How to produce antihelium from dark matter

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We show how to produce antihelium and other antinuclei in large fractions from the decays of a new particle \( \phi \) that carries baryon number. Close to threshold, the production of nuclear bound states is preferred over the decay into individual nucleons, effectively decoupling antihelium and antiproton fluxes and allowing the former to dominate. \( \phi \) can either form dark matter itself or be produced by it, and can give rise to a flux of antihelium that is large enough to explain the preliminary antihelium events reported by AMS-02.

I. INTRODUCTION

Most of our universe is composed of matter, with only a tiny fraction of antimatter produced in highly energetic cosmic events, measured at experiments such as PAMELA [1] and AMS-02 [2]. These experiments also probe Dark Matter (DM) models that can lead to an enhanced number of positron and antiproton events [3, 4], and even antideuteron [5] has become a promising target [4–6]. We will postulate the existence of a new particle \( \phi \) which may well be related to DM. We will therefore often take it as paradigm-shifting evidence for antimatter regions.

Future experiments such as GAPS [7] have the potential to improve these measurements significantly. Preliminary AMS-02 results report six \(^3\)He events in seven years of observations, compared to about 700 million helium events [8]. If confirmed, this observation of cosmic antimatter would be a spectacular discovery, hard to reconcile with known astrophysics and even most physics beyond the Standard Model (SM).

In this situation, the nucleons can be treated as non-relativistic and will have a significant probability of forming antinuclei with a large mass number \( A \leq k \).

We will explore this basic idea and show that we can indeed build models of new physics where a new particle \( \phi \) decays with detectable production rates for antinuclei. We will focus mostly on producing \(^4\)He but these models could easily be generalized to the production of other antinuclei.

The rest of this article is organized as follows: in Sec. [I] we discuss models that lead to \(^4\)He and the possible relation to DM. Sec. [II] gives an analogous discussion of models that produce \(^3\)He. We conclude in Sec. [IV].

II. MODELS FOR \(^3\)He PRODUCTION

As an explicit realization of \(^3\)He production let us introduce a new SM-singlet complex scalar \( \phi \) with baryon number \( -3 \) and lepton number \( -1 \) (the quantum numbers of the antitritium atom). We assume these symmetries to be sufficiently conserved so that they govern the \( \phi \) decay modes, with one of the lowest-dimensional operators given by

\[
\mathcal{L} = \frac{\phi \pi^+ \gamma_5 n \gamma_5 e}{\Lambda^3} + \text{h.c.} \tag{1}
\]

In the non-relativistic limit this corresponds to a coupling \( \phi n_p n_p e \), \( f \) being the spin up/down states of particle \( f \). Operators similar to Eq. (1) but without the \( \gamma_5 \) would create neutrons in the same spin state and thus be suppressed by Fermi statistics in the non-relativistic limit; Eq. (1) is then expected to be the dominant operator of interest for \( \phi \) decay.

A. Decay channels

We will consider the situation where \( m_\phi \) is close to the total mass of the decay products, i.e. \( m_\phi \sim 2m_n + m_p \), neglecting the electron mass \( m_e \) in the following for simplicity. In this situation, the nucleons can be treated as point particles and we can ignore the underlying quark structure. The possible final states then depend strongly on \( m_\phi \); if \( m_\phi \) is below the threshold \( 2m_n + m_p \), the decay has to go into a bound state of the nucleons in order to be energetically allowed. The relevant thresholds are given in Tab. 1 which shows that in the range \( 0.5 \text{ MeV} < m_\phi - m_{^4\text{He}} < 6.8 \text{ MeV} \), the dominant kinematically allowed
TABLE I. Lowest lying states with baryon number 3 and electric charge 1, with mass $m$ relative to $^{3}$He$^{+}$. The deuteron and triton nuclei are as usual denoted as $d^{+}$ and $t^{+}$ instead of $^{2}$H$^{+}$ and $^{3}$H$^{+}$. We will drop the ionization superscripts in the following since we always refer to nuclei.

| State | $m - m_{^{3}He}$ [MeV] |
|-------|--------------------------|
| $t^{+}$ | 0.5 |
| $d^{+}n$ | 6.8 |
| $pnn$ | 9.0 |

Decay induced by Eq. (3) is $\phi \to \pi n$, followed by the beta decay of the antitritium nucleus $t$ into $^{3}$He within $12 \text{yr}$. This is then a source for production of $^{3}$He as long as the lifetime for $\phi$ is not too long, which can indeed be the case despite the small available phase space, as we will now show.

