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Vector-like technineutron Dark Matter: is a QCD-type Technicolor ruled out by XENON100?

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Abstract One of the specific predictions of a new strongly coupled dynamics at a TeV scale is the existence of stable vector-like technibaryon states so that the lightest neutral one could serve as a Dark Matter candidate. We study the latter hypothesis in the QCD-type technicolor with $SU(3)_{TC}$ confined group and one $SU(2)_{W}$ doublet of vector-like techniquarks consistent with electroweak precision constraints and test it against the existing Dark Matter astrophysics data. We discuss the most stringent Dark Matter constraints on weak interactions of technibaryons in $SU(3)_{TC}$ technicolor and possible implications of these findings for the cosmological evolution of relic technineutrons. We conclude that vector-like techniquark sectors with an odd group of confinement $SU(2n+1)_{TC}$, $n = 1, 2, \ldots$ and with ordinary vector-like weak $SU(2)_{W}$ interactions are excluded by XENON100 data under the assumption of technibaryon number conservation in the modern Universe allowing for an even technicolor group $SU(2n)_{TC}$ only.

1 Introduction

The undoubtful existence of the Dark Matter (DM) comprising about a third (or more precisely, about 27 % [1]) of the total mass of the Universe today remains the strongest phenomenological evidence for New Physics beyond the Standard Model (SM) required by astrophysics measurements. The hypothetical weakly interacting massive particles (WIMPs), which the DM is possibly composed of, and their properties are yet undiscovered at the fundamental level, while the DM itself is being regarded as one of the major cornerstones of modern theoretical astrophysics and cosmology [2]. Such an uneasy situation motivates ongoing search for appropriate Particle Physics candidates for WIMPs away from constantly improving observational bounds.

Traditionally, the lightest supersymmetric particles (LSPs) predicted by supersymmetry (SUSY) [3,4] such as the neutralino are often referred to as the best DM candidates [5], and this is considered to be one of the major advantages of SUSY-based SM extensions (for an overview of existing DM candidates, see e.g. Refs. [6,7] and references therein). Direct SUSY searches are currently ongoing at the LHC and major direct and indirect DM detection experiments, so that the parameter space of the simplest SUSY scenarios is getting more and more constrained (for the most recent exclusion limits and their effects on SUSY DM candidates see e.g. Refs. [8–11]).

Here, we are focused on one of the alternatives to SUSY-based DM candidates predicted by dynamical electroweak symmetry breaking (EWSB) and compositeness scenarios, the lightest heavy neutral technibaryon (or T-baryon) state N. In case of the odd QCD-type $SU(3)_{TC}$ group of confinement extending the SM gauge group such a candidate is often referred to as the Dirac T-neutron in analogy to ordinary neutrons from low-energy hadron physics. The idea of composite DM candidates has a long history, starting from the mid-1980s and is from Refs. [12,13] where it has been claimed that an excess of T-baryons possibly built up in the early Universe can explain the observed missing mass. So far, a number of different models of composite DM candidates and hypotheses about their origin and interactions have been proposed. Generic DM signatures from technicolor-based models with stable T-baryons were discussed e.g. in Refs. [14–18] (for a review see also Ref. [19] and references therein). In particular, well-known minimal dynamical EWSB mechanisms predict relatively light T-baryon states as pseudo Nambu–Goldstone bosons of the underlying gauge theory [20,21]. The latter can naturally provide asymmetric DM candidates if one assumes the existence of a T-baryon asymmetry in Nature similarly to ordinary baryon asymmetry [22]. Having similar mechanisms for ordinary matter and DM formation

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in the early Universe one would expect the DM density to be of the same order of magnitude as that of baryons. Depending on a particular realization of the dynamical EWSB mechanism such composite DM candidates may be self-interacting, which helps in avoiding problematic cusp-like DM halo profiles [23].

All of the existing composite DM models rely on the basic assumption about New Physics extension of the SM by means of the extra confined matter sectors. These ideas were realized in a multitude of technicolor (TC) models developed so far [24,25] (for a detailed review on the existing TC models, see e.g. Refs. [19,26]). Historically, the first TC models with dynamical EWSB are based upon the idea that the Goldstone degrees of freedom (technipions or T-pions) appearing after the global chiral symmetry breaking $SU(2)_L \otimes SU(2)_R \to SU(2)_V$ are absorbed by the SM weak gauge bosons, which thereby gain masses. The dynamical EWSB mechanism is then triggered by the condensate of fundamental technifermions (or T-quarks) in confinement, $\langle \tilde{Q}\bar{\tilde{Q}} \rangle \neq 0$. Traditional TC models with dynamical EWSB are faced with the problem of the mass generation of standard fermions, which was consistently resolved in the Extended TC model [27,28]. However, many of the existing TC-based models have got severely constrained or often are ruled out by the EW precision data [29,30]. Generally, in these schemes noticeable contributions to strongly constrained Flavor Changing Neutral Current (FCNC) processes appear, together with too large contributions to the Peskin–Takeuchi (especially, to the S) parameters. Further developments of the TC ideas have resulted in the Walking TC and the vector-like (or chiral-symmetric) TC, which succeeded in resolving the above-mentioned problems and there remain viable scenarios of the dynamical EWSB [31–34].

In this paper, we continue the investigation of promising phenomenological implications of the vector-like TC model proposed recently in Ref. [34]. This is one of the simplest successful realizations of the bosonic TC scenarios—an extension of the SM above the electroweak (EW) scale which includes both a Higgs doublet H and a new strongly coupled *vector-like* techniquark sector (for different realizations of the bosonic TC ideas, see e.g. Refs. [35–39]). In contrast to conventional (Extended and Walking) TC models, in the vector-like TC model the mechanism of the EWSB and generation of SM fermions masses is driven by the Higgs vacuum expectation value (vev) in the standard way, irrespectively of the (elementary or composite) nature of the Higgs field itself. Similarly to other bosonic TC models, the Higgs field H develops a vev which in turn is induced by the T-quark condensate. Thus, it is possible to assimilate the SM-like Higgs boson while the Higgs vev acquires a natural interpretation in terms of the T-quark condensates. This means that the Higgs mechanism is not the primary source of the EWSB, but that it is effectively induced by unknown TC dynamics at high scales.

The vector-like TC model [34] is based upon the phenomenologically successful gauged linear σ model (GLσM) initially proposed in Ref. [40] and further elaborated in Refs. [41–44]. It is well known that in the low-energy limit of QCD and in the limit of massless u and d quarks, the resulting QCD Lagrangian with switched-off weak interactions of u, d quarks possesses *exact global chiral* $SU(2)_L \otimes SU(2)_R$ symmetry. The physical degrees of freedom in this Lagrangian are given by a superposition of initially chiral quark fields—*the Dirac u,d-quark fields*. Global $SU(2)_L \otimes SU(2)_R$ is then considered as *a classification symmetry of composite states* giving rise to the lightest hadrons in the physical spectrum and nicely predicting their properties. This model predicts the lightest *physical* pseudoscalar T-pion $\tilde{\pi}$, scalar T-sigma $\tilde{\sigma}$ fields as well as T-baryon states classified according to representations of the gauged vector subgroup $SU(2)_{V\equiv L+R}$ of original global chiral $SU(2)_L \otimes SU(2)_R$ symmetry. Its complete gauging is also possible at the composite level, giving rise to an effective field theory describing the 'chiral-gauge' interactions between bound states in adjoint (e.g. composite vector/pseudovector fields and pions) and fundamental (e.g. composite baryons, constituent 'dressed-up' quarks) representations. But this gauging makes sense *only* at the level of bound states, but not at the fundamental level—in the high-energy (UV) limit of this effective field theory the chiral group remains global.

As usual, we assume that the spontaneous chiral symmetry breaking in the T-hadron sector happens in the chiral-symmetric (vector-like) way,

$$SU(2)_L \otimes SU(2)_R \to SU(2)_{V\equiv L+R}. \tag{1}$$

In Ref. [34] it was argued that in the low-energy limit the vector-like gauge group $SU(2)_V$ can be approximately *identified* with the weak isospin group $SU(2)_W$ of the SM. *The latter identification at the bound-state level should be understood simply as a natural phenomenological trick to introduce local weak interactions into the T-hadron spectrum valid in the low-energy effective field theory limit only*, and it can be schematically written as

$$SU(2)_{V\equiv L+R} \simeq SU(2)_W. \tag{2}$$

Such a 'gauging' of the vector subgroup $SU(2)_{V\equiv L+R}$ in the GLσM sense and its identification with the SM gauge isospin group do not mean that one introduces extra elementary gauge bosons into the existing fundamental theory, e.g. to the SM or its possible high-scale gauge extensions. In our context, the procedure (2) means a very simple thing: both T-quarks and T-hadrons interact with *already existing* gauge bosons in the SM in the low-energy effective field theory limit with local gauge couplings [34]. In the high-energy limit of the theory, the global chiral symmetry $SU(2)_L \otimes SU(2)_R$ is restored, while Dirac vector-like T-quark fields, along with chiral SM fermion fields, reside in fundamental representations of the SM gauge $SU(2)_W$ group of the SM. As one of the important features of the VLTC model, after the chiral symmetry breaking in the T-quark sector the left and right components of the original Dirac T-quark fields can interact with the SM weak $SU(2)_W$ gauge bosons with vector-like couplings, in opposition to ordinary SM fermions, which interact under $SU(2)_W$ by means of their left-handed components only.

The VLTC scenario represents the very first step focusing on the low-energy implications of a new strongly coupled dynamics with chiral-symmetric UV completion—the first relevant step for searches for such a dynamics at the LHC and in astrophysics—*formally* keeping the elementary Higgs boson as it is in the one-doublet SM, which does not satisfy the naturalness criterion. From the theoretical point of view, the model points out a promising path towards a consistent formulation of composite Higgs models in extended chiral-gauge theories with chiral-symmetric UV completion. The latter can be further exploited in the composite Higgs models as well as in attempts to attain Grand-like TC unification with the SM at high scales (see e.g. Refs. [45–48] and references therein). Such a strongly coupled sector survives the EW precision tests with minimal vector-like confined sector (U and D T-quarks) without any extra assumptions. Excluding the naturalness criterion, three other important points which are considered to be primary achievements of the VLTC model [34] can be summarized as follows:

- The effective Higgs mechanism of the dynamical EW symmetry breaking naturally emerges in this approach and is especially pronounced in the limit of approximate conformal symmetry. Namely, the Higgs vev is induced by the T-quark condensate so that the EW symmetry is broken simultaneously with the chiral symmetry. No T-pions are eaten and remain physical, they escape current detection limits due to suppressed loop-induced couplings to, at least, two (in odd $SU(3)_{TC}$ case) or three (in even $SU(2)_{TC}$ case) gauge bosons at the leading order, and they can remain light.
- The phenomenologically consistent minimal vector-like UV completion is based upon the linear σ model with the global chiral $SU_L(2) \otimes SU_R(2)$ symmetry group. In the minimal formulation, the model operates with two vector-like T-flavors passing the EW constraints and results in (almost) standard Higgs couplings in the limit of small $h\tilde{\sigma}$ mixing. The model can be, in principle, used to explain possibly small Higgs couplings deviations from the standard ones if confirmed experimentally.
- There are specific phenomenological consequences of such a new dynamics at the LHC, e.g. light hardly detectable technipions with multi-boson final states produced via a suppressed VBF only, and possible distortions of the Higgs boson couplings and especially self-couplings, vector-like T-baryon states at the LHC with a large missing-E_T and asymmetry signatures as well as implications for cosmology (vector-like T-baryon Dark Matter).

The proposed scenario, at least in its simplest form discussed in Ref. [34], does not attempt to resolve the naturalness/hierarchy problem of the SM and does not offer a mechanism for the generation of the current T-quark masses. It is considered as a low-energy phenomenologically consistent limit of a more general strongly coupled dynamics, which is yet to be constructed (it has the same status as the low-energy effective field theories existing in hadron physics).

