Bosonic-seesaw portal dark matter

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We discuss a new type of Higgs-portal dark matter (DM) production mechanism, called the bosonic-seesaw portal (BSP) scenario. The BS provides the dynamical origin of the electroweak symmetry breaking, triggered by mixing between the elementary Higgs and a composite Higgs generated by a new-color strong dynamics, hypercolor (HC). At the HC strong coupling scale, the classical-scale invariance assumed in the model is dynamically broken, as well as the “chiral” symmetry present in the HC sector. In addition to the composite Higgs, HC baryons emerge to potentially be stable because of the unbroken HC baryon number symmetry. Hence the lightest HC baryon can be a DM candidate. Of interest in the present scenario is that HC pions can be as heavy as the HC baryon due to the possibly enhanced explicit “chiral”-breaking effect triggered after the BS mechanism, so the HC baryon pair cannot annihilate into HC pions. As in the standard setup of the freeze-in scenario, it is assumed that the DM was never in the thermal equilibrium, which ends up with no thermal abundance. It is then the non-thermal BSP process that crucially comes into the game below the HC scale: the HC baryon significantly couples to the standard-model Higgs via the BS mechanism, and can non-thermally be produced from the thermal plasma below the HC scale, which turns out to allow the TeV mass scale for the composite baryonic DM, much smaller than the generic bound placed in the conventional thermal freeze-out scenario, to account for the observed relic abundance. Thus the DM can closely be related to the mechanism of the electroweak symmetry breaking.

1. Introduction

A quarter of our universe is constituted by unknown matter, called dark matter (DM). Several cosmological and astrophysical observations have so far suggested that the DM should be electrically neutral, and cold enough and stable enough to be long-lived compared with the age of the universe. The abundance left in the present universe is thought to have been produced via interactions with the standard model (SM) particles in the early universe, involving mediators such as the SM Higgs (Higgs portal scenario). For some basic references, see Refs. [1,2].

In this paper, we discuss a new type of Higgs-portal DM-production mechanism, which we call the bosonic-seesaw portal (BSP) scenario. The DM candidate will be identified as a composite baryonic state, arising as a bound state of new fermions strongly coupled in a new-color dynamics—hypercolor (HC)—which triggers the electroweak symmetry breaking (EWSB) via the BS mechanism [3–5], which also generates the portal coupling between the DM candidate and the 125 GeV Higgs boson, dubbed as the BSP coupling. It is the BSP process that produces the DM relic abundance, the right amount of which can be achieved in accordance with the BS mechanism.
Composite baryonic dark matter has been discussed in Refs. [3,6,7]. The DM abundance in those studies is thermally produced through the annihilation to the HC pions, where, by a naive analogy to QCD, the thermal relic abundance can be estimated as \( \Omega_{\text{DM}} h^2 \approx 10^{-5} \) when we suppose the order of TeV mass for the DM. This implies that to realize the correct amount of the presently observed DM abundance, a general limit on the composite baryonic DM mass is set by the thermal freeze-out relic abundance, \( m_{\text{DM}} = \mathcal{O}(100) \text{ TeV} \) [3,6,7].

In the scenario addressed here, in contrast to the conventional freeze-out scenario, the DM abundance is not thermally produced from the strong HC sector itself. The production of the DM takes place non-thermally by the BSP process (which is à la freeze-in scenario [8,9]), outside of the HC sector. It then turns out that the BSP production allows the DM mass to be of the order of TeV. Thus the relic abundance of the somewhat light composite DM directly links to the EWSB mechanism.

The present scenario would generically be operative as long as the following two setups are at hand:

(I) the classical scale invariance is present to be dynamically broken by the HC sector weakly coupled to the SM sector, yielding the BS mechanism.

(II) HC pions become massive after the BS mechanism (i.e., below the strong HC coupling scale) due to the weak coupling to another hidden sector, which dynamically enhances the explicit breaking effect of “chiral” symmetry present in the HC sector, to make the HC pion mass as large as the HC dynamical scale, i.e. the HC baryon mass scale (DM mass scale). In this setup, the HC baryon cannot annihilate into HC pions, and the inclusive annihilation cross section would yield no thermal relic abundance (which would be the initial condition for the DM thermal history).