Assuming $m_{\phi}$ to be close to $m_{^{3}He}$ ensures that the nucleons and nuclei involved in the process are non-relativistic. It is then convenient to calculate the decay rate non-relativistically with Fermi’s golden rule,

$$\Gamma(\phi \to \pi n) = 2\pi |V(\phi \to \pi n)|^2 \rho,$$

(2)

$$\rho = (m_{\phi} - m_{t})^2/(2\pi^2)$$

being the phase-space density. $V(\phi \to \pi n)$ is the appropriate transition element, here approximated as

$$|V(\phi \to \pi n)|^2 \simeq \sum_{\text{spins}} |M(\phi \to \pi \pi \pi \pi)|^2 / 2m_{\phi} 2E_{m} 2E_{n} 2E_{\pi} 2E_{\pi},$$

(3)

where the first factor is the hard process and the second factor is the overlap of the nucleon wave function with the tritium nucleus. The hard process $\phi(k) \to \pi(k_{1})\pi(k_{2})\pi(k_{3})\pi(k_{4})$ has a squared matrix element

$$\sum_{\text{spins}} |M(\phi \to \pi \pi \pi \pi)|^2 \simeq \frac{32}{A^6} (k_{1}k_{2} + m_{\phi}^2)(k_{3}k_{4}).$$

(4)

The overlap $|\langle \pi \pi \pi \pi |\rangle|$ of the nucleons with the tritium nucleus wave function can be estimated by treating the nucleons as moving in a mean field potential. Once the center of mass motion is factored out, the nucleus is described by a product of two wave functions $\psi(x_{1} - x_{2})$ for the two relative coordinates. Since the decay of $\phi$ is assumed to produce the three nucleons initially at a single point, we can estimate the overlap as

$$|\langle \pi \pi \pi \pi |\rangle^2 \approx (|\psi(0)|^2)^2.$$  

In simple shell models, the wave function $\psi(0)$ can be approximately related to the radius of $t$, $\psi(0) \sim r_{t}^{-3/2}$, or simply $|\psi(0)| \sim (100 \text{MeV})^{3/2}$ as a nuclear-physics energy scale. A more accurate calculation of this matrix element is desirable and will be left to future work.

1 The direct decay of $\phi$ into $^{3}$He is further suppressed by $G_{F}$ and will not be discussed here.

FIG. 1. Decay rates of $\phi$ to $\pi n$, $\bar{d} \pi \pi$, and $\pi \pi \pi \pi$ for $\Lambda = 10^{4} \text{TeV}$. We again denote $d \equiv ^{2}H$ and $t \equiv ^{3}H$. All nuclear matrix elements are taken to be 100 MeV to the appropriate power. The rates are calculated assuming $m_{\phi} - m_{^{3}He} \ll m_{^{3}He}$.

The nucleons are non-relativistic but the electrons relativistic, leading to the decay rate

$$\Gamma(\phi \to \pi n) \simeq \frac{2(m_{\phi} - m_{t})^2}{\pi m_{\phi} A^6} |\psi(0)|^4,$$

(5)

which we expect to be valid for $m_{\phi} - m_{t} < m_{t}$. With $|\psi(0)|^4 \simeq (100 \text{MeV})^4$ this yields the lifetime

$$\tau \simeq 10^{8} \text{Gyr} \left(\frac{10 \text{MeV}}{m_{\phi} - m_{^{3}He}}\right)^2 \left(\frac{\Lambda}{10^{4} \text{TeV}}\right)^6,$$

(6)

illustrated in Fig. 1 which can easily surpass the age of the universe, $t_{\text{universe}} \simeq 14 \text{Gyr}$.

For masses $m_{\phi} - m_{^{3}He} > 7 \text{MeV}$, the decay channel $\phi \to \bar{d} \pi \pi$ into antideuterons opens up, which we calculate in analogy to before using Fermi’s golden rule. The nuclear matrix element $|\langle \bar{d} \pi \pi |\rangle^2 \simeq |\psi(0)|^2$ is assumed to be $(100 \text{MeV})^2$ in our numerical evaluations. For slow hadrons the decay rate then reads

$$\Gamma(\phi \to \bar{d} \pi \pi) \simeq \frac{16}{105\pi^{5}m_{\phi}A^{6}(m_{d} + m_{n})^{3/2}} (m_{\phi} - m_{d} - m_{n})^{7/2} |\psi(0)|^2,$$

(7)

which starts to dominate over the two-body decay for $m_{\phi} - m_{^{3}He} \gtrsim 500 \text{MeV}$ (Fig. 1).