Even though the EW precision constraints are satisfied for any $SU(n)_{TC}$ group with vector-like weak interactions [34], it is still an open question, if astrophysics constraints are satisfied for any $SU(n)_{TC}$ group as well. The DM exclusion limits, therefore, become an extra important source of information about TC dynamics which has the power to constrain the parameter space of the vector-like TC model even more. One of the unknowns we would like to consider here is the rank of the confined group n. In particular, we will discuss for which $SU(n)_{TC}$ groups in confinement it is possible to make the identification of the gauge groups (2) in the T-quark/T-baryon sectors, and for which it is not, based upon existing constraints from DM astrophysics. The latter will be our main conclusion of this work.

The paper is organized as follows. Section 2 provides a brief overview of vector-like $SU(3)_{TC}$ TC model with Dirac T-baryons with generic weak-type $SU(2)$ interactions (before the identification (2)). Section 3 contains a discussion of the T-baryon mass spectrum; in particular, an important mass splitting between T-proton and T-neutron. In Appendix A, we consider typical T-baryon annihilation processes in the cosmological plasma in two different cases—in the high- and low-symmetry phases. Section 4 is devoted to a discussion of cosmological evolution of T-neutrons in two cases of symmetric and asymmetric DM. In Sect. 5, major implications of direct detection limits to the vector-like T-neutron DM model are outlined. It was shown that weakly $SU(2)_W$ interacting Dirac vector-like T-baryons are excluded by recent XENON100 data [49], which poses an important constraint on the rank of the confined group in the T-quark sector under the condition (2). Finally, Sect. 6 contains basic concluding remarks.

2 Vector-like T-baryon interactions

For simplicity, here we adopt the simplest version of the Standard Model with one Higgs doublet, and the question whether it is elementary or composite is not critical for our further considerations [34]. The new heavy physical states of the model (additional to those in the SM) are the singlet T-sigma $\tilde{\sigma}$, triplet of T-pions $\tilde{\pi}_a$, $a = 1, 2, 3$, $\mathrm{SU}(2)_{\mathrm{W}}$ doublets of constituent T-quarks $\tilde{Q}$ and T-baryons $\tilde{N}$,

$$\tilde{Q} = \begin{pmatrix} U \\ D \end{pmatrix}, \qquad \tilde{N} = \begin{pmatrix} P \\ N \end{pmatrix}. \tag{3}$$

The latter acquire masses via the T-quark condensate, acting as an external source in the potential and the T-sigma vev. We focus on phenomenological studies of such a low-energy effective field theory at typical momentum transfers squared with $Q^2 \ll \Lambda_{\mathrm{TC}}^2$ in the linear σ-model framework. The corresponding Lagrangian responsible for the Yukawa-type interactions of the T-quarks and T-baryons, Eq. (3), reads

$$\mathcal{L}_Y^{\mathrm{TC}} = -g_{\mathrm{TC}}^{Q} \bar{\tilde{Q}}(S + i\gamma_5 \tau_a P_a)\tilde{Q} - g_{\mathrm{TC}}^{N} \bar{\tilde{N}}(S + i\gamma_5 \tau_a P_a)\tilde{N},$$
$$g_{\mathrm{TC}}^{Q} \neq g_{\mathrm{TC}}^{N}, g_{\mathrm{TC}}^{Q,N} > 1, \tag{4}$$

where τ_a, $a = 1, 2, 3$ are the Pauli matrices, and techni-strong Yukawa couplings g_{TC}^{Q} and g_{TC}^{N} are introduced in a complete analogy to low-energy hadron physics, they absorb unknown non-perturbative strongly coupled dynamics and can be chosen to be different. Typically, the perturbativity condition requires them to be bounded, $g_{\mathrm{TC}}^{Q,N} < \sqrt{4\pi}$, in order to trust predictions of the linear model. After the EWSB phase, the Yukawa interactions (4) will play an important role in determining the strength of the T-neutron N self-interactions leading to specific properties of the associated DM; this will be discussed below.

In the SM, the ordinary gauge boson–hadron interactions are usually introduced by means of gauge bosons hadronization effects. In the case of a relatively large T-confinement scale $\Lambda_{\mathrm{TC}} \sim 0.1$–$1$ TeV relevant for our study, the effect of T-hadronization of light W, Z bosons into heavy composite states is strongly suppressed by large constituent masses of T-quarks $M_Q \sim \Lambda_{\mathrm{TC}}$. Following the arguments of Ref. [34], the vector-like interactions of $\tilde{Q}$, $\tilde{N}$ and P_a fields with initial $U(1)_{\mathrm{Y}}$ and $\mathrm{SU}(2)_{\mathrm{V}}$ gauge fields B_μ, V_μ^a, respectively, can be safely introduced in the local approximation via covariant derivatives over the local $\mathrm{SU}(2)_{\mathrm{V}} \otimes U(1)_{\mathrm{Y}}$ group in the same way as ordinary SM gauge interactions, i.e.

$$\mathcal{L}_{\mathrm{kin}}^{\mathrm{TC}} = \frac{1}{2}\partial_\mu S\, \partial^\mu S + \frac{1}{2} D_\mu P_a\, D^\mu P_a + i\bar{\tilde{Q}}\hat{D}\tilde{Q} + i\bar{\tilde{N}}\hat{D}\tilde{N}. \tag{5}$$

Here, covariant derivatives of $\tilde{Q}$, $\tilde{N}$ and P_a fields with respect to $\mathrm{SU}(2)_{\mathrm{V}} \otimes U(1)_{\mathrm{Y}}$ interactions read

$$\hat{D}\tilde{Q} = \gamma^\mu \left(\partial_\mu - \frac{iY_Q}{2} g_1 B_\mu - \frac{i}{2} g_2^V V_\mu^a \tau_a\right)\tilde{Q},$$
$$\hat{D}\tilde{N} = \gamma^\mu \left(\partial_\mu - \frac{iY_N}{2} g_1 B_\mu - \frac{i}{2} g_2^V V_\mu^a \tau_a\right)\tilde{N}, \tag{6}$$
$$D_\mu P_a = \partial_\mu P_a + g_2^V \epsilon_{abc} V_\mu^b P_c,$$

respectively; besides that $\tilde{Q}$ is also assumed to be confined under a QCD-like $\mathrm{SU}(n)_{\mathrm{TC}}$ group. Below, for the sake of simplicity we discuss a particular case with the number of T-colors $n = 3$; we analyze a possible implementation of EW interactions into T-quark/T-baryon sectors according to the

$$\mathrm{SU}(2)_{\mathrm{V}} \to \mathrm{SU}(2)_{\mathrm{W}}, \quad V_\mu^a \to W_\mu^a, \quad g_2^V \to g_2 \tag{7}$$

replacement rule in Eq. (6). A consistency test of the latter scenario against the DM relic abundance and direct DM detection data for the rank-2 confined group will enable us to draw important conclusions about the properties of the TC sectors.

The theory in its simplest formulation as discussed here, of course, does not predict particular values for the elementary and composite T-quark hypercharges Y_Q and Y_N. These, together with the number of T-quark generations, the respective properties of interactions, the group of the confinement, etc. should be ultimately constrained in extended chiral-gauge or grand-unified theories along with the coming (collider and astrophysics) experimental data.

Employing further analogies with the SM and QCD, in what follows we fix the hypercharge of the elementary T-quark doublet to be the same as that of quark doublets in the SM, i.e. $Y_Q = 1/3$, and the hypercharge of the T-nucleon doublet—to be the same as that of the nucleon doublet in the SM, i.e. $Y_N = 1$. Thus, the T-baryon states in Eq. (3) become T-nucleons composed of three elementary T-quarks, i.e. $P =$ (UUD), $N =$ (DDU) in analogy to proton and neutron in QCD. Other assignments with different hypercharges and a different number of T-colors are also possible and would lead to other possible types of T-baryons. The basic qualitative results for the DM properties in odd $\mathrm{SU}(3)_{\mathrm{TC}}$ confined group of QCD-type are generic for other odd $\mathrm{SU}(2n+1)_{\mathrm{TC}}$, $n = 2, 3, \ldots$ groups. In this work we stick to a direct analogy with QCD for simplicity and test it against available DM constraints.

One interesting alternative opportunity would be to consider the $Y_Q = 0$ case, so that an integer electric charge of T-baryons would only be possible for even TC groups $\mathrm{SU}(2n)_{\mathrm{TC}}$ with the simplest $\mathrm{SU}(2)_{\mathrm{TC}}$. Here, T-baryons are two-T-quark systems. In the non-perturbative T-hadron vacuum the UD state with zero electric charge is energetically favorable since an extra binding energy appears due to exchanges of collective excitations with T-pion quantum numbers (in usual hadron physics the effect of ud-coupling brings up extra 70 MeV into the binding energy) making the neutral di-T-quark UD state absolutely stable and thus rendering it an appealing DM candidate. This case has certain advantages and will be considered elsewhere.

A standard option proposed in Ref. [34] would be to adopt the approximate identification (7) without introducing any gauge bosons at the composite level. Note that it can only be approximate since in the minimal formulation the model does not account for a possible (very small) mixing with heavy vector T-mesons. Alternatively, if one does not impose a straightforward identification (7), after local $SU(2)_V \otimes U(1)_Y$ symmetry breaking an extra set of heavy gauge Z', $W'^{\pm}$ bosons interacting with constituent T-quark and T-baryons appears. As a practically important example, such extra Z', $W'^{\pm}$ bosons can be identified with composite $\tilde{\rho}^{0,\pm}$ mesons which are often treated as 'gauge bosons' with respect to the local chiral subgroup $SU(2)_V$ in the GLσM framework [40–44]. For the sake of generality, we will keep Z' and W' notations in what follows, without a thorough discussion of their particular origin. In the latter case, T-quarks and T-baryons may still interact with ordinary SM gauge fields via a (very small) mixing between Z' and Z, $W'^{\pm}$ and $W^{\pm}$ bosons, respectively. Of course, such a mixing must be tiny not to spoil the EW precision tests. In one way or another, the vector-like SM gauge interactions of T-baryons with Z, $W^{\pm}$ bosons are controlled by the following part of the Lagrangian:

$$L_{\bar{\tilde{N}}\tilde{N}Z/W} = \delta_W \frac{g_2}{\sqrt{2}}\,\bar{P}\gamma^\mu N \cdot W^+_\mu + \delta_W \frac{g_2}{\sqrt{2}}\,\bar{N}\gamma^\mu P \cdot W^-_\mu + \delta_Z \frac{g_2}{c_W}\, Z_\mu \sum_{f=P,N} \bar{f}\gamma^\mu \big(t_3^f - q_f\, s_W^2\big) f. \qquad (8)$$

Here, $\delta_{W,Z}$ are the generic parameters which control EW interactions of T-baryons, $e = g_2 s_W$ is the electron charge, t_3^f is the weak isospin ($t_3^P = 1/2$, $t_3^N = -1/2$), and $q_f = Y_N/2 + t_3^f$ is the T-baryon charge. The two consistent options for introducing weak interactions into the T-fermion sectors dictated by EW precision tests can be summarized as follows:

$$\begin{array}{llll} \text{I.} & \delta_{W,Z} \simeq 1, & SU(2)_V \simeq SU(2)_W, & \\ \text{II.} & \delta_{W,Z} \ll 1, & SU(2)_V \neq SU(2)_W, & m_{Z',W'} \gg 100\,\text{GeV}. \end{array} \qquad (9)$$

In the first case, one deals with pure EW vector interactions of T-quarks/T-baryons corresponding to the transition (7), while in the second case $\delta_{W,Z}$ are related to a very small mixing between SM vector bosons and extra *different* $SU(2)_V$ bosons tagged as Z' and $W'^{\pm}$. In both cases, couplings with photons are not (noticeably) changed and are irrelevant for Dirac T-neutron DM studies, so they are not shown here. We will test both options above against the available constraints on T-neutron DM, implying the existence of $SU(3)_{TC}$ group in confinement.

