D \right)/\text{MeV}^{1/2}$ is the electro-dissociation cross-section of deuterium [35]. Further, $f_c(E, T)$ is the “steady state” dark-matter annihilation product spectrum. By arguments similar to the ones presented in Ref. [36] it is written as

$$f_c(E, T) \approx \sigma_n \left( \frac{T}{m_X} \right)^n \frac{s}{n_s \sigma_{e\gamma}(E)} \left( \frac{n_X}{s} \right)^2 \times \frac{m_X \theta[E_{\text{max}}(T) - E]}{\sqrt{E^3 E_C(T)}}, \quad (D3)$$

where $\sigma_{e\gamma}(E)$ is the Compton scattering cross section and $E_{\text{max}}(T) = \text{Min}[E_C(T), m_X]$. Note that the primary particle’s “first” shower spectrum has roughly the same shape for $\gamma$ or $e$ initiated showers [35 36].

We have numerically solved the integral in Eq. (D1) by changing to the temperature variable and assuming a radiation dominated universe. The effects we found are completely negligible both for $s$- and $p$-wave annihilations.

Note, however, that for showers induced by secondary photons one should replace in Eq. (D3) the huge factor $n_s$, with $n_c$, and the electro-dissociation cross section with the photo-dissociation one, while a penalty factor in the energy spectrum would enter. Making some extreme assumption one could thus gain up to a factor $10^{12}$ of the previous estimates. Even so, it is not enough to affect by more than a few percent our simplified nuclides’ predictions. The effects are even less pronounced for the preferred case $n = 1$ (p-wave annihilation).

\[\text{APPENDIX D: DEUTERIUM PHOTO-DISSOCIATION}\]

We show that the dissociation of fragile nuclides, principally deuterium, by late dark-matter annihilation is completely negligible. To this end we first follow standard works on this subject (e.g. Ref. [35]) and note that the “first-generation” of nonthermal photons is rapidly degraded to have a spectrum with a high energy cut $E_C \approx m_X^2/22T$ because of the highly efficient pair creation and subsequent reactions on the background medium. This means that in order for some photons to survive and to be able to dissociate nuclei the universe has to drop at least to values $T \lesssim m_X^2/22B_D \approx 5 \text{ keV}$, where $B_D \approx 2.2 \text{ MeV}$ is the binding energy of deuterium. This implies that neglecting this phenomenon during the BBN epoch is certainly a good approximation. Note also that in our preferred models, MeV-mass $X$ particles would not annihilate directly into photons, so the “first-generation” photon spectrum is already degraded in energy because it is produced by secondary effects of the primary electrons and positrons.

Therefore, by simply looking at the expected energy spectrum, the most dangerous process would be the electro-disintegration of deuterium, for which a further 0.5 MeV penalty in the energy should be considered as the electron rest mass would not be released. Without invoking a detailed analysis we conclude that our results up to $m_X \lesssim 2.7 \text{ MeV}$ remain unchanged.

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$$\exp \left[ - \int_{t_i}^{t_f} dt \Gamma_{eD}(t) \right], \quad (D1)$$

where $t_i$ is the time at the onset of dissociation, $t_f$ is some late time where the D abundance is observed, and

$$\Gamma_{eD} = s \int_{B_D}^{E_C(T)} dE f_c(E, T) \sigma_{eD}(E). \quad (D2)$$

Here, $\sigma_{eD}(E) \approx 13.3 \mu b \left( E - B_D \right)/\text{MeV}^{1/2}$ is the electro-dissociation cross-section of deuterium [35]. Further, $f_c(E, T)$ is the “steady state” dark-matter annihilation product spectrum. By arguments similar to the ones presented in Ref. [36] it is written as

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