Finally, when $m_{n}$ is above the threshold for decay into free nucleons, $m_{\phi} > 2m_{n} + m_{p}$, but still small enough to not resolve the underlying quark structure or heavier baryons such as $\Sigma$ and $\Delta$, $\phi$ will decay into free nucleons. Using the matrix element from Eq. (4) we find the decay rate near threshold

$$\Gamma(\phi \to \pi \pi \pi \pi) \simeq \frac{137}{10468\pi^{5}A^{6}(2m_{n} + m_{p})^{5/2}} (m_{\phi} - 2m_{n} - m_{p})^{3/2} \times (m_{\phi} - m_{d} - m_{n})^{7/2} |\psi(0)|^2,$$

(8)

which is the dominant decay mode for high $m_{\phi}$. The direct decay of $\phi$ into $^{3}$He is further suppressed by $G_{F}$ and will not be discussed here.
This decay to free nucleons dominates for large \( m_\phi \) and will eventually lead to constraints from antiproton-flux measurements that plague most DM models that aim to produce heavier antinuclei \([9,10,15]\).

The main conclusion from this analysis is that \( \phi \) decays can produce almost exclusively \(^{3}\text{He} \) for \( 0 \lesssim m_\phi - m_{^{3}\text{He}} \ll 100 \text{ MeV} \) (see Fig. 1), while for \( m_\phi - m_{^{3}\text{He}} \approx 500 \text{ MeV} \) we expect similar amounts of \(^{3}\text{He} , \d, \) and \( \overline{p} \). For even larger \( \phi \) masses the \(^{3}\text{He} \) fraction will decrease and the amount of \( \overline{p} \) will become significant, eventually leading back to coalescence fractions for the antinuclei. In order to dominantly produce \(^{3}\text{He} \) we therefore need \( m_\phi \approx 3 \text{ GeV} \)\(^2\).

**B. \( \phi \) as dark matter**

Since the total lifetime of the new particle \( \phi \) can be longer than the age of the universe, it is conceivable that \( \phi \) is DM or a subcomponent of it. The appropriate relic density could be produced through interactions like the Higgs portal interaction \(|\phi|^2 H^2\) or new gauge interactions, which allow for \( \phi \) pair production without affecting its decay. Alternatively, as a result of the assigned baryon number it is reasonable to expect \( \phi \) to have an asymmetry similar to the baryon asymmetry, which is consistent with the \( \sim 3 \text{ GeV} \) DM mass \([19]\).

Assuming \( \phi \) to be DM, we can in principle calculate the antinuclei flux. As a conservative lower limit on the lifetime we take \( \tau > 10^{25} \text{ s} \) \([20]\) in order to suppress energy injection during CMB formation from the fast positron that is emitted in \( \phi \) decays. For \( m_\phi = 3 \text{ GeV} \) this corresponds to a lower limit \( \Lambda > 7 \times 10^3 \text{ TeV} \). There are about \( 10^{12} M_\odot / (3 \text{ GeV}) \approx 4 \times 10^{68} \) DM particles in our galaxy, a fraction \( \eta \tau / \tau < 3 \times 10^{-18} \) of which decay into \(^{3}\text{He} \) per year. Even with such a long lifetime we can thus produce up to \( 10^{51} \) \(^{3}\text{He} \) nuclei per year in this scenario.

However, all of these \(^{3}\text{He} \) are by construction non-relativistic, with kinetic energy \( (m_\phi - m_{^{3}\text{He}})^2 / (2m_{^{3}\text{He}}) = 6.5 \text{ MeV} \), far below the kinetic-energy threshold for AMS detection of around GeV/nucleon. Still, a small fraction of these antinuclei might be accelerated by astrophysical processes such as supernovae shock waves, leading to a relativistic flux of antihelium that could explain the preliminary AMS events. Most of the remaining non-relativistic antinuclei will eventually come in contact with normal matter and annihilate, giving rise to characteristic photon spectra that could be detected by the Fermi-LAT experiment. A determination of this fraction is necessary to evaluate whether our model is better constrained by AMS or Fermi-LAT. Note that we ignored induced-decay processes such as \( e^+ \rightarrow \gamma \) that could increase the total number of \(^{3}\text{He} \) and even boost them, but require a dedicated astrophysical simulation.