As agreed above, we choose $Y_N = 1$ in analogy to the SM, thus $q_P = 1$ and $q_N = 0$ as anticipated. The Yukawa-type interactions of T-baryons with scalar (h and $\tilde{\sigma}$) and pseudoscalar ($\tilde{\pi}^{0,\pm}$) fields are driven by

$$L_{\bar{\tilde{N}}\tilde{N}h} + L_{\bar{\tilde{N}}\tilde{N}\tilde{\sigma}} + L_{\bar{\tilde{N}}\tilde{N}\tilde{\pi}} = -g^N_{TC}\,(c_\theta\tilde{\sigma} + s_\theta h)\cdot(\bar{P}P + \bar{N}N) - i\sqrt{2}g^N_{TC}\,\tilde{\pi}^+\bar{P}\gamma_5 N - i\sqrt{2}g^N_{TC}\,\tilde{\pi}^-\bar{N}\gamma_5 P - i g^N_{TC}\,\tilde{\pi}^0(\bar{P}\gamma_5 P - \bar{N}\gamma_5 N). \qquad (10)$$

The gauge and Yukawa parts of the Lagrangian (8) and (10) completely determine the T-baryon interactions at relatively low kinetic energies $E_{\text{kin}} \ll M_{B_T}$ typical for equilibrium reactions (scattering, production and annihilation) processes in the cosmological plasma before DM thermal freeze-out epoch (see below). Note that due to the vector-like nature of extra virtual T-baryon states they do not produce any noticeable contributions to the oblique corrections and FCNC processes preserving internal consistency of the model under consideration [34]. The latter is true for both models I and II, (9).

In considering the vector-like TC model, the T-baryon mass scale $\sqrt{s} \lesssim M_{B_T}$ should be considered as an upper cut-off of the model considered, which contains only the lightest physical d.o.f. $\tilde{\pi}$ and $\tilde{\sigma}$. The latter are sufficient in the current first analysis of the vector-like T-baryon DM in the non-relativistic limit $v_B \ll 1$. Certainly, at higher energies $\sqrt{s} \gtrsim M_{B_T}$ the theory should involve higher (pseudo)vector and pseudoscalar states (e.g. $\tilde{\rho}, \tilde{a}_0, \tilde{a}_1$ etc.). The latter extension of the model will be done elsewhere if required by phenomenology.

3 T-proton–T-neutron mass splitting

A typical (but optional) assumption about a dynamical similarity between color and T-color in the case of confined $SU(3)_{TC}$ enables us to estimate characteristic masses of the lightest T-hadrons and constituent T-quarks through the scale transformation of ordinary hadron states via an approximate scale factor $\zeta = \Lambda_{TC}/\Lambda_{QCD} \gtrsim 1{,}000$ following from a relative proximity of EW scale and $\Lambda_{TC} \sim 0.1$–1 TeV, i.e.

$$m_{\tilde{\pi}} \gtrsim 140\,\text{GeV}, \quad M_{\tilde{\sigma}} \gtrsim 500\,\text{GeV}, \quad M_Q \gtrsim 300\,\text{GeV}, \quad M_{B_T} \equiv M_P \simeq M_N \gtrsim 1\,\text{TeV}, \qquad (11)$$

for the T-pion $m_{\tilde{\pi}}$, T-sigma $M_{\tilde{\sigma}}$ and constituent T-quark M_Q and T-baryon M_{B_T} mass scales. In QCD, the constituent quark masses roughly take a third of the nucleon mass, so it is reasonable to assume that the same relation holds in T-baryon spectrum,

$$M_{B_T} \equiv M_P \simeq M_N \simeq 3M_Q = 3g^Q_{TC}u, \qquad (12)$$

where $u \sim \Lambda_{TC} = 0.1 - 1$ TeV is the T-sigma vacuum expectation value (vev) which spontaneously breaks the local chiral symmetry in the T-quark sector down to weak isospin group (1) (for more details on the chiral and EW symmetries breaking in the considered model, see Ref. [34]). In the chiral

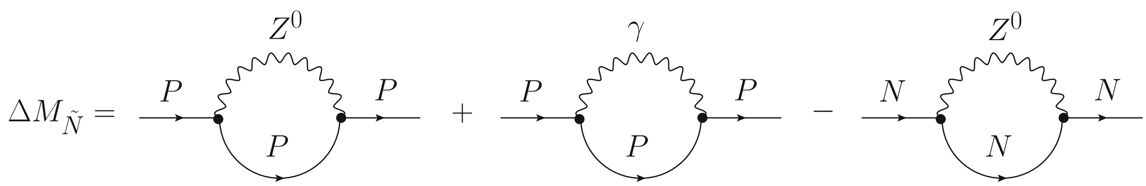

Fig. 1 One-loop EW radiative corrections causing positive mass splitting between T-proton and T-neutron in the chiral limit of the underlined theory, $\Delta M_{B_T}^{EW} > 0$. Other corrections from $W^{\pm}$, T-pion, T-sigma, and Higgs boson loops enter symmetrically into P and N self-energies and thus do not contribute to the mass splitting $\Delta M_{B_T}^{EW}$ and not shown here

limit of the theory, T-sigma vev u has the same quantum-topological nature as the SM Higgs vev $v \simeq 246\,\text{GeV}$, i.e. $u,\ v \sim |\langle\bar{\tilde{Q}}\tilde{Q}\rangle|^{1/3}$ in terms of the T-quark condensate $|\langle\bar{\tilde{Q}}\tilde{Q}\rangle| \neq 0$, providing the dynamical nature of the EWSB mechanism in the SM.

Also, with respect to interactions with known particles at typical 4-momentum squared transfers $Q^2 \ll l_{TC}^{-2} \gtrsim 2.3\,\text{TeV}^2$, where l_{TC} is the characteristic length scale of the non-perturbative T-gluon fluctuations estimated by rescaling of that from QCD, Eq. (11), the T-hadrons behave as elementary particles with respect to EW interactions. Besides DM astrophysics, not very heavy vector-like T-baryons can also be relevant for the LHC phenomenology as well, which is an important subject for further studies.

Adopting the hypothesis of T-baryon number conservation in the modern Universe in analogy to ordinary baryon number, let us find constraints on the vector-like TC model parameters providing an inverse mass hierarchy between T-neutron and T-proton, i.e. $M_N < M_P$. In this case, T-neutron becomes indeed the lightest T-baryon state and hence stable, which makes it an appealing DM candidate.

In usual hadron physics it is well known that the isospin SU(2) symmetry at the level of current quark masses is strongly broken—the current mass difference between u and d quarks is of the order of their masses. Such a symmetry is restored to a good accuracy at the level of constituent quarks and nucleons. This restoration is a direct consequence of the smallness of the current quark masses compared to contributions from the non-perturbative quark–gluon vacuum to the hadron masses. A small mass splitting in hadron physics is typically estimated in the baryon–meson theory which operates with hadron-induced corrections (in particular, ρ-meson loops with a ρ–γ mixing).

In the case of the local vector-like chiral subgroup $SU(2)_V$ in both models I and II we neglect T-rho $\tilde{\rho}$ mediated contributions to respective DM annihilation cross sections assuming for simplicity that $\tilde{\rho}$ mixing with γ and Z is very small due to a strong mass hierarchy between them. In this simplified approach we can evaluate the *lower bound* on the T-baryon mass splitting induced by pure EW corrections only (other EW-like gauge interactions and non-local effects may only increase it). The techni-strong interactions do not distinguish isotopic components in the T-baryon doublet, and thus do not contribute to the mass splitting between P and N.

When the loop momentum becomes comparable to the T-baryon mass scale $M_{B_T} \gg m_Z$ or higher, a T-baryon parton substructure starts to play an important role. In particular, the local approximation for EW T-baryon interactions does not work any longer and one has to introduce non-local Pauli form factors instead of the local gauge couplings. The latter must, in particular, account for the non-zero anomalous magnetic moments of the T-baryons. In the EW loop corrections to the T-baryon mass splitting, however, typical momentum transfers q which dominate the corresponding (finite) Feynman integral are at the EW scale, the $M_{EW} \sim 100$ GeV scale. At such scales we can safely neglect non-local effects and use ordinary local gauge couplings renormalized at $\mu^2 = M_{B_T}^2$ scale. The latter approximation is sufficient for a rough estimate of the mass splitting and, most importantly, its sign.

Note the coincidence of T-isotopic $SU(2)_V$ symmetry at the fundamental T-quark level with the weak isospin $SU(2)_W$ of the SM providing arbitrary but *exactly equal* current T-quark masses in the initial Lagrangian, $m_U = m_D$, so that the fundamental T-quark and hence T-baryon spectra are degenerate at tree level. Note that in the initial SM Lagrangian current u, d-quark masses are equal to zero due to the chiral asymmetry of the weak interactions. After spontaneous EW symmetry breaking very different current u, d-quark masses emerge as a consequence of the absence of SU(2) interactions for right-handed u, d quarks.

In the considered case of a degenerate vector-like T-quark mass spectrum $\Delta M_Q \equiv M_U - M_D = 0$, the EW radiative corrections dominate the mass splitting between T-proton and T-neutron for suppressed heavy T-rho $\tilde{\rho}$ contributions and a small $\tilde{\rho}$-gauge bosons mixing, $\Delta M_{B_T}^{EW} \lesssim \Delta M_{B_T} \equiv M_P - M_N \ll M_{B_T}$. The corresponding EW ($Z, \gamma$-mediated) one-loop diagrams are shown in Fig. 1.

In model I, Eq. (9), the T-baryon mass splitting $\Delta M_{B_T}^{EW}$ due to EW corrections is given in terms of the difference between the T-baryon mass operators on mass shell, which takes the form of the following finite integral:

$$\Delta M_{B_{\rm T}}^{\rm EW} = -\frac{ie^2 M_Z^2}{8\pi^4} \int \frac{(\hat{q} - M_{B_{\rm T}})\,{\rm d}q}{q^2(q^2 - M_Z^2)\left[(q+p)^2 - M_{B_{\rm T}}^2\right]}, \quad (13)$$

given by γ and Z corrections shown in Fig. 1 only. Note that logarithmic divergences explicitly cancel out in the difference between the P and N mass operators, providing us with a finite result for $\Delta M_{B_{\rm T}}^{\rm EW}$. Other corrections from $W^{\pm}$, scalar, and pseudoscalar loops enter symmetrically into the P and N self-energies and thus are canceled out too. In the realistic case of heavy T-baryons $M_{B_{\rm T}} \gg M_Z$ we arrive at the following simple relation:

$$\Delta M_{B_{\rm T}}^{\rm EW} \simeq \frac{\alpha(M_{B_{\rm T}}) M_Z}{2} > 0, \quad (14)$$

where the fine structure constant $\alpha(\mu)$ is fixed at the T-baryon scale $\mu = M_{B_{\rm T}} \sim 1$ TeV. Numerically, we find to a good accuracy

$$\Delta M_{B_{\rm T}}^{\rm EW} \simeq 360 \text{ MeV}. \quad (15)$$

This can be considered as the EW contribution to the T-baryon mass splitting and provides a conservative lower limit to it. The T-proton is not stable and weakly decays into T-neutron and light SM fermions ($e,\ \mu,\ \nu_{e,\mu},\ u,\ d,\ s$) as follows: $P \to N + (W^* \to f_i \bar{f}_j)$. Remarkably enough, the EW radiative corrections appear to work in the right direction, making the T-proton slightly heavier than the T-neutron, so that the latter turns out to be stable and viable as a heavy DM candidate. With the mass splitting value (15), we find the following approximate vector-like T-proton lifetime:

$$\tau_P \simeq \frac{15\pi^3}{G_F^2} (\Delta M_{B_{\rm T}}^{\rm EW})^{-5} \simeq 0.4 \times 10^{-9} \text{ s}. \quad (16)$$

In model II, the Z contributions die out in the limit $\delta_{\rm EW} \ll 1$, so it can only be induced by extra heavy Z' exchange. The corresponding contribution to the P–N mass splitting $\Delta M_{B_{\rm T}}^{\rm V}$ is obtained from Eq. (14) by a replacement $m_Z \to m'_Z$,

$$\Delta M_{B_{\rm T}}^{\rm V} \simeq \frac{\alpha(M_{B_{\rm T}}) M_{Z'}}{2} > 0, \qquad M_{Z'} < M_{B_{\rm T}}, \quad (17)$$

so that $\Delta M_{B_{\rm T}}^{\rm V} \gg \Delta M_{B_{\rm T}}^{\rm EW}$ and the T-proton lifetime would even be shorter. Of course, the estimate (17) should be taken with care for $m_{Z'} \gtrsim M_{B_{\rm T}}$ when non-local effects become important, and a radiative mass splitting between constituent U and D T-quarks would determine the actual mass difference between P and N.