To demonstrate the point, as a concrete example we shall take a BS model discussed in Refs. [4,5].

2. Scenario description: part (I)

We begin by employing the former two sectors in the above part (I), which are constructed from the classically scale-invariant SM and the HC gauge sector of \( SU(N_{\text{HC}}) \). In the model we have the HC gluon as well as the HC fermions \( F_{L/R} \) forming the \( SU(3)_{F_{L/R}} \) flavor triplet, \( F_{L/R} = (\chi_{L/R}, \psi_{L/R})^T \), where \( \chi \) denotes the \( SU(2)_{FL/R} \) doublet. These HC fermions carry the vector-like charges under the \( SU(3)_c \times SU(2)_W \times U(1)_Y \times SU(N_{\text{HC}}) \) gauge groups, \( \psi_{L/R} \sim (1, 1, 0, N_{\text{HC}}) \) and \( \chi_{L/R} \sim (1, 2, 1/2, N_{\text{HC}}) \). The HC fermions couple to the elementary Higgs doublet \( H \) with the small coupling \( y \) as

\[
-y \cdot \bar{F}_L \begin{pmatrix} 0_{2 \times 2} & H \\ H^\dagger & 0 \end{pmatrix} F_R + \text{h.c.} = -y \cdot \bar{\chi} \bar{H} \psi + \text{h.c.} \quad (1)
\]

At around the scale \( \Lambda_{\text{HC}} \), the HC gauge coupling gets so strong that the “chiral” \( SU(3)_{F_L} \times SU(3)_{F_R} \) symmetry is spontaneously broken by the “chiral” (but SM vector-like) condensate \( \langle \bar{F}^a F^b \rangle \sim -4\pi f_{\pi_{\text{HC}}}^3 / \sqrt{N_{\text{HC}}C_{ab}} (a,b = 1, \ldots, 8) \), down to the vectorial \( SU(3)_F \), where \( f_{\pi_{\text{HC}}} \sim \sqrt{N_{\text{HC}}} \cdot \Lambda_{\text{HC}} / (4\pi) \) is the HC pion decay constant. After the HC confinement, furthermore, the HC fermions form composite HC hadrons in a manner similar to the QCD case. Among those HC hadrons, a composite HC scalar doublet \( \Theta \sim \bar{\psi} \chi (\Theta^\dagger \sim \bar{\chi} \bar{\psi}) \), embedded in the \( SU(3)_F \) flavor nonet, has the same quantum numbers as the elementary Higgs doublet \( H \). Hence this \( \Theta \) mixes with the elementary
Higgs doublet $H$, due to the presence of the above Yukawa term, so that one finds the mixing form

$$-y \cdot f_\Theta m_\Theta (H^\dagger \Theta) + \mathrm{h.c.},$$

where $m_\Theta$ denotes the mass of the composite Higgs doublet $\Theta$, which is of $\mathcal{O}(\Lambda_{\mathrm{HC}})$, and $f_\Theta$ is the decay constant associated with the scalar current $\bar{\chi} \psi$ coupled to the $\Theta$, defined as $\langle 0 | \bar{\chi} \psi(0) | \Theta \rangle = f_\Theta m_\Theta$. The induced mixing term yields the mass matrix for $(H, \Theta)^T$:

$$
\begin{pmatrix}
0 & y\Lambda_{\mathrm{HC}}^2 \\
y\Lambda_{\mathrm{HC}}^2 & \Lambda_{\mathrm{HC}}^2
\end{pmatrix},
$$

with the mixing strength $y \ll 1$ controlling the coupling between the SM and the HC sectors. In Eq. (3) we have simply taken $f_\Theta \simeq m_\Theta \simeq \Lambda_{\mathrm{HC}}$, which can be expected from the QCD case. Note that the determinant of the mass matrix is negative, so the negative mass-squared of the SM Higgs is dynamically generated by the seesaw mechanism (the BS mechanism [10–12]) to trigger the EWSB. (For detailed potential analysis, see Ref. [5].)