\( \phi \) could also be detectable in large underground detectors such as Super-Kamiokande through exotic events of the form \( \phi p \rightarrow \overline{\text{p}} \text{He} \). This process can be interpreted as an effective proton lifetime, albeit with an energy release \( E > m_\phi \), that is determined by the interaction rate \([21,22]\)

\[
\tau^\text{eff}_p = \frac{1}{n_\phi \sigma v(\phi p \rightarrow \overline{\text{p}} \text{He})} ,
\]

with DM number density at Earth \( n_\phi = \rho_\phi / m_\phi \approx 0.1 / \text{cm}^3 \), and \( \sigma v \) the cross section between the non-relativistic \( \phi \) and proton. We calculate this cross section with CalcHEP \([23]\) for \( m_\phi = 3 \text{ GeV} \), ignoring formation of final state bound states for simplicity, as

\[
\sigma v(\phi p \rightarrow \overline{\text{p}} \text{He}) \approx 3 \times 10^{-61} \frac{\text{cm}^3}{\text{s}} \left( \frac{7 \times 10^3 \text{ TeV}}{\Lambda} \right)^6 ,
\]

which yields an effective proton lifetime \( \tau^\text{eff}_p \) above \( 10^{24} \text{ yr} \). This is comfortably larger than typical proton lifetime limits, but of course no dedicated search for this dramatic process exists.

**C. \( \phi \) produced by dark matter**

In order to not rely on an exotic astrophysical \(^{3}\text{He} \) acceleration mechanism we can imagine \( \phi \) to be produced boosted by a heavier DM particle \( \chi \), either via annihilations or decays (e.g. \( \chi \rightarrow \phi \phi \) if \( \chi \) carries baryon number \( -6 \) or \( \chi \rightarrow \phi \phi \phi \) if \( \chi \) has baryon number \( -9 \)). In this basic setup there are no CMB constraints on the lifetimes of \( \chi \) or \( \phi \), other than \( \tau_\chi > \tau_\text{Universe} \) by assumption. In particular, the decay rate of \( \phi \) is only constrained by the AMS antihelium flux and UV-considerations for \( \Lambda \), which we address later.

Consider for example \( \chi \rightarrow \phi \phi \) in our galaxy with \( m_\chi = 80 \text{ GeV} \) and \( m_\phi = 3 \text{ GeV} \), which gives a flux of \( \phi \) particles with \( E_\phi = 40 \text{ GeV} \) at Earth of

\[
J_\phi = \frac{2}{4\pi m_\chi \tau_\chi} \int_0^\infty \text{d} s \rho_\text{Halo}[r(s)] \sim 10^6 \left( \frac{t_\text{Universe}}{\tau_\chi} \right) \left( \frac{E_\phi}{\Lambda} \right)^6 ,
\]

where the integral is over the line of sight \([24]\) and we ignored extragalactic contributions. Only a fraction \( d/(\tau_\phi / m_\phi) \) of these boosted \( \phi \) particles will decay on the typical \( d \sim 8 \text{ kpc} \) journey from the center of our galaxy to us, leading to the antihelium flux

\[
J_{^{3}\text{He}} \approx 4 \times 10^{-14} \left( \frac{t_\text{Universe}}{\tau_\chi} \right)^6 \left( \frac{10^3 \text{ TeV}}{\Lambda} \right)^6 ,
\]

with \( E_{^{3}\text{He}} \approx 40 \text{ GeV} \). While this two-body decay mode gives a monochromatic spectrum of antinuclei, more complicated spectra can be obtained from other processes, e.g. \( \chi \rightarrow \phi \phi \phi \). The resulting \(^{3}\text{He} \) flux is clearly large enough to be detectable in AMS as long as \( \Lambda \lesssim 10^3 \text{ TeV} \).

\( ^2 \) An actual comparison of the various antimatter channels has to include the spectral information, so it is entirely possible to also consider the region \( m_\phi \gg 3 \text{ GeV} \), which will not be done here.
in this example, or larger if we increase $m_\phi$. The viability of this in UV-complete models will be discussed below.

On a side note, the heavier DM particle $\chi$ could easily be envisioned to have annihilation channels into $b$-quarks or $\tau$ particles that can produce the $\gamma$-ray excesses observed in the galactic center [25–29] and Andromeda [30].