Note that in the most natural and the simplest model, model I, the properties of the vector-like T-baryon spectrum are very similar to the properties of the vector-like Higgsino LSP (e.g. splitting between chargino and neutralino) spectrum due to practically the same structure of EW interactions [50,51]. The key difference between the lightest Higgsino and T-neutron DM candidates is in the capability of T-neutrons to self-interact (e.g. we have enhanced self-annihilation and elastic scattering rates) driven essentially by techni-strong Yukawa terms (10), which make them specifically interesting for DM phenomenology and astrophysics.

4 Cosmological evolution of vector-like T-baryons

The T-proton lifetime (16) has an age of the order of the Universe at the EW phase transition epoch $t_{\rm EW}$. So after this epoch the T-proton remnants quickly decay into T-neutrons and light SM fermions without leaving significant traces on the cosmological evolution of cold DM and ordinary matter. Indeed, due to a rather small T-nucleon mass splitting (15), the amount of T-baryons frozen off the cosmological plasma will be approximately equal to the T-neutron/anti-T-neutron DM density extrapolated back to the freeze-out epoch. So T-proton/anti-T-proton decay processes and a relative fraction of T-protons and T-neutrons at freeze-out time scale can be disregarded in analysis of the DM relic abundance, at least, to a first approximation. Let us consider two possible scenarios of T-baryon DM relic abundance formation—symmetric and asymmetric DM cases.

4.1 Scenario I: Thermal freeze-out of symmetric T-baryon Dark Matter

We consider the thermal evolution of the T-baryon density in the cosmological plasma in the case of symmetric DM, i.e. when number densities of T-baryons and anti-T-baryons are (at least, approximately) equal at all stages of Universe evolution until the present epoch, i.e.

$$n_{B_{\rm T}} \simeq n_{\bar{B}_{\rm T}}, \qquad n_{B_{\rm T}} - n_{\bar{B}_{\rm T}} \ll n_{B_{\rm T}}, \quad (18)$$

so the chemical potential of the T-baryon plasma can be neglected to a sufficiently good accuracy. In the considered model the T-baryons are Dirac particles so $B_{\rm T}$ does not coincide with $\bar{B}_{\rm T}$, while these particles are always produced and annihilate in $B_{\rm T}\bar{B}_{\rm T}$ pairs, at least, at relevant temperatures $T < M_{B_{\rm T}}$.

For details of the thermal evolution and freeze-out of symmetric heavy relic, see e.g. Ref. [59] and references therein, while here we repeat just a few relevant formulas. The irreversible T-baryon annihilation process in the cosmological plasma is initiated at the moment t_0 and temperature $T_0 = T(t_0)$, when the mean energy of relativistic quarks and leptons is comparable with T-baryon mass scale $M_{B_{\rm T}}$ in the radiation-dominated epoch, i.e. $\bar{\varepsilon}_f \simeq 3\,T_0 \simeq M_{B_{\rm T}}$ and the Hubble parameter is $H = 1/2t$. The chemical equilibrium with respect to T-baryon annihilation/production processes breaks down at this moment, and the residual density of T-baryons in the plasma $n_{B_{\rm T}} = n_{B_{\rm T}}(t)$ at later times $t \gg t_0$

is then described by a standard solution of the Boltzmann evolution equations [59].

4.1.1 High-symmetry phase

At temperatures $T > T_{\rm EW} \sim 200$ GeV the Higgs condensate $\langle H\rangle \equiv v$ is melted, i.e. $v = 0$, and thus weak isospin $\mathrm{SU}(2)_{\rm W}$ of the SM is restored, while the T-sigma condensate does not vanish $\langle S\rangle \equiv u \neq 0$, $u \gg T_{\rm EW}$, so that the chiral symmetry in the fundamental T-quark sector is broken: $\mathrm{SU}(2)_{\rm L} \otimes \mathrm{SU}(2)_{\rm R} \to \mathrm{SU}(2)_{\rm V}$. In what follows, we refer to this period in the cosmological evolution as *the high-symmetry (HS) phase of the cosmological plasma* with the characteristic temperature $T_{\rm EW} < T \lesssim u$. This means that the T-baryon mass scale should be well above the EW scale for the DM relic abundance to be formed entirely in the HS phase, i.e. $M_{B_{\rm T}} \gg 200\,\mathrm{GeV}$.

In the HS case with $\mathrm{SU}(2)_{\rm V} = \mathrm{SU}(2)_{\rm W}$ (model I), the equilibrium number densities of (anti)T-neutrons and (anti)T-protons are equal to each other, $n_N = n_P$ ($n_{\bar N} = n_{\bar P}$) since the T-baryon mass spectrum is degenerate, i.e. $\Delta M_{B_{\rm T}} = 0$ (or more precisely, $T \gg \Delta M_{B_{\rm T}}$), so the total T-baryon number density is $n_{B_{\rm T}} \simeq 2(n_N + n_{\bar N})$, at least before the T-baryon freeze-out epoch. Practically, the P and N states are dynamically equivalent in this phase and participate in all reactions as components of the isospin $\mathrm{SU}(2)_{\rm V}$ doublet $\tilde N$, Eq. (3), with Yukawa and gauge interactions determined by Eqs. (4) and (5), respectively. Consequently, the masses of all SM fermions and gauge bosons vanish in this phase (more precisely, $m_f,\ M_{W,Z} \ll T$), while the T-sigma $M_{\tilde\sigma}$ and T-pion $m_{\tilde\pi}$ masses (related as $M_{\tilde\sigma} \simeq \sqrt{3}m_{\tilde\pi}$ in the limiting 'no $\tilde\sigma h$-mixing' case) do not vanish but are likely to be much smaller than the T-baryon mass scale $M_{B_{\rm T}} \sim u$ since $u \gg v$, i.e. $M_{\tilde\sigma},\ m_{\tilde\pi} \ll M_{B_{\rm T}}$. Thus, all the masses except for the T-baryon mass can be neglected in practical calculations to a good approximation. So, in the HS phase we effectively end up with the single T-baryon mass scale parameter $M_{B_{\rm T}}$, which has to be constrained together with the strong Yukawa coupling $g_{\rm TC}^N$ from astrophysics data.

In the HS phase of the evolution of the Universe, the T-baryon annihilation epoch effectively begins at the moment of physical time

$$t_0 \simeq \frac{1}{4T_0^2}\left(\frac{3}{2\pi G w}\right)^{1/2} \ll t_{\rm EW}, \qquad T_0 \simeq \frac{M_{B_{\rm T}}}{3}, \tag{19}$$

where G is the gravitational constant, $w = g_*(T)\pi^2/30$ is the statistical weight of cosmological plasma at the T-baryon annihilation epoch, and $g_* = g_*(T)$ is the effective number of relativistic d.o.f. in the plasma. At high temperatures $T > T_{\rm EW} \sim 200\,\mathrm{GeV}$ before the EW phase transition, in the SM with one Higgs doublet we have $g_* = 106.75$. At lower temperatures, $T < m_W$, the value $g_* = 86.25$ can be used. The steepest change in g_* occurs around the QCD phase transition $T_{\rm QCD} \sim 100\,\mathrm{MeV}$ when it drops down to $g_* \simeq 10$, but the latter does not affect the heavy T-baryons evolution at late times, since they have already decoupled from the plasma while the annihilation rate is practically negligible on average.

Assuming that the annihilation epoch occurs entirely in the HS phase, it should terminate before the EW phase transition time as soon as T-baryons drop off of the chemical equilibrium. Then the freeze-out of heavy T-baryons happens at $t = t_1$ when the temperature of the Universe is $T_1 = T_1(t_1) \gtrsim T_{\rm EW}$, given by the standard formula

$$T_1 \simeq \frac{M_{B_{\rm T}}}{\log\left(\frac{g_{B_{\rm T}} M_{B_{\rm T}} M_{\rm PL}^*(\sigma v_B)_{\rm ann}}{(2\pi)^{3/2}}\right)} \ll M_{B_{\rm T}}, \quad M_{B_{\rm T}} \ll M_{\rm PL}^*, \tag{20}$$

valid to a logarithmic accuracy. Here, $M_{\rm PL}^* = M_{\rm PL}/1.66\sqrt{g_*}$ is the reduced Planck mass, and $g_{B_{\rm T}}$ is the number of T-baryon d.o.f. The phase of the cosmological plasma where the DM gets effectively frozen out; i.e. the actual relation between T_1 and the EW phase transition temperature $T_{\rm EW}$, depends on the details of the DM scenario, or on the typical $M_{B_{\rm T}}$ and $(\sigma v_B)_{\rm ann}$ values in our case. The typical weak interactions strength leads to a crude order-of-magnitude estimate: $(\sigma v_B)_{\rm ann} \sim \alpha_{\rm W}^2/M_{\rm EW}^2$, $\alpha_{\rm W} \simeq 1/30$, so that $T_1 \simeq M_{B_{\rm T}}/20$. Additional techni-strong annihilation channels may affect this estimate but only logarithmically in the respective cross sections (or masses).

Under the basic assumption that the DM in the present epoch $t = t_{\rm U}$ consists mostly of heavy particles of one type, e.g. T-neutrons, the condition on the current mass density of the DM,

$$\rho_{B_{\rm T}}(t_{\rm U}) = M_{B_{\rm T}} n_{B_{\rm T}}(t_{\rm U}) \simeq \rho_{\rm DM}(t_{\rm U}), \tag{21}$$

provides the canonical constraint on the thermally averaged kinetic annihilation cross section known from Ref. [59]

$$(\sigma v_B)_{\rm ann}^{\rm DM} \simeq 2.0 \times 10^{-9}\ \mathrm{GeV}^{-2}. \tag{22}$$

In terms of the latter, the relic DM abundance is

$$\Omega_{\rm DM} \simeq 0.2\Big[(\sigma v_B)_{\rm ann}^{\rm DM}/(\sigma v_B)_{\rm ann}^{\rm th}\Big]. \tag{23}$$

This formula can be used forthe determination of the T-baryon mass scale $M_{B_{\rm T}}$ as long as a theoretical prediction for the annihilation cross section $(\sigma v_B)_{\rm ann}^{\rm th}$ is given.

By comparing the above astrophysical constraint with the theoretical cross section given by Eq. (45) one extracts the lower bound for the T-baryon mass scale,

$$M_{B_{\rm T}} \gtrsim 5\ \mathrm{TeV}, \qquad g_{\rm TC}^N \gtrsim 1.0. \tag{24}$$

Clearly, this bound is consistent with naive QCD scaling hypothesis (11). An actual T-baryon mass estimate may vary in a very broad range from a few TeV to a few tens of TeV due to a strong dependence of the cross section on the techni-strong Yukawa coupling, $(\sigma v_B)_{\rm ann} \sim (g_{\rm TC}^N)^4$. Then typically the irreversible T-baryon annihilation starts at the temperature $T_0 \gtrsim$ 1–10 TeV and terminates at $T_1 \gtrsim$ 0.1–1 TeV or higher. So indeed the annihilation epoch occurs entirely in the HS phase for model I, Eq. (9), demonstrating consistency of the considered scenario of T-baryon relic formation. For model II, a high-scale $\mathrm{SU}(2)_{\rm V}$ symmetry is likely to be broken at temperatures $T < M_{B_{\rm T}}/3$ so in this case the annihilation occurs in the low-symmetry phase (low w.r.t. $\mathrm{SU}(2)_{\rm V}$) . Note that somewhat lower $M_{B_{\rm T}}$ estimates are also possible for smaller $g_{\rm TC}^N$ values in both models I and II, but then mixed EW+TC and pure EW channels become important, so that the mean value $M_{B_{\rm T}}$ never goes below 3 TeV, corresponding to 'switched off' Yukawa TC interactions, i.e. $g_{\rm TC}^N \to 0$.