In addition to the composite Higgs, the HC sector generically involves rich composite spectra such as the HC pions and baryons, as in the case of QCD. Among those HC hadrons, some of HC baryons, presumably the lightest ones, can be stable due to the conserved HC baryon number, and hence can be a DM candidate. Here we shall suppose that the HC sector possesses $N_{\mathrm{HC}} = 4$, i.e. the $SU(4)_{\mathrm{HC}}$. The HC baryon can be realized as a complex scalar, having the HC scalar-baryon charge. (Our argument is substantially unchanged even if we employ the case other than the $SU(4)_{\mathrm{HC}}$ in which fermionic HC baryons can be present, as will be clarified below.) Among the HC baryons, the EW-singlet HC scalar-baryon, $\varphi \sim \psi \psi \psi \psi$, can be the lightest, i.e., the DM candidate. The $\varphi$ mass is expected to be of the order of $\mathcal{O}((N_{\mathrm{HC}}/3)\Lambda_{\mathrm{HC}})$. In the present article, we take $\Lambda_{\mathrm{HC}} = \mathcal{O}(\mathrm{TeV})$, and hence $m_\varphi = \mathcal{O}(\mathrm{TeV})$ as well.

The DM candidate $\varphi$, the EW-singlet complex scalar baryon, strongly and minimally couples to the composite HC Higgs doublet, $\Theta$, like

$$a \cdot \varphi^\dagger \varphi \Theta^\dagger \Theta,$$

with the order one (or larger) coefficient $a$. By the BS mechanism, the $\Theta$ starts to mix with the elementary Higgs doublet $H$ below the scale $\Lambda_{\mathrm{HC}}$. This dynamically generates a Higgs portal coupling between the DM $\varphi$ and the SM Higgs $H_1$:

$$\kappa_{\varphi H} \cdot \varphi^\dagger H_1^\dagger H_1, \quad \text{with} \quad \kappa_{\varphi H} = ay^2,$$

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1 The composite Higgs doublet $\Theta$ should not be confused with pseudo-Nambu–Goldstone bosons as in extant composite Higgs models: all “chiral” HC pions in the present scenario are CP-odd pseudoscalars ($\pi_{\mathrm{HC}}^{ab} \sim \tilde{F}^a i\gamma_5 F^b$ with $J^P = 0^-$), not CP-even scalars.

2 In terms of the $SU(3)$ flavor nonet in QCD, the $\Theta$ can be viewed as an analogue of the scalar meson $K_0^0(1430)$ with $I(J^P) = \frac{1}{2}(0^+)$. 

3 The HC baryon number associated with the unbroken $U(1)_{F_Y}$ is necessarily conserved, as long as the HC dynamics is vector-like and the HC fermions are vector-like charged, as in the present model.

4 In Refs. [13–15], a similar composite dark-bosonic baryon to the DM candidate, called stealth DM, has been discussed in a context different than the bosonic seesaw.

5 The coefficient $a$ does not scale with $N_{\mathrm{HC}}$, i.e., $a = \mathcal{O}(N_{\mathrm{HC}}^0)$, because it is the coupling of the baryon–meson two-body scattering.
where the factor $y^2$ has come from the BS mixing strength $y$ [$\Theta \approx y H_1 + H_2$, which can be understood by diagonalizing the mass matrix Eq. (3)] between the SM Higgs $H_1$ and heavy Higgs $H_2$. The mixing strength $y$ is supposed to be much smaller than $O(1)$, so that the Higgs portal coupling $\kappa_{\psi H}$ can naturally be small to be consistent with the present relic abundance of dark matter, as will be clarified later on.