D. UV completion

As discussed above, Eq. (1) with an effective scale $\Lambda \sim 10^3$–$10^4$ TeV can easily give rise to an observable $^4\text{He}$ flux without other antimatter contributions. However, a full UV-complete model has to be based on quark couplings rather than nucleons, which increases the dimension of the underlying operator. Naively, a three-quark operator will hadronize into one nucleon $N$ via $qqq \rightarrow A_{\text{QCD}}^3 N$, which implies that the $\Lambda$ in our Eq. (1) is related to a quark-level effective-field-theory scale $\Lambda_{\text{UV}}$ via

$$\frac{1}{\Lambda^3} \sim \frac{\Lambda_{\text{QCD}}^3}{\Lambda_{\text{UV}}^3}.$$  \hspace{0.5cm} (13)

$\Lambda_{\text{UV}}$ is thus parametrically suppressed compared to $\Lambda$ and of order $\Lambda_{\text{UV}} \sim 10$–30 GeV in our region of interest, using $\Lambda_{\text{QCD}} \sim 200$ MeV. It is this $\Lambda_{\text{UV}}$ that is related to the masses of the integrated-out mediators, which should be below $4\pi\Lambda_{\text{UV}}$ on account of perturbative unitarity. Such electroweak-scale colored and charged mediators are likely to be excluded by collider searches, but it has to be stressed that we are making very conservative estimates here.

Let us consider variations of our model that do not suffer from this potentially dangerous UV suppression. Instead of $\phi \rightarrow \bar{\tau}\bar{\tau}$ consider $\phi_1 \rightarrow \phi_2 \bar{\tau}\bar{\tau}$, again close to phase-space closure $m_{\phi_1} \sim m_t + m_{\phi_2}$ in order to suppress antiproton production. The discussion is essentially analogous to before upon replacing $m_{\phi} \rightarrow m_{\phi_1} - m_{\phi_2}$ but now we are able to push $m_{\phi_2}$ to the TeV scale. Assuming all new particles, including the colored mediators, to be around the TeV scale makes it possible to enhance the overall $\phi_1 \rightarrow \phi_2 \bar{\tau}\bar{\tau}$ via resonances, i.e. nearly on-shell mediator particles. This is illustrated in Fig. 2 which shows a UV-complete origin of the nine-quark operator relevant for $\phi_1 \rightarrow \phi_2 \bar{\tau}\bar{\tau}$. Here, the scalars $\xi_j$ can lead to a resonant enhancement if their masses are close to $m_{\phi_1}$ because the fermions only carry away small amounts of energy compared to $m_{\phi_1}$. This can compensate for the unwelcome $\Lambda_{\text{QCD}}^9$ suppression that is generic for nine-quark operators and makes it possible to have UV-complete realizations of heavy antimatter production that are consistent with collider constraints.

III. MODELS FOR $^4\text{He}$ PRODUCTION

Above we have discussed simple ways to produce $^3\text{He}$ without accompanying $\bar{d}$ or $\bar{p}$. The same idea can be applied to other antimatter, $^4\text{He}$ being particularly intriguing considering the potential observation of two events in AMS [8]. As an explicit realization of $^4\text{He}$ production in complete analogy to Sec. II let us introduce a new SM-singlet complex scalar $\phi$ with baryon number $-4$ and lepton number $-2$ (the quantum numbers of the $^4\text{He}$ atom). We assume these symmetries to be sufficiently conserved so that they govern the $\phi$ decay modes, with one of the lowest-dimensional operators given by

$$\mathcal{L} = \frac{1}{(2\Lambda)^6} \phi \bar{\pi}^c \gamma_5 n \bar{p} \gamma_5 p \bar{\pi}^c \gamma_5 e + \text{h.c.}.$$  \hspace{0.5cm} (14)

In the non-relativistic limit this corresponds to a coupling $n_\uparrow f_1 p_\uparrow f_1 e_\downarrow$ being the spin up/down states of particle $f$. Operators similar to Eq. (14) but without the $\gamma_5$ would create fermions in the same spin state and thus be suppressed by Fermi statistics in the non-relativistic limit; operators with more than two fermions of the same type, e.g. $\bar{\pi}^c (\gamma_5) n \bar{\pi} (\gamma_5) n \bar{p} (\gamma_5) \nu$ will be suppressed as well. Eq. (14) is then expected to be the dominant operator of interest for $\phi$ decay.