Note that the estimate (24) should be taken with care since it has been obtained under a few generic conditions which can be summarized as follows:

- the DM is made entirely of heavy T-neutrons (i.e. contributions to the DM mass density from possible lighter components is negligibly small);
- the T-baryon number is conserved similarly to the baryon number;
- the T-protons decay very fast and soon after the EW phase transition epoch, and their decays do not affect the cosmological evolution of the neutrino gas and CMB;
- the neutrino gas evolution is adiabatic.

Also, we have a few model-specific assumptions:

- there is a significant splitting between the chiral symmetry breaking scale and the EW symmetry breaking scale allowing for a strong hierarchy between $M_{B_{\rm T}}$ and $M_{\rm EW}$;
- the number densities of T-baryons and anti-T-baryons are the same to a good approximation;
- the T-baryon distribution is roughly homogeneous so a possible extra loss of the DM due to its annihilation in dense regions of the Universe during structure formation epoch is neglected.

Clearly, a more involved analysis of the relic T-baryon abundance evolution, lifting out one or more of the specific assumptions would be necessary.

4.1.2 Low-symmetry phase

At lower temperatures, $T < T_{\rm EW}$, often referred to as to *the low-symmetry (LS) phase of the cosmological plasma*, the EW symmetry is broken and all the SM fermions and gauge bosons acquire non-zero masses due to non-trivial Higgs vev $v \simeq 246$ GeV. The latter scenario is conventional for low-scale SUSY-based DM models with e.g. a neutralino LSP [56]. This type of DM annihilation dynamics would happen entirely after the EW phase transition epoch for rather low mass T-baryons, i.e. $M_{B_{\rm T}} < 3\,T_{\rm EW} \sim 600$ GeV giving rise to cosmological consequences specific to the considered vector-like TC model. The vector boson masses $m_{W,Z}$, top-quark mass m_t together with $m_{\tilde\pi}$ and $M_{\tilde\sigma}$ cannot be considered as negligible compared to the T-baryon mass scale $M_{B_{\rm T}}$ any longer and have to be included making respective calculations more involved than the ones in the HS phase. Corresponding contributions to the kinetic annihilation cross section $(\sigma v_B)_{\rm ann}$ in this case are listed in Fig. 4.

Here, only T-neutrons participate in the annihilation processes. Indeed, T-protons have very small mean lifetimes (16), so they rapidly decay soon after the EW phase transition epoch and cannot substantially contribute to the T-baryon annihilation processes, and thus to DM relic abundance formation in the low-symmetry phase. The annihilation reactions shown in Fig. 4 also happen at later stages of the Universe evolution during the structure formation including the present epoch and thus are relevant for ongoing indirect DM detection measurements.

In distinction to the HS phase, there are no pure EW annihilation channels of the T-neutrons since vector boson channels may go also through intermediate h and $\tilde\sigma$ exchanges involving techni-strong Yukawa couplings and a small $\tilde\sigma h$ mixing angle θ. A straightforward calculation, however, reveals that s-channel Higgs boson and T-sigma contributions in EW annihilation channels are of the order of v_B^2, including interference terms, and can be neglected.

Now, consider the formation of the relic symmetric DM density in the LS phase of the cosmological plasma for model I set-up. It appears to be very hard to realize such a scenario with $M_{B_{\rm T}} \lesssim 600$ GeV due to a large $(\sigma v_B)_{\rm ann}^{\tilde\pi^+\tilde\pi^-}$ contribution for $g_{\rm TC}^N \gtrsim 1$ and $M_{B_{\rm T}} \gg m_{\tilde\pi}$ (see also Fig. 2). Formally, it may still be possible to tune $m_{\tilde\pi}$ and $g_{\rm TC}^N$ in a special way by minimizing the total LS cross section given by

$$(\sigma v_B)_{\rm ann}^{\rm LS} \simeq (\sigma v_B)_{\rm ann}^{\tilde\pi^+\tilde\pi^-} + (\sigma v_B)_{\rm ann}^{\tilde\pi^0\tilde\sigma} + (\sigma v_B)_{\rm ann}^{\tilde\pi^0 h}, \qquad (25)$$

but the corresponding TC model would be unnatural. In model II, there is a larger freedom in relative (rather fine) tuning between large $m_{W'}$, $m_{Z'}$ and $M_{B_{\rm T}}$ parameters capable of reducing the annihilation cross section. This, however, cannot bring $M_{B_{\rm T}}$ down significantly without spoiling EW precision constraints, and it keeps it roughly at the same level, Eq. (24). Another possibility would be to have the annihilation epoch in the mixed HS/LS phase by relaxing the constraint on the T-baryon mass scale $M_{B_{\rm T}} \lesssim 600$ GeV. The LS and HS annihilation cross sections become comparable for $M_{B_{\rm T}} \gtrsim 3$ TeV, so that in this case most of the T-baryon annihilation epoch occurs in the HS phase, again justifying the above estimate (24).

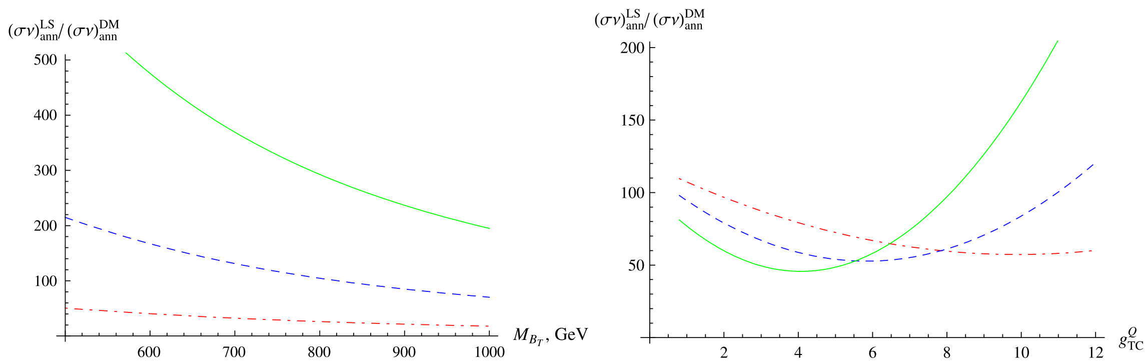

Fig. 2 Dependence of the ratio of kinetic T-baryon cross section in the LS phase for model I $(\sigma v_B)^{\rm LS}_{\rm ann}$, Eq. (25), to the kinetic cross section of the DM $(\sigma v_B)^{\rm DM}_{\rm ann} \simeq 2.0 \times 10^{-9}\ {\rm GeV}^{-2}$ required by observations on T-neutron mass $M_{B_{\rm T}}$ (*left*) and on T-quark Yukawa coupling $g^{Q}_{\rm TC}$ (*right*). In both panels we adopt the 'no $\tilde{\sigma}h$-mixing' limit $s_\theta = 0$, $M_{\tilde{\sigma}} = \sqrt{3}m_{\tilde{\pi}}$ for simplicity. In the *left panel*, the three different $g^{N}_{\rm TC} = 1.0, 1.4, 1.8$ values correspond to *dash-dotted*, *dashed*, and *solid lines*, respectively, and we have fixed $m_{\tilde{\pi}} = 250\,{\rm GeV}$ and $g^{Q}_{\rm TC} = 2.0$. In the *right panel*, $m_{\tilde{\pi}} = 150, 200, 250\,{\rm GeV}$ correspond to *dash-dotted*, *dashed*, and *solid lines*, respectively, and we have fixed $M_{B_{\rm T}} = 400\,{\rm GeV}$ and $g^{N}_{\rm TC} = 1.0$

Thus, we come to the conclusion that the most likely scenario with symmetric T-baryon DM can only be realized for the T-baryon mass scale of at least a few TeV or more, so the corresponding DM particles should be very heavy.

4.2 Scenario II: Asymmetric vector-like T-baryon Dark Matter

As was argued above, the LS annihilation cross section (25) (model I) becomes too high for low-mass T-baryons. Indeed, in Fig. 2 we show the typical parameter dependencies of the ratio $(\sigma v_B)^{\rm LS}_{\rm ann}/(\sigma v_B)^{\rm DM}_{\rm ann}$. It turns out that the LS cross section is far larger than is required by observations, at least by a factor of 10–100 or more. This means that relatively light T-neutrons and anti-T-neutrons, $M_{B_{\rm T}} \lesssim 600\,{\rm GeV}$, most probably have quickly annihilated off in the cosmological plasma by the time of their freeze-out shortly after the EW phase transition epoch. In this case, in order to provide the observable DM abundance, it is natural to assume the existence of *a T-baryon asymmetry* in a complete analogy with the typical baryon asymmetry. Thus, the bulk of the observable DM density is essentially given in terms of the T-baryon asymmetry $\Delta n_{B_{\rm T}}$, i.e.

$$\rho_{\rm DM}(t_{\rm U}) \simeq \Delta n_{B_{\rm T}} M_{B_{\rm T}}, \qquad \Delta n_{B_{\rm T}} \equiv n_{B_{\rm T}} - n_{\bar{B}_{\rm T}}, \tag{26}$$

even though it has been negligible in the beginning of the T-baryon annihilation epoch, $\Delta n_{B_{\rm T}} \ll n_{B_{\rm T}}(t_0)$. This is the so-called asymmetric DM (ADM) model, which has been previously studied in other TC/compositeness scenarios (see e.g. Refs. [60–62]), and it can be realized in the considered vector-like TC model as well.

Introducing a fractional asymmetry of T-neutrons N and anti-T-neutrons $\bar{N}$ in the cosmological plasma by

$$r \equiv \frac{n(\bar{N})}{n(N)}, \qquad 0 < r < 1, \tag{27}$$

one could estimate its time evolution via a detailed analysis of the system of coupled Boltzmann equations, as performed in Refs. [63,64]. The late-time fractional asymmetry r_∞ for light T-neutrons turns out to be exponentially suppressed in most of the vector-like TC parameter space, i.e.

$$r_\infty \sim \exp\left\{-2\frac{(\sigma v_B)^{\rm LS}_{\rm ann}}{(\sigma v_B)^{\rm DM}_{\rm ann}}\right\} \ll 1 \tag{28}$$

for the dominant s-wave Dirac T-neutron annihilation processes considered above in the LS annihilation (light T-neutron, model I). Having typically large ratios of the cross sections illustrated in Fig. 2 one concludes that practically no light anti-T-neutrons remain in the modern Universe similar to ordinary antibaryons. The T-neutrons by themselves cannot produce any annihilation-like signal expected to be constrained by indirect detection measurements, while they may have a certain impact on direct measurements [65].

Similarly to the HS phase annihilation, in the vector-like T-neutron ADM scenario the T-baryon mass scale is expected to be above the EW scale $M_{B_{\rm T}} \gtrsim 200\,{\rm GeV}$, so that the chemical decoupling of the ADM occurs when DM is non-relativistic, while SM fermions are still relativistic. Then one could assume a tight relation between ordinary baryon and T-baryon asymmetries as typically considered in ADM scenarios, which translates into a relation between the number densities of the visible matter and T-neutron abundances. In this case depending on the details of the chemical equilibrium the T-baryon mass scale is expected to be rather high, $M_{B_{\rm T}} \sim {\rm TeV}$, as was advocated in Ref. [66], which is consistent with the suggested vector-like ADM scenario having the T-neutron annihilation epoch in the LS (or mixed HS+LS) phase.

For relatively low $g^{N}_{\rm TC} \sim 1$ and high T-neutron masses $M_{B_{\rm T}} \gtrsim 1\,{\rm TeV}$ the ratio of the cross sections goes down,

$$1 < \frac{(\sigma v_B)^{\rm LS}_{\rm ann}}{(\sigma v_B)^{\rm DM}_{\rm ann}} \lesssim 10, \tag{29}$$

and it can therefore accommodate the *partially ADM scenario* $0 < r_\infty < 1$ with heavy T-neutrons allowing for many attractive features. The same is true for the case of HS annihilation in model I and LS annihilation in model II. In particular, having a small but non-zero relic density of anti-T-neutrons would open up immediate opportunities for indirect detection measurements of DM annihilation products from galactic cores and compact stars. Also, this scenario is a particular case of the self-interacting DM model, which allows one to avoid any problematic cusp-like DM density profile in the central regions of galactic haloes, leading to the core-like DM distribution favored by astrophysical observations [67]. In the considered scenario, the intensive annihilation together with elastic NN scattering plays an important role in the self-interactions of the DM particles, allowing for necessary adjustments of the DM density profile. Indeed, in the epoch of early structure formation an overdensity of DM in cusp-like regions has been eliminated by intensive annihilation processes, so that they do not exist today.