3. Scenario description: part (II)

The spontaneous breaking of the “chiral” $SU(3)_F \times SU(3)_F$ symmetry gives rise to eight Nambu–Goldstone bosons, HC pions ($\pi_{\text{HC}}$). The “chiral” symmetry is explicitly broken by the $y$-Yukawa interaction in Eq. (1), which makes HC pions in part massive with the tiny mass $\Delta m^2_{\pi_{\text{HC}}} |_{\psi} \sim O(y \cdot v \Lambda_{\text{HC}})$, where $v$ is the EW scale $\sim 246$ GeV ($= \mathcal{O}(\Lambda_{\text{HC}}/5)$). Also, the EW gauge interactions would slightly lift up the masses for some of the HC pions charged under the EW, $\Delta m^2_{\pi_{\text{HC}}} |_{\text{EW}} \sim O(\alpha_{\text{em}} \Lambda_{\text{HC}}^2)$. Note that by those explicit breaking effects, the HC pions cannot have a mass as large as the HC scale. To make the HC pions heavy enough by explicit breaking outside of the HC and EW dynamics, one needs another hidden sector explicitly breaking the “chiral” symmetry, as described in part (II) in Sect. 1.

In the BS model proposed in Refs. [4,5], such an explicit breaking term is supplied by a pseudoscalar ($S$) in the scale-invariant form as

$$g_S \cdot (\bar{F}i\gamma_5 F) S,$$

with the weak coupling $g_S (\ll 1)$, which ensures conservation of the approximate “chiral” symmetry in the theory. Below the HC confinement scale $\sim \Lambda_{\text{HC}}$, the pseudoscalar $S$ dynamically develops its vacuum expectation value, $v_S$, due to the pseudoscalar seesaw mechanism between $S$ and the HC $\eta'$ meson [4,5], and makes all the HC pions massive by the $g_S$-Yukawa coupling as [5]

$$m_{\pi_{\text{HC}}} \sim \left( \frac{\Lambda_{\text{HC}}}{f} \right) g_S v_S,$$

where $f = f_{\pi_{\text{HC}}} / \sqrt{N_{\text{HC}}/3} (\sim \Lambda_{\text{HC}}/(4\pi))$ with the HC pion decay constant $f_{\pi_{\text{HC}}}$. Thus the large $v_S$ enhances the explicit breaking effect to make it possible to lift the HC pion mass to reach a scale as large as $O(\Lambda_{\text{HC}})$, i.e., the HC baryon ($\psi$) mass scale:

$$m_{\pi_{\text{HC}}} \sim m_{\psi} \sim O(\Lambda_{\text{HC}}).$$

Actually, as discussed in Ref. [5], the $v_S$ is required to be much larger than $\Lambda_{\text{HC}}$ to realize the EWSB through the potential analysis, where the HC pion mass is indeed lifted up to be of the order of $\Lambda_{\text{HC}}$, even for the small $g_S$ coupling. Crucial to note here is that this explicit-breaking enhancement is triggered after the HC confinement/the “chiral” symmetry breaking (as well as the BS mechanism) takes place, above which scale the $g_S$ term does not play any role, so the HC fermions are (almost) massless. Hence this effect becomes relevant only for the HC pions, not for HC fermions; or for HC baryons, either.\footnote{This point is also different from the heavy pseudoscalar in QCD, such as pseudoscalar mesons including charm and bottom quarks, which acquire the large mass because of the large current quark masses $m_c, m_b$.}

\footnote{The amplification of the explicit breaking effect irrespective of bare fermion mass has been addressed in a different strong dynamics in Ref. [16], and other references therein.}
Note also that the presence of the large $\nu_S$ is necessary to realize the EWSB scale in the present BS model, as was demonstrated in Ref. [5].

The heavy HC pion will significantly affect the thermal history of the lightest HC baryon $\phi$: the $\phi$ was generated when the temperature cools down to the critical scale of $O(\Lambda_{HC})$. Since the $\phi$ mass is of the same order as $\Lambda_{HC}$, analogously to the case of QCD light baryons (e.g. neutron and proton), the number density of $\phi$ gets diluted no sooner than the HC confinement. Due to the possible breaking enhancement as in Eq. (8), such a $\phi$ cannot kinematically annihilate into HC pions, so that the conventional thermal freeze-out would not happen, which could drastically alter the composite baryonic DM scenario, as will be discussed in detail below.\footnote{The heavy HC pions, having a mass of $O(\text{TeV})$ as in Eq. (8), including electromagnetically neutral ones, completely decay to EW gauge bosons with a width of $O(100)\text{ GeV}$ [5], to quickly disappear below the HC scale $\Lambda_{HC}$.}