A. Decay channels

As in the previous case, we will consider the situation where $m_\phi$ is close to the total mass of the decay products, i.e. $m_\phi \sim 2m_n + 2m_p$, neglecting again the electron mass $m_e$ for simplicity. The nucleons can be treated as point particles and we can ignore the underlying quark structure. Again, if $m_\phi$ is below the threshold $2m_n + 2m_p$, the decay has to go into a bound state of the nucleons in order to be energetically allowed. The relevant thresholds are given in Tab. II which shows that in the range $0 < m_\phi - m_{^4\text{He}} < 19.81$ MeV, the only kinematically allowed decay induced by Eq. (14) is $\phi \rightarrow ^4\text{He} \bar{\pi} \bar{\pi}$. This is then a source for production of $^4\text{He}$ as long as the lifetime for $\phi$ is not too long. This can indeed be the case despite the small available phase space, as we will now show.

We calculate the decay rate non-relativistically with Fermi’s golden rule,

$$\Gamma(\phi \rightarrow ^4\text{He} \bar{\pi} \bar{\pi}) = 2\pi|V(\phi \rightarrow ^4\text{He} \bar{\pi} \bar{\pi})|^2 \rho,$$  \hspace{0.5cm} (15)

$\rho = (m_\phi - m_{^4\text{He}})^5/(120\pi^4)$ being the phase-space density, well-known from beta decays. $V(\phi \rightarrow ^4\text{He} \bar{\pi} \bar{\pi})$ is the

FIG. 2. UV-realization of a relevant operator for $\phi_1 \rightarrow \phi_2 \bar{\tau}\bar{\tau}$. 

\hspace{0.5cm}
TABLE II. Lowest lying ground states with baryon number 4 and electric charge 2, with mass \( m \) relative to \( ^{4}\text{He}^{++} \). The deuteron and triton nuclei are as usual denoted as \( d^+ \) and \( t^+ \) instead of \(^{2}\text{H}^+\) and \(^{3}\text{H}^+\). We will drop the ionization superscripts in the following. Not shown are excited states, which are instead given in Tab. III.

| \( E - m_{^{4}\text{He}} \) [MeV] | \( J^+ \) | \( \Gamma \) [MeV] | main decay |
|-----------------|----------|---------|----------|
| 0               | 0+       | 0.50    | \( t \, p \)       |
| 20.21           | 0+       | 0.84    | \( t \, p \)       |
| 21.01           | 0-       | 2.01    | \( t \, p \)       |
| 23.33           | 2-       | 5.01    | \( t \, p , ^{3}\text{He} \) |
| 23.64           | 1-       | 6.20    | \( t \, p , ^{3}\text{He} \) |
| 24.25           | 1-       | 6.10    | \( t \, p , ^{3}\text{He} \) |
| 25.28           | 0-       | 7.97    | \( t \, p , ^{3}\text{He} \) |
| 25.95           | 1-       | 12.66   | \( t \, p , ^{3}\text{He} \) |
| 27.42           | 2+       | 8.69    | \( d \, d \)       |
| 28.31           | 1+       | 9.89    | \( t \, p , ^{3}\text{He} \) |
| 28.37           | 1-       | 3.92    | \( d \, d \)       |
| 28.39           | 2-       | 8.75    | \( d \, d \)       |
| 28.64           | 0-       | 4.89    | \( d \, d \)       |
| 28.67           | 2+       | 3.78    | \( d \, d \)       |
| 29.89           | 2+       | 9.72    | \( d \, d \)       |

TABLE III. \(^{4}\text{He}\) states with energy relative to the ground state, \( J^+ \) quantum numbers, decay width, and dominant decay channel. Adopted from Ref. 31.

The appropriate transition probability, here approximated as

\[
|V(\phi \to ^{4}\text{He} \, \pi \, \pi \, \pi)|^2 \approx \sum_{\text{spins}} |M(\phi \to \pi \, \pi \, \pi \, \pi)|^2 \frac{2m_e}{3840 \pi m_\phi} \frac{1}{A^2} \prod_{j=n,p,e} \frac{1}{2E_{j}},
\]

where the first factor is the hard process and the second factor is the overlap of the nucleon wave function with the helium nucleus. The hard process \( \phi(k) \to \pi(k_1) \pi(k_2) \pi(k_3) \pi(k_4) \) has a squared matrix element

\[
\sum_{\text{spins}} |M(\phi \to \pi \, \pi \, \pi \, \pi)|^2 = \prod_{j=n,p,e} \frac{k_{j_1} k_{j_2} m^2}{A^4}.
\]

The overlap \( (\pi \, \pi \, \pi \, \pi)^{4}\text{He}\) of the nucleons with the helium nucleus wave function can be estimated by treating the nucleons as moving in a mean field potential. Once the center of mass motion is factored out, the nucleus is described by a product of three wave functions \( \psi(x_i - x_j) \) for the three relative coordinates. Since the decay of \( \phi \) is assumed to produce the four nucleons initially at a single point, we can estimate the overlap as \( |\langle \pi \, \pi \, \pi \, \pi | ^{4}\text{He} \rangle|^2 \approx |\langle \psi(0) \rangle|^3 \). As before, we take \( |\psi(0)| \sim (100 \text{MeV})^{3/2} \) as a nuclear-physics energy scale. A more accurate calculation of this matrix element is desirable and will be left to future work.