Another interesting astrophysical implication of the partially ADM scenario in general is the possible thermalization of the cosmological medium in the beginning of the structure formation epoch. Indeed, the growth of structures is accompanied by a substantial increase in DM density in the central regions of haloes. Having a large cross section (25) the intensity of $N\bar{N}$ annihilation processes has gone up in that epoch and further reduced the amount of anti-T-neutrons which have survived after the T-neutron freeze-out. The annihilation products could be capable of thermalization of the medium at $z \sim 10$ or somewhat earlier which may have serious observational consequences. Some small remnants of anti-T-neutron density could have survived such a 'second T-baryon annihilation epoch' and could remain today providing possible observational signatures of their annihilation with T-neutrons. Certainly, a thorough analysis involving simulations of the structure formation epoch together with T-baryon annihilation and DM formation details is required.

Finally, intensive vector-like T-neutron DM annihilation allows one to explain why the DM in the Galactic halo is much more tepid than ordinary CDM models predict. The corresponding problem is typically tagged as the missing satellite problem. Due to a much higher temperature of the DM in the Galactic halo, the observed number of dwarf galaxies is by an order of magnitude smaller, while the DM density in the halo cores is much smaller than CDM WIMP-based simulations predict [68–70]. In the considered vector-like TC model in the case of partial ADM one finds an interesting opportunity for such a 'tepid' DM. Indeed, at the initial stages of structure formation slower (colder) DM particles in the central cusp-like regions have annihilated off, while faster particles moving in less dense regions at the Galactic periphery according to the Jeans instability criterion could have survived until today. So this proposal could be an efficient alternative to warm DM models with low-mass WIMPs $\mathcal{O}(1)$ GeV aimed at resolving the above issues.

Thus, the vector-like ADM scenario with a relatively high annihilation rate of heavy T-neutrons offers a few appealing possibilities compared to traditional practically non-interacting WIMP-miracle. The basic opportunities mentioned above inherent to the considered T-neutron DM scenario should further be explored quantitatively.

5 Direct T-neutron detection constraints

Finally, let us consider one of the most important constraints on the vector-like TC model with Dirac T-baryons—the direct DM detection limits. Among them the data on the spin-independent (SI) component of the elastic WIMP-nucleon cross section from CDMS II [71] and, especially, XENON100 [49] experiments provide the most stringent model-independent exclusion limits. Indeed, the elastic T-neutron–nucleon scattering goes via a T-neutron vector coupling to the Z boson defined in Eq. (8). This leads to a sizable SI scattering cross section off nuclei, which needs to be compared to the data.

Let us consider Dirac T-neutron–nucleon scattering in the non-relativistic limit $v_N \ll 1$. Previously, a similar process has been investigated in the case of Dirac vector-like neutralino DM (see e.g. Refs. [50,51,72,73]), so we do not go into details of explicit calculations here. Following Ref. [73] the SI spin-averaged T-neutron–nucleus cross section reads

$$\sigma_{\rm SI} = \frac{\mu^2}{16\pi M_{B_{\rm T}}^2 m_A^2}\left(\frac{1}{4}\sum_{\rm spins}|\mathcal{M}_{\rm SI}|^2\right), \tag{30}$$

where μ is the reduced mass and m_A is the mass of the target nucleus. In the considered case at small momentum transfers $q^2 \ll m_Z^2$ the effective operator for T-neutron scattering off quarks through Z-exchange is described by vector couplings only,

$$\mathcal{O}_Z^q = \delta_Z\,\lambda_{qV}\,\frac{ig_2^2}{2c_W^2 m_Z^2}\,\bar{N}\gamma^\mu N\cdot\bar{q}\gamma_\mu q, \tag{31}$$

where $q = u, d$ are quarks in a nucleon, and the standard vector Zq couplings are

$$\lambda_{qV} = \frac{g_2}{c_W}\left[t_{qL}^3 - 2Q_q s_W^2\right]. \tag{32}$$

This leads to a squared matrix element (cf. Ref. [73])

$$\frac{1}{4}\sum_{\text{spins}}|\mathcal{M}_{\text{SI}}|^2 \simeq 16\,\frac{M_{B_{\text{T}}}^2 m_A^2}{m_Z^4}\,\frac{g_2^4}{4c_W^4}\,\delta_Z^2 \times \left[\sum_{q=u,d}\lambda_{qV}[Z B_{qV}^p + (A-Z)B_{qV}^n]\right]^2, \tag{33}$$

where $B_{uV}^p = B_{dV}^n = 2$ and $B_{uV}^n = B_{dV}^p = 1$ are u,d quark multiplicities in a proton, p and neutron, n, and Z and A correspond to the atomic number and atomic mass of the target nucleus, respectively. For the individual elastic N–p and N–n cross sections one obtains

$$\sigma_{\text{SI}}^{N-p,n} = \delta_Z^2\,\frac{g_2^4}{16\pi c_W^4}\,\frac{\mu^2}{m_Z^4}\,\left(c_V^{p,n}\right)^2, \qquad c_V^p = 1-4c_W^2, \quad c_V^n = -1, \tag{34}$$

where $c_V^{p,n}$ are the vector form factors of the proton and neutron. The elastic T-neutron–nucleus cross section reads

$$\sigma_{\text{SI}}^{N-A} = \delta_Z^2\,\frac{g_2^4}{16\pi c_W^4}\,\frac{\mu^2}{m_Z^4}\left[Z\,c_V^p + (A-Z)\,c_V^n\right]^2. \tag{35}$$

For practical use, it is instructive to represent the T-neutron–nucleon cross sections in numerical form, i.e.

$$\sigma_{\text{SI}}^{N-p} = 1.5\times 10^{-40}\,\text{cm}^2\;\times\delta_Z^2\left(\frac{\mu}{m_p}\right)^2, \qquad \sigma_{\text{SI}}^{N-n} = 2.5\times 10^{-38}\,\text{cm}^2\;\times\delta_Z^2\left(\frac{\mu}{m_n}\right)^2. \tag{36}$$

These cross sections are rather large and strongly constrained by direct detection experiments. At the moment, XENON100 [49] provides the most stringent limit on $\sigma_{\text{SI}}^{\text{nucleon}}$ (per nucleon in the case of a xenon target) for high mass Dirac T-neutrons, which is roughly

$$-\log_{10}\left(\frac{\sigma_{\text{SI}}^{\text{nucleon}}}{\text{cm}^2}\right) \simeq 44.6 - 43.4, \tag{37}$$

corresponding to a T-neutron mass range of about

$$M_{B_{\text{T}}} \simeq 0.1 - 2\,\text{TeV}, \tag{38}$$

respectively. This limit immediately provides a strong bound on Dirac vector-like T-neutron coupling to the Z-boson, namely,

$$\delta_Z \lesssim 2\times 10^{-3}, \qquad M_{B_{\text{T}}} \lesssim 2\,\text{TeV}, \tag{39}$$

and a little bit weaker constraint for $M_{B_{\text{T}}} \gtrsim 2\,\text{TeV}$. This means that if the corresponding direct DM detection limits are confirmed the vector-like Dirac T-baryons with standard EW interactions, i.e. model I with $\delta_Z \simeq 1$ introduced above in Eq. (9), are firmly ruled out. Note, however, that if the T-baryon mass scale is $M_{B_{\text{T}}} < 5\,\text{TeV}$ for $g_{\text{TC}}^N \gtrsim 1.0$ and if T-baryon symmetry is exact in Nature, Dirac T-neutrons have all annihilated off by their freeze-out leaving practically none of them in the present Universe—only under this condition odd $\text{SU}(2n+1)_{\text{TC}}$ confined symmetries would not be excluded by the direct DM detection, while they are excluded for the asymmetric DM case in a wide range of $M_{B_{\text{T}}}$ values and for symmetric DM in the case of larger scales $M_{B_{\text{T}}} \gtrsim 5\,\text{TeV}$.

The other appealing possibility to accommodate the Dirac T-baryons is in model II with $\delta_Z \ll 1$, which, however, requires the extra assumption of the existence of an extra vector $\text{SU}(2)_{\text{V}} \neq \text{SU}(2)_{\text{W}}$ gauge symmetry in the constituent T-quark/T-baryon sector. According to the standard GLσM approach, the local vector subgroup $\text{SU}(2)_{\text{V}}$ of the chiral group can be effectively associated with the heavy composite vector T-rho sector $\tilde{\rho}^{0,\pm}$ treated as 'gauge bosons' of the local chiral group. Model II with the $\delta_Z \ll 1$ condition can then be naturally motivated by a possibly much stronger coupling of T-baryons to the heavy vector T-mesons compared to the weak coupling to ordinary Z, $W^{\pm}$ bosons. Since the vector-like TC model in its present minimal form does not incorporate heavy (pseudo)vector T-mesons, the latter opportunity offered by model II for QCD-like theories with Dirac T-neutrons remains a promising direction for further studies.

Note that the constraint (39) is the upper limit on the mixing parameter (e.g. the sine of the mixing angle) between Z and new Z' bosons from an unspecified high-scale $\text{SU}(2)_{\text{V}}$, which should also be additionally constrained by EW precision tests and extra gauge boson (Z' and W') searches at the LHC. The latter analysis can therefore be performed together with the direct DM detection limits which, however, goes beyond the present scope.

6 Summary and conclusions

In this work we have investigated basic properties of the Dirac vector-like T-neutron DM predicted by the vector-like technicolor model [34] and have found important limitations on the structure of techni-strong dynamics from the direct DM detection data (in particular, by the XENON100 experiment [49]). This has been done in the simplest QCD-like setting of the T-confinement $\text{SU}(3)_{\text{TC}}$ group and one generation of T-quarks. We have shown that under the natural assumption of T-baryon number conservation, the local chiral symmetry breaking gives rise to the vector $\text{SU}(2)_{\text{V}}$ gauge symmetry, which acts on constituent T-quark and T-baryon sectors, additional to the SM sectors. Whether the $\text{SU}(2)_{\text{V}}$ group is identified with the weak isospin $\text{SU}(2)_{\text{W}}$ symmetry of the SM or not provides us with the two possible scenarios for the structure of (vector) weak interactions in the T-quark/T-baryon sectors, respectively.

As a specific prediction of the vector-like TC model, the constituent T-quark and hence the T-baryon mass spectrum is degenerate at tree level by virtue of vector-like weak interactions. The EW radiative corrections with vector T-baryon–Z, γ couplings effectively split the T-baryon sector making the T-proton $P = (\text{UUD})$ (slightly) heavier than the T-neutron $N = (\text{UDD})$, irrespectively of the nature of $\text{SU}(2)_\text{V}$ group. The lightest T-baryon state is then the T-neutron, which provides us with the prominent heavy self-interacting DM candidate with many appealing features. Two scenarios with symmetric and partially asymmetric T-neutron DM have been considered and limits on the T-baryon mass scale have been derived from the DM relic abundance data. Together with the naive QCD scaling hypothesis, this provides an effective lower bound on the new T-confinement scale in a few TeV range. We conclude that the sufficient amount of T-neutron DM can only be formed in the high-symmetry phase of the cosmological plasma at temperatures above the EW symmetry breaking scale.

As was discussed thoroughly in Ref. [34], the EW precision constraints at the fundamental level can only be satisfied for the vector-like T-quarks under the weak isospin $\text{SU}(2)_\text{W}$. Alternatively, one could introduce the vector-like weak interactions via a small mixing $\delta_Z \ll 1$, which can be limited so as not to upset the SM tests. While both scenarios can be satisfied by all existing EW and collider constraints, it turned out that only the latter scenario with QCD-type TC group $\text{SU}(3)_\text{TC}$ can be consistent with the DM astrophysics constraints. This provides an extra important constraint on the structure of TC sectors and high-scale strongly coupled dynamics.