4. Non-thermal DM production by BSP

Now we discuss the production mechanism for the composite baryonic DM based on the scenario descriptions (I) and (II). Of importance to note is the presence of the Higgs portal coupling in Eq. (5), dynamically induced from the BS mechanism at $T = \Lambda_{HC}$. Although the number density of the $\phi$ has been diluted, due to the induced portal coupling, the $\phi$ can be produced unilaterally via the SM sector such as SM $\rightarrow \phi\phi^\dagger$. Note that, although the composite DM strongly couples in the HC sector, the BSP interaction is perturbative due to the suppression by $y \ll 1$, hence can be computed reliably enough by the usual perturbation theory. It thus turns out that the BSP interaction is non-thermal for the whole time scale in the thermal history and the produced abundance saturates slightly after the generation of $\phi$: at around $x = m_{\phi}/T \simeq \Lambda_{HC}/T \simeq 5$—see Figs. 1 and 2. (If the HC confinement happened at a somewhat lower temperature, say lower by factor of 10 than the HC baryon mass, $T_c \sim 1/10\Lambda_{HC}$, as in the case of QCD, the thermal history described in Fig. 2 would be changed just by making the starting position of the rising up shift to $x = m_{\phi}/T \sim 10$.\footnote{Then the HC and EW phase transitions might happen almost simultaneously in the thermal history. In that case the abundance estimate for the BSP process produced from SM particles could be more complicated than that done in the present analysis.})

As noted in the previous section, the HC pions become heavy enough that the DM population is not sufficiently created by the heavy HC pions, and the DM annihilation into two HC pions is kinematically blocked. Even in that case, however, the thermal DM abundance produced by the total inclusive annihilation cross section might not be negligibly small because of the strong interaction of the HC sector. Then the thermal history of the DM would be changed from the one in Fig. 2 and reduced to be a conventional freeze-out scenario, where the freeze-out takes place at around $x \sim 20–30$, which is much later than the BSP production at $x \sim 5$. To separate our present scenario from the conventional freeze-out scenario, we shall therefore assume that the DM was never in the thermal equilibrium, as done in the usual freeze-in scenario as a natural setting [8,9]. With this assumption at hand, we may then safely set the initial condition for the number density per comoving volume, $Y(T = m_{\phi}) = 0$.

We thus evaluate the production cross sections arising from the BSP coupling in Eq. (5). Since the $\phi$ is generated at $T = \Lambda_{HC}$ and the saturation point ($T_f$) at around the EW scale ($T_f = m_{\phi}/5 = O(\nu)$; see Fig. 2), the relevant processes governed by the population in the thermal bath are: $hh, f\bar{f}, WW, ZZ \rightarrow \phi\phi^\dagger$. Those cross sections are computed at the tree level of the...
Fig. 1. The ratio of the BSP reaction rate to the Hubble parameter, $$\Gamma_{\text{BSP}}/H$$, with the number density of the (non-relativistic and bosonic) DM $$\varphi$$ ($$n_\varphi \simeq 1/\pi^2 (m_\varphi T/2\pi)^3 e^{-m_\varphi/T}$$) and the thermal-averaged cross section estimated by summing the production cross sections in Eq. (9). In the plot we have assumed radiation dominance for the Hubble parameter $$H$$ ($$H \simeq 0.33 \sqrt{g^*_\varphi (T)} T^2/M_p$$) with the effective degree of freedom $$g^*_\varphi (T)$$ set to 100, and taken the BSP coupling $$\kappa_{\varphi H} = 10^{-10}$$ consistent with the estimate in Eq. (11).