The nucleons are non-relativistic but the electrons relativistic, leading to the decay rate

\[
\Gamma(\phi \to ^{4}\text{He} \, \pi \, \pi \, \pi) \approx \frac{(m_\phi - m_{^{4}\text{He}})^5}{3840 \pi m_\phi A^2} |\psi(0)|^6.
\]

We expect this expression to be valid for \( m_\phi - m_{^{4}\text{He}} \ll m_{^{4}\text{He}} \). With \( |\psi(0)| \approx (100 \text{MeV})^{3} \) this yields the lifetime

\[
\tau \approx 9 \times 10^9 \text{Gyr} \left( \frac{10 \text{MeV}}{m_\phi - m_{^{4}\text{He}}} \right)^5 \left( \frac{A}{100 \text{GeV}} \right)^{12},
\]

shown in Fig. 3 (top). \( \tau \) depends strongly on the mass splitting \( m_{^{4}\text{He}} - m_{^{4}\text{He}} \) and the effective scale \( A \), and can easily surpass the age of the universe, \( t_{\text{universe}} \approx 14 \text{Gyr} \).

For masses \( m_\phi - m_{^{4}\text{He}} > 20 \text{MeV} \), additional final states become energetically allowed, in particular the excited \(^{4}\text{He}\) states of Tab. III. All of these excited states are unbound and unstable 31, but we still expect some of them to contribute as resonances to the decay \( \phi \to ^{4}\text{He} \, \pi \, \pi \, \pi \), followed by the fast decay of \( ^{4}\text{He} \) into \(^{3}\text{He} \, \pi \, \pi \, \pi \), \( ^{3}\text{He} \), \( ^{3}\text{He} \), \( ^{3}\text{He} \), \( ^{3}\text{He} \), \( \Lambda \), and \( \Lambda \), depending on the excited state. This is then a production mode for \(^{3}\text{He} \) and \(^{3}\text{He} \). Due to the resonant enhancement, the decay rates for \( \phi \to ^{4}\text{He} \, \pi \, \pi \pi \) should be formally similar to Eq. 18, naively replacing \( m_\phi \) by \( \psi(0) \) and \( \psi(0) \) by the relevant excited mass and wave function. The latter is unknown to us, but we expect at least some of the excited states to have a similar matrix element as the ground state \(^{3}\text{He} \). As shown in Fig. 3 (bottom), this allows some of the excited states to catch up to the \( \phi \to ^{4}\text{He} \, \pi \, \pi \pi \) rate. While the exact details depend on the excited-state wave functions and branching ratios, we generically expect \( \phi \) decays to produce similar amounts of \(^{4}\text{He} \, ^{3}\text{He} \) and \(^{3}\text{He} \, ^{3}\text{He} \) for \( m_\phi \approx m_{^{4}\text{He}} + 100 \text{MeV} \approx 3.8 \text{GeV} \).

Aside from the three-body decays into excited \(^{4}\text{He} \) states there are four-body decays into ground-state nuclei that should be considered, namely \( \phi \to \overline{\Lambda} \, \pi \, \pi \, \pi \), \(^{3}\text{He} \, \pi \, \pi \, \pi \), and \(^{3}\text{He} \, \pi \, \pi \, \pi \) (Tab. 1). All these four-body decays \( \phi \to \overline{A} \, B \, \pi \, \pi \) with non-relativistic \( A \) and \( \overline{B} \) can be calculated using Fermi’s golden rule,

\[
\Gamma(\phi \to \overline{A} \, B \, \pi \, \pi ) \approx \frac{|\langle \psi(0) \rangle|^4}{270720 \pi^4 A^2 m_\phi} \left( \frac{m_{A} m_{B}}{m_{A} + m_{B}} \right)^{3/2} \times (m_\phi - m_{A} - m_{B})^{1/2},
\]