Indeed, the Dirac T-baryons originated under the simplest assignment of an odd T-confined group in the T-quark sector having rank three, i.e. $\text{SU}(3)_\text{TC}$. Thus, we conclude that in this case or, more generally, in the case of any odd group $\text{SU}(2n+1)_\text{TC}$, $n = 1, 2, \ldots$, it is not possible to introduce the standard EW interactions over the SM $\text{SU}(2)_\text{W}$ gauge group in a phenomenologically consistent way as suggested by the XENON100 constraint (39). The only way to satisfy the existing phenomenological (EW, collider, and astrophysics) constraints is to consider an even T-confinement $\text{SU}(2n)_\text{TC}$, $n = 1, 2, \ldots$ group, for example, the simplest $\text{SU}(2)_\text{TC}$, where the lightest stable neutral T-baryon state $B_0 = (\text{UD})$ is scalar and does not interact with the Z-boson, thus evading the direct detection limits. The corresponding analysis is ongoing.

Note that $\text{SU}(2n+1)_\text{TC}$ confined groups with EW vector-like T-fermion (constituent T-quark and T-baryon) sectors are excluded by direct DM detection if and only if the T-confinement scale is well above a TeV scale and/or there existed a mechanism for T-baryon asymmetry generation in the early Universe, in analogy to that of the baryon asymmetry generation (i.e. the Sakharov conditions were at work for T-baryons as well as for baryons). In the opposite case, if the T-baryon symmetry is exact in Nature and the T-baryon mass scale is not very large, $\lesssim$5 TeV, then all T-baryons would have annihilated off by the freeze-out time due to rather large annihilation cross sections (into T-pions mostly). Modern DM would then be made of something else, and the direct detection constraints would not be decisive for the considered model. In the latter case, the vector-like $\text{SU}(3)_\text{TC}$ model would not have any direct astrophysical implications in the late Universe.

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Appendix A: Annihilation of vector-like Dirac T-baryons

Assuming that the DM consists of composite T-baryons (mostly T-neutrons with probably a small fraction of anti-T-neutrons), let us discuss extra possible constraints on the vector-like TC parameter space coming from astrophysical observations and cosmological evolution of the DM (for a review of the current (in)direct DM detection measurements and constraints, see e.g. Refs. [52–55]).

In order to estimate the T-baryon mass scale M_{B_T} from the DM relic abundance data [1] one has to consider the evolution of the T-baryon density in the early Universe, which is largely determined by the kinetic $B_\text{T}\bar{B}_\text{T}$ annihilation cross section $(\sigma v_B)_\text{ann}$. As is typical for the cold DM formation scenarios one naturally assumes that the residual T-baryon abundance is formed at temperatures $T \ll M_{B_\text{T}}$ when a non-relativistic approximation is applied.

It is reasonable to consider the irreversible annihilation of T-baryons in two different phases of the cosmological plasma separately—before and after the EW phase transition epoch, $T_\text{EW} \sim 200$ GeV. Consequently, we will end up with two different scenarios for the DM relic abundance formation.

A.1 Annihilation of vector-like T-baryons: the high-symmetry phase

Let us evaluate the vector-like T-baryon annihilation cross section, $(\sigma v_B)_\text{ann}$, in the high-symmetry (HS) phase of the cosmological plasma in model I, Eq. (9). All relevant contributions are schematically depicted in Fig. 3. In comparison with the Higgsino LSP scenario in the SU(5) split SUSY model [50,51], the T-baryon annihilation in the HS phase is

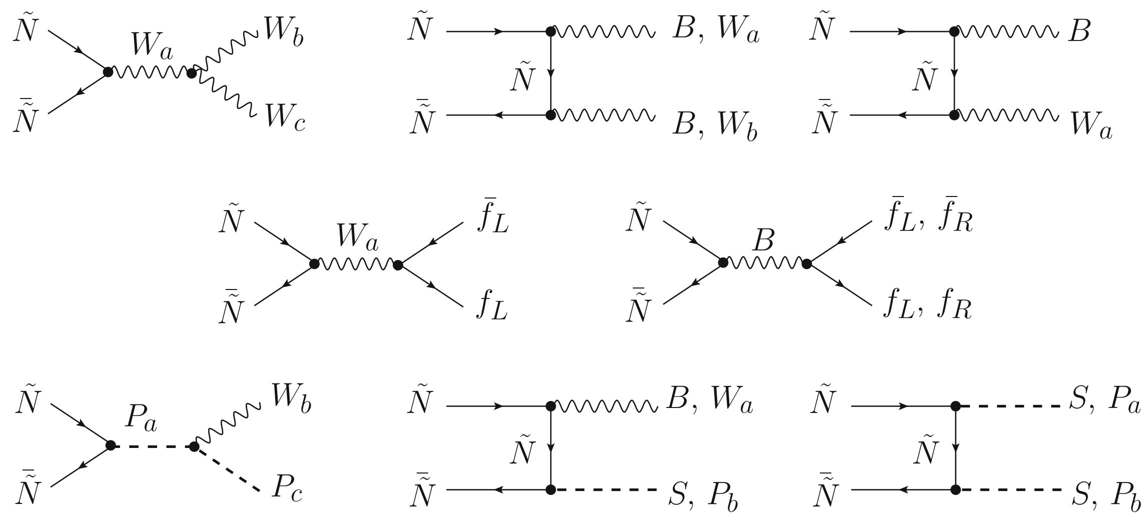

Fig. 3 Typical diagrams contributing to the T-baryon DM annihilation in the high-symmetry phase of the cosmological plasma corresponding to $T_{\rm EW} < T \lesssim u$. Model I, Eq. (9), is implied

given by essentially the same EW amplitudes, due to the same vector-like structure of T-baryon and Higgsino EW interactions, i.e.

$$\begin{aligned}&\tilde{N}\bar{\tilde{N}} \to BB, \quad \tilde{N}\bar{\tilde{N}} \to BW_a, \quad \tilde{N}\bar{\tilde{N}} \to W_aW_b,\\ &\tilde{N}\bar{\tilde{N}} \to B^* \to l_L\bar{l}_L,\ q_L\bar{q}_L,\ e_R\bar{e}_R,\ u_R\bar{u}_R,\ d_R\bar{d}_R, \qquad (40)\\ &\tilde{N}\bar{\tilde{N}} \to W^*_a \to l_L\bar{l}_L,\ q_L\bar{q}_L,\end{aligned}$$

where $l_L,\ q_L$, and $e_R,\ u_R,\ d_R$ are the $SU(2)_W$ doublet and singlet (chiral) leptons and quarks, respectively, in each of the three generations. The corresponding EW contribution to the total T-baryon annihilation cross section in the HS phase for non-relativistic T-baryons $v_B \ll 1$ is found to be

$$(\sigma v_B)^{\rm EW}_{\rm ann} = \frac{21g_1^4 + 6g_1^2g_2^2 + 39g_2^4}{512\pi M^2_{B_{\rm T}}}. \qquad (41)$$

Here $g_1 = g_1(\sqrt{s}),\ g_2 = g_2(\sqrt{s})$ are the EW gauge couplings fixed at the scale $\sqrt{s} \simeq 2M_{B_{\rm T}}$.

In addition to the pure EW channels listed above, there are a few important T-strong channels with primary T-pion P_a and T-sigma S in the intermediate and final states. In particular, the annihilation channels into a (pseudo)scalar and a massless gauge boson involving additional Yukawa interactions in the T-hadron sector are

$$\tilde{N}\bar{\tilde{N}} \to P^*_a \to P_bW_c, \qquad \tilde{N}\bar{\tilde{N}} \to P_aB,\ SB,\ SW_a, \qquad (42)$$

and the corresponding total cross section in the limit $M_{B_{\rm T}} \gg M_{\tilde{\sigma}},\ m_{\tilde{\pi}}$ is

$$(\sigma v_B)^{\rm EW+TC}_{\rm ann} \simeq \frac{\left(g^N_{\rm TC}\right)^2\left(2g_1^2 + 3g_2^2\right)}{32\pi M^2_{B_{\rm T}}}. \qquad (43)$$

In order to turn to model II, Eq. (9), in the limit $m_{Z'} \ll M_{B_{\rm T}}$, corresponding to unbroken $SU(2)_{\rm V} \neq SU(2)_{\rm W}$, one has to perform the replacement $g_2 \to g_2^{\rm V}$ in Eqs. (41) and (43). This would provide a rough estimate for the annihilation cross sections into $B_\mu,\ V^a_\mu$ bosons. For a more precise analysis of the gauge $SU(2)_{\rm V}$ part of the cross sections one should consider details of the broken phase of $SU(2)_{\rm V}$ and evaluate them for massive $Z',\ W'^{\pm}$ bosons for kinematically allowed channels, i.e. for $m_{W',Z'} \lesssim M_{B_{\rm T}}/2$ (for more details, see the calculations in the low-symmetry phase below). The latter, however, do not affect our conclusions here since the corresponding cross sections are relatively small compared to those in the T-strong channels.

For pure T-strong channels

$$\tilde{N}\bar{\tilde{N}} \to P_aP_b,\ SP_a,\ {\rm SS}, \qquad (44)$$

we have the total cross section

$$(\sigma v_B)^{\rm TC}_{\rm ann} \simeq \frac{9\left(g^N_{\rm TC}\right)^4}{32\pi M^2_{B_{\rm T}}}. \qquad (45)$$

The latter comes essentially from the T-pion channels P_aP_b and SP_a, while the T-sigma one is suppressed by the relative velocity squared, i.e. $(\sigma v_B)^{\rm SS}_{\rm ann} \sim v_B^2$.

Based upon a QCD analogy the T-strong Yukawa interactions are much more intensive than the EW interactions, i.e. $g^N_{\rm TC} \gtrsim 1,\ g^N_{\rm TC} \gg g_{1,2},\ g_2^{\rm V}$, leading to strong dominance of pure TC (T-pion induced) annihilation channels in both models I and II, so that

$$(\sigma v_B)_{\rm ann} \simeq (\sigma v_B)^{\rm TC}_{\rm ann}. \qquad (46)$$

This makes the considered T-baryon DM model specific compared to other standard SUSY-based DM models where $(\sigma v_B)_{\rm ann} \sim \alpha_{\rm W}^2/M_\chi^2,\ \alpha_{\rm W} \simeq 1/30$, given by weak interactions only. Thus, more intense T-baryon annihilation with extremely weak interactions with ordinary matter (see below)

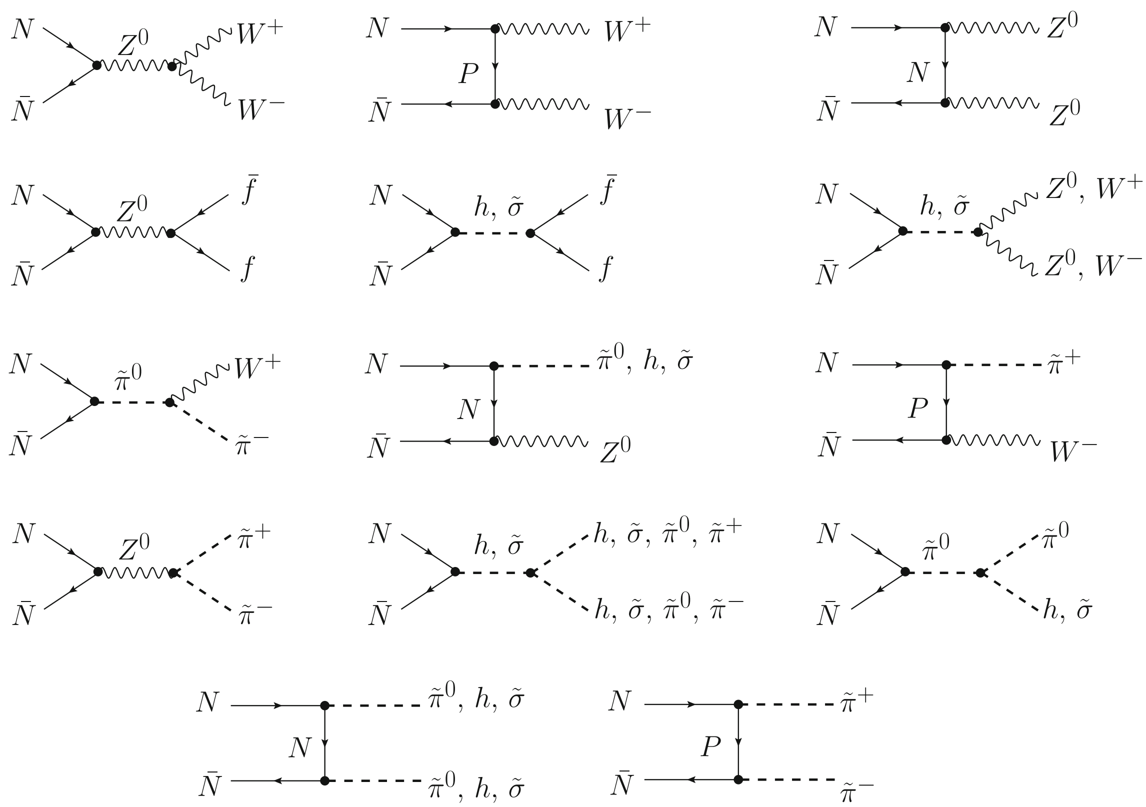

Fig. 4 Typical diagrams contributing to the T-neutron DM annihilation in the low-symmetry phase of the cosmological plasma $T < T_{\rm EW}$ including much later stages of structure formation and the present epoch. Model I, Eq. (9), is implied

makes them promising DM candidates, alternative to standard WIMPs. Note that the value (45) behaves as a fourth power of $g_{\rm TC}^N$ leading to a large sensitivity of the T-baryon mass scale extracted from the DM relic abundance data to this parameter (see below).