Fig. 2. A sample of the thermal history of the HC scalar baryon $$\varphi$$ yield (number density per entropy density) $$Y = n/s$$ against the temperature $$x = m_\varphi/T$$ along the bosonic seesaw portal production, with the mass $$m_\varphi = 1$$ TeV (solid curve) and 5 TeV (dashed curve), and the coupling $$\kappa_{\varphi H} = 10^{-10}$$ fixed (for $$N_{\text{HC}} = 4$$). The effective degree of freedom $$g_*(T)$$ has been set to 100 in the plot.

perturbation in couplings to be

\[
\sigma(hh \to \varphi^\dagger \varphi) = \frac{\kappa_{\varphi H}^2}{16\pi} \frac{s}{(s - m_h^2)^2} \frac{1 - 4m_W^2/s}{\sqrt{1 - 4m_h^2/s}} \left( \frac{m_W^2/s}{1 + \sqrt{1 - 4m_W^2/s}} \right)^{5/2} 
\]

\[
\sigma(WW/ZZ \to \varphi^\dagger \varphi) = \frac{9\kappa_{\varphi H}^2}{64\pi} \frac{s}{(s - m_h^2)^2} \frac{1 - 4m_W^2/s}{\sqrt{1 - 4m_h^2/s}} \left( \frac{m_W^2/s}{1 + \sqrt{1 - 4m_W^2/s}} \right)^2 \left[ 2 + \left( \frac{s}{2m_W^2} \right)^2 \right]
\]

\[
\sigma(f\bar{f} \to \varphi^\dagger \varphi) = \frac{N_f^2 \kappa_{\varphi H}^2}{32\pi} \frac{s}{(s - m_h^2)^2} \frac{1 - 4m_f^2/s}{\sqrt{1 - 4m_h^2/s}} \left( \frac{m_f^2/s}{1 + \sqrt{1 - 4m_h^2/s}} \right) \left[ 1 - \frac{4m_f^2/s}{1 + \sqrt{1 - 4m_h^2/s}} \right] 
\]
Fig. 3. The constraints on the BSP coupling $\kappa_{\psi H}$ from the presently observed relic abundance ($\simeq 0.1$ drawn as the red horizontal line in the plot) for cases of the bosonic baryon ($N_{\text{HC}} = 4$) and the fermionic baryon ($N_{\text{HC}} = 3$). The plots have been produced by taking the mass to be 1 TeV (solid cyan curves) and 5 TeV (dashed black curves).

with $N_{\ell} = 3(1)$ for quarks (leptons) and $\sqrt{s}$ is the center of mass energy. Here we have neglected the Higgs width because the effective range of $\sqrt{s}$ is much above the Higgs mass scale. The number density per entropy density today, $Y(T_0) = n(T_0)/s(T_0)$, can be calculated by integrating the Boltzmann equation with the production cross sections $\sigma(ij \to \phi^\dagger \phi)$ and the boundary condition $Y(T = m_\phi \simeq \Lambda_{\text{HC}}) = 0$ to be

$$Y(T_0) \sim Y(T_f) = \frac{135\sqrt{10} M_p \xi^2(3)}{32\pi^7} \int_{T_f}^{\Lambda_{\text{HC}}} dT \sum_{i,j} \frac{g_i g_j \eta_i \eta_j}{[g_*(T)]^{3/2}} \int_{(m_i + m_j)/T}^{\infty} dx \times x^4 K_1(x) \sigma(ij \to \phi^\dagger \phi) \left(1 - \frac{(m_i + m_j)^2}{x^2 T^2}\right) \left(1 - \frac{(m_i - m_j)^2}{x^2 T^2}\right), \quad (10)$$

where $g_*(T)$ stands for the effective degree of freedom for relativistic particles, $g_i = 2(1)$ and $\eta_i = 3/4(1)$ for fermions (bosons), $M_p \simeq 10^{18}$ GeV (reduced Planck mass), $K_1(x)$ denotes the modified Bessel function of the first kind, $\xi(3) \simeq 1.202$, $x \equiv \sqrt{s}/T$.