3 For the first excited state Refs. 32 and 33 calculate a nuclear radius around 3 times larger than that of the ground state, \( r_{^{4}\text{He}} \approx 1.5 \text{fm} \), which we could interpret as a correspondingly smaller matrix element. Since no calculations are available for the other excited states we will not show the result here.
Using the matrix element from Eq. [17] and CalcHEP [23] to calculate the six-body phase space numerically we find the decay rate described to excellent degree by
\[
\Gamma(\phi \to \bar{p}p\bar{p}p\bar{p}p) \simeq \frac{3.3 \times 10^{-19}}{m_\phi A_{12}} (m_\phi + 2m_n + 2m_p)^{19/2} 
\times (m_\phi - 2m_n - 2m_p)^{19/2}.
\]  
(21)

The decay to free nucleons dominates for large \(m_\phi\) and will eventually lead to constraints from antiproton-flux measurements that plague most DM models that aim to produce heavier antimatter [9, 10].

The main conclusion from this analysis is that \(\phi\) decays can produce almost exclusively \(^4\text{He}\) for \(0 < m_\phi - m_{4\text{He}} < 100 \text{ MeV}\), while for \(m_\phi - m_{3\text{He}} \sim 100 \text{ MeV}\) we expect similar amounts of \(^4\text{He}, \, ^4\text{He}, \, \text{and } \bar{p}\). For even larger \(\phi\) masses, where our approximations become invalid, the \(^4\text{He}\) fraction will decrease and the amount of \(\bar{p}\) will become significant, eventually leading back to coalescence fractions for the antimatter.

B. \(\phi\) connection to dark matter and UV completion

In complete analogy to Sec. [11B] we can imagine \(\phi\) to be DM: the CMB bound \(\tau > 10^{11} \text{ s}\) [20] now corresponds to \(\Lambda > 300 \text{ GeV}\) for \(m_\phi = 3.8 \text{ GeV}\), otherwise the same comments apply regarding potential acceleration mechanisms. Exotic events of the form \(\phi p \to \bar{p} p p p p\) in Super-Kamiokande are once again sufficiently suppressed.

Alternatively, in order not to rely on an exotic astrophysical \(^4\text{He}\) acceleration mechanism we can imagine \(\phi\) to be produced boosted by a heavier DM particle \(\chi\), either via annihilations or decays (e.g. \(\chi \to \phi \phi\) if \(\chi\) carries baryon number \(-8\) or \(\chi \to \phi \phi\) if \(\chi\) has baryon number \(-12\)). Analogous to Sec. [11C] this can easily give a large \(^4\text{He}\) flux without accompanying antiproton or antideuteron flux, but requires sub-TeV \(\Lambda\).

This brings us to the main issue of \(^4\text{He}\) production from our model: Eq. [14] has to come from a quark-level operator of dimension \(\geq 22\) with scale \(A_{UV} \lesssim (\Lambda_{QCD}/\Lambda)^2\), which is at the GeV scale for the \(\Lambda\) values of interest to us. \(A_{UV}\) is naively of order of the (colored) mediator masses and thus most likely excluded by collider searches. It is then more useful to study variations of our model along the lines of Sec. [11D] i.e. considering decays \(\phi_1 \to \phi_2 \bar{p}p\bar{p}p\bar{p}\) near threshold with TeV-scale \(\phi_1\) that can be resonantly enhanced by TeV-scale colored mediators. While on the baroque side, models along these lines are so far the only explanation for \(^4\text{He}\) events in AMS that do not rely on hidden antimatter regions [13].

IV. CONCLUSIONS

Standard astrophysical models and even most DM models predict antimatter fluxes that are strongly suppressed for antihelium nuclei, making it very difficult
to explain the preliminary AMS-02 observation of $^3\text{He}$ events and impossible to explain $^4\text{He}$. We have shown here that it is possible to introduce new particles $\phi$ that dominantly decay into $^4\text{He}$ without the typical accompanying antiproton flux. This new particle $\phi$ could in principle be DM with a long lifetime, or be produced by the decay of heavier DM. Models along these lines could potentially explain the AMS-02 $^4\text{He}$ signals.

Future work could proceed in various directions. The UV-complete models have to be studied to evaluate the signatures of the underlying colored mediator particles at the LHC. Improved nuclear-physics calculations are necessary to pin down the relevant matrix elements that determine the fluxes and branching ratios of $^4\text{He}$, $^3\text{He}$, $^d$, and $^p$. Dedicated astrophysical simulations for antinuclei propagation are required to derive the final antimatter fluxes and compare them to indirect constraints e.g. from gamma-ray searches.

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