A.2 Annihilation of vector-like T-baryons: the low-symmetry phase

Consider now the T-baryon annihilation in the low-symmetry phase of the cosmological plasma in the phenomenologically appealing model I, Eq. (9), in detail. The EW channels with SM fermion and gauge boson final states shown in the first two lines in Fig. 4 are

$$N\bar{N} \to Z^{0*}, h^*, \tilde{\sigma}^* \to l\bar{l},\ q\bar{q}, \quad N\bar{N} \to W^+W^-,\ Z^0Z^0, \tag{47}$$

respectively. The annihilation cross section into fermions is

$$(\sigma v_B)_{\rm ann}^{f\bar{f}} \simeq \frac{1}{192\pi M_{B_{\rm T}}(4M_{B_{\rm T}}^2 - m_Z^2)^2}\Big\{2M_{B_{\rm T}}^3(103g_1^4 + 6g_1^2g_2^2 + 63g_2^4) + (17g_1^4 - 6g_1^2g_2^2 + 9g_2^4)(2M_{B_{\rm T}}^2 + m_t^2)\sqrt{M_{B_{\rm T}}^2 - m_t^2}\Big\}, \tag{48}$$

into a W^+W^- pair,

$$(\sigma v_B)_{\rm ann}^{W^+W^-} \simeq \frac{g_2^4}{64\pi M_{B_{\rm T}} m_W^4}\frac{(M_{B_{\rm T}}^2 - m_W^2)^{3/2}}{(2M_{B_{\rm T}}^2 - m_W^2)^2(4M_{B_{\rm T}}^2 - m_Z^2)^2} \times\Big\{4M_{B_{\rm T}}^4(12m_W^4 - 4m_W^2m_Z^2 + m_Z^4) + 4M_{B_{\rm T}}^2m_W^2(20m_W^4 - 24m_W^2m_Z^2 + 5m_Z^4) + m_W^4(12m_W^4 - 12m_W^2m_Z^2 + 5m_Z^4)\Big\}, \tag{49}$$

and into a Z^0Z^0 pair

$$(\sigma v_B)_{\rm ann}^{Z^0Z^0} \simeq \frac{g_2^4(M_{B_{\rm T}}^2 - m_Z^2)^{3/2}}{64\pi c_W^4 M_{B_{\rm T}}(2M_{B_{\rm T}}^2 - m_Z^2)^2}. \tag{50}$$

Above we have neglected all the fermion masses, assuming for simplicity $m_{l,q} \ll M_{B_{\rm T}}$, except for the top-quark mass: $m_t \simeq 173\,{\rm GeV}$. Also, the vector boson and (pseudo)scalar masses cannot be neglected and are kept here. The $N\bar{N} \to h^*, \tilde{\sigma}^* \to l\bar{l},\ q\bar{q},\ W^+W^-$ processes and their interference with pure EW channels are of the order of $\sim v_B^2$. So the $\tilde{\sigma}$- and h-mediated diagrams were neglected in the integrated cross sections (48), (49), and (50).

The channels with mixed (gauge and (pseudo)scalar) final states (third line in Fig. 4) are

$$N\bar{N} \to W^\pm\tilde{\pi}^\mp, \quad N\bar{N} \to Z^0\tilde{\pi}^0, \quad N\bar{N} \to Z^0\tilde{\sigma}, \quad N\bar{N} \to Z^0h, \tag{51}$$

leading to the following contributions:

$$(\sigma v_B)_{\rm ann}^{W^\pm\tilde\pi^\mp} \simeq \frac{g_2^2(g_{\rm TC}^N)^2\Big[16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(m_{\tilde\pi}^2+m_W^2)+(m_{\tilde\pi}^2-m_W^2)^2\Big]^{3/2}}{32\pi M_{B_{\rm T}}^4(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2)^2(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2-m_W^2)^2} \times\Big\{(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2)^2+2M_{B_{\rm T}}^2m_W^2\Big\}, \quad (52)$$

$$(\sigma v_B)_{\rm ann}^{Z^0\tilde\pi^0} \simeq \frac{g_2^2(g_{\rm TC}^N)^2\Big[16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(m_{\tilde\pi}^2+m_Z^2)+(m_{\tilde\pi}^2-m_Z^2)^2\Big]^{3/2}}{128\pi c_W^2 M_{B_{\rm T}}^4(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2-m_Z^2)^2}, \quad (53)$$

$$(\sigma v_B)_{\rm ann}^{Z^0\tilde\sigma} \simeq \frac{g_2^2(g_{\rm TC}^N)^2\,c_\theta^2\sqrt{16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(M_{\tilde\sigma}^2+m_Z^2)+(M_{\tilde\sigma}^2-m_Z^2)^2}}{128\pi c_W^2 M_{B_{\rm T}}^4(4M_{B_{\rm T}}^2-M_{\tilde\sigma}^2-m_Z^2)^2} \times\Big\{16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(M_{\tilde\sigma}^2-2m_Z^2)+(M_{\tilde\sigma}^2-m_Z^2)^2\Big\}, \quad (54)$$

$$(\sigma v_B)_{\rm ann}^{Z^0h} \simeq \frac{g_2^2(g_{\rm TC}^N)^2\,s_\theta^2\sqrt{16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(m_h^2+m_Z^2)+(m_h^2-m_Z^2)^2}}{128\pi c_W^2 M_{B_{\rm T}}^4(4M_{B_{\rm T}}^2-m_h^2-m_Z^2)^2} \times\Big\{16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(m_h^2-2m_Z^2)+(m_h^2-m_Z^2)^2\Big\}, \quad (55)$$

where $m_h \simeq 126$ GeV is the Higgs boson mass [57,58].

Finally, the last two lines in Fig. 4 represent diagrams with T-strong final states (T-pions and T-sigma) as well as the Higgs boson:

$$N\bar N \to \tilde\pi^+\tilde\pi^-,\quad N\bar N\to\tilde\pi^0\tilde\pi^0,\quad N\bar N\to\tilde\pi^0\tilde\sigma,\quad N\bar N\to\tilde\pi^0 h,\quad N\bar N\to h\tilde\sigma,\quad N\bar N\to\tilde\sigma\tilde\sigma,\quad N\bar N\to hh. \quad (56)$$

The relevant contributions to the total cross section are

$$(\sigma v_B)_{\rm ann}^{\tilde\pi^+\tilde\pi^-} \simeq \frac{(M_{B_{\rm T}}^2-m_{\tilde\pi}^2)^{3/2}}{16\pi M_{B_{\rm T}}}\left(\frac{g_2^4}{(4M_{B_{\rm T}}^2-m_Z^2)^2}+\frac{4(g_{\rm TC}^N)^4}{(2M_{B_{\rm T}}^2-m_{\tilde\pi}^2)^2}\right), \quad (57)$$

$$(\sigma v_B)_{\rm ann}^{\tilde\pi^0h} \simeq \frac{(g_{\rm TC}^N)^2\,s_\theta^2\sqrt{16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(m_h^2+m_{\tilde\pi}^2)+(m_h^2-m_{\tilde\pi}^2)^2}}{256\pi M_{B_{\rm T}}^4M_Q^2(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2)^2(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2-m_h^2)^2} \times\Big(g_{\rm TC}^QM_{B_{\rm T}}(m_{\tilde\pi}^2-m_h^2)(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2-m_h^2) +2g_{\rm TC}^NM_Q(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2)(4M_{B_{\rm T}}^2+m_{\tilde\pi}^2-m_h^2)\Big)^2, \quad (58)$$

$$(\sigma v_B)_{\rm ann}^{\tilde\pi^0\tilde\sigma} \simeq \frac{(g_{\rm TC}^N)^2\,c_\theta^2\sqrt{16M_{B_{\rm T}}^4-8M_{B_{\rm T}}^2(M_{\tilde\sigma}^2+m_{\tilde\pi}^2)+(M_{\tilde\sigma}^2-m_{\tilde\pi}^2)^2}}{256\pi M_{B_{\rm T}}^4M_Q^2(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2)^2(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2-M_{\tilde\sigma}^2)^2} \times\Big(g_{\rm TC}^QM_{B_{\rm T}}(m_{\tilde\pi}^2-M_{\tilde\sigma}^2)(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2-M_{\tilde\sigma}^2) +2g_{\rm TC}^NM_Q(4M_{B_{\rm T}}^2-m_{\tilde\pi}^2)(4M_{B_{\rm T}}^2+m_{\tilde\pi}^2-M_{\tilde\sigma}^2)\Big)^2, \quad (59)$$

while scalar ($\tilde\sigma\tilde\sigma$, hh, and $\tilde\sigma\tilde\sigma$) and pseudoscalar $\tilde\pi^0\tilde\pi^0$ channels are suppressed as $\sim v_B^2$, due to a cancelation between t- and u-channel diagrams. This is similar to the HS phase, where the T-pion and mixed T-pion/T-sigma channels dominate in the total cross section. In order to turn to model II with broken $\mathrm{SU(2)_V} \neq \mathrm{SU(2)_W}$ symmetry, one has to make the following replacements: $m_Z \to m_{Z'}$, $m_W \to m_{W'}$, $g_2 \to g_2^{\rm V}$ in the above formulas (48)–(59).

In the above expressions, $M_{B_{\rm T}} \simeq 3M_Q$, $m_{\tilde\pi}$, $M_{\tilde\sigma}$, and $g_{\rm TC}^{N,Q}$ are kept as free parameters to be constrained from (collider and astrophysics) phenomenology. Due to the rather strong inequality $g_{\rm TC}^{N,Q} \gg g_{1,2}, g_2^{\rm V}$ characteristic for the new strongly coupled dynamics under discussion, the $\tilde\pi^+\tilde\pi^-$-channel $(\sigma v_B)_{\rm ann}^{\tilde\pi^+\tilde\pi^-}$ and mixed scalar–pseudoscalar channels $(\sigma v_B)_{\rm ann}^{\tilde\pi^0h}$ and $(\sigma v_B)_{\rm ann}^{\tilde\pi^0\tilde\sigma}$ dominate the total T-baryon annihilation cross section in the LS phase for not very large $M_{B_{\rm T}} \lesssim 600$ GeV. Note that all the annihilation cross sections in the HS and LS phases behave as $(\sigma v_B)_{\rm ann} \sim M_{B_{\rm T}}^{-2}$ in the limit of large $M_{B_{\rm T}}$.

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