The relic abundance, $\Omega_\psi h^2 = Y(T_0) \cdot m_\phi s(T_0)/(\rho_c h^{-2})$, turns out to actually be almost independent of the $\phi$ mass as far as $\phi$ masses of $\mathcal{O}(\text{TeV})$ is concerned (see Fig. 3). The BSP coupling $\kappa_{\psi H}$ is then constrained by the presently observed DM relic abundance $\simeq 0.1$. Figure 3 shows the constraint plot on the portal coupling for the scalar baryon DM ($N_{\text{HC}} = 4$), as well as the fermionic baryon case ($N_{\text{HC}} = 3$) in which $\phi \sim \psi \bar{\psi} \bar{\psi}$. The figure tells us the upper bounds,

$$\kappa_{\psi H} \lesssim 10^{-10} (10^{-11}),$$

or

$$y \lesssim 10^{-5} (10^{-6}) \times (1.0/a)^{1/2}, \quad (11)$$

for $N_{\text{HC}} = 4(3)$, where $g_*(T)$ in Eq. (10) has been taken to be $\simeq 100$. The smallness of the BSP coupling, $\kappa_{\psi H}$ or $y$ in Eq. (11), is consistent with the BS mechanism, and indeed implies that the

In the case of $N_{\text{HC}} = 3$ the $\phi \sim \psi \bar{\psi} \bar{\psi}$ would actually be a spin 3/2 baryon, not the Dirac fermion of the ground state in terms of the spin statistics, though it can be stable by the HC-baryon number conservation.
BSP is non-thermal. Note also that the size of the small $y$ does not affect the realization for the Higgs boson mass of 125 GeV: after solving the potential problem and fluctuating the Higgs field around the EW vacuum, one finds that the Higgs mass is given as $m_h \simeq \sqrt{2\lambda_H} v$ with the quartic coupling of the elementary Higgs $\lambda_H$, which is precisely the same mass formula as in the SM (for details, see Ref. [5]).

5. Conclusion and discussion

In conclusion, the bosonic-seesaw portal scenario, based on the model descriptions (I) and (II) proposed in the present paper, provides a dark matter candidate having a coupling to the standard model Higgs, which is dynamically generated by the seesaw and essentially related to the origin of the electroweak symmetry breaking. In this scenario the dark matter candidate dynamically arises as the hypercolor baryon with the conserved hypercolor-baryon charge. Sufficient composite baryonic dark matter can be non-thermally produced, due to the significantly small coupling to the standard model Higgs as a consequence of the bosonic seesaw mechanism, to realize the observed relic amount which allows for the composite dark matter to have TeV-scale mass, in contrast to the conventional thermal freeze-out scenario.

Having the Higgs portal coupling, the composite baryonic dark matter $\phi$ can be detected by direct detection experiments such as LUX [17], PandaX-II [18,19], and the upcoming XENON1T and LZ [20]. The spin-independent (SI) cross section is computed as $\sigma_{SI}(\phi N \rightarrow \phi N) \simeq \frac{\kappa_H^2}{16\pi m_h} m_{\phi}^2 \left( N, \phi \right) g_{NN}^2$, where $g_{NN} \simeq 0.25 \text{GeV}/v$ [21–23], and $m_{\phi} = m_N m_{\phi}/(m_N + m_{\phi})$ is the reduced mass with $m_N \simeq 940 \text{MeV}$. Using the upper bound for the portal coupling $\kappa_{\phi H}$ in Eq. (11), and taking the dark matter mass 1–5 TeV for the reference value, we find the upper bound on the SI cross section, $\sigma_{SI} \lesssim 10^{-63}(10^{-65}) \text{cm}^2$ for $N_{HC} = 4(3)$. These values are far below the current limit most stringently set by LUX2016 [17], and the sensitivity in the anticipated XENON1T or LZ, $\sigma_{SI} \lesssim 10^{-47} \text{cm}^2$ at the TeV range [20], which will actually be overlapped with the expected neutrino background [24].

Thus, having an extremely small portal coupling, the bosonic-seesaw portal dark matter is fairly insensitive to direct detection experiments, which would imply some extension for the present scenario, or other detection proposals.

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