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

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Abstract

We continue to explore a question about the existence of a new strongly coupled dynamics above the electroweak scale. The latter has been recently realized in the simplest consistent scenario, the vector-like (or chiral-symmetric) Technicolor model based upon the gauged linear σ-model. One of the predictions of a new strong dynamics in this model, the existence of stable vector-like technibaryon states at a TeV scale, such that the lightest neutral one could serve as a Dark Matter candidate. Here, we consider the QCD-type Technicolor with $SU(3)_{\rm TC}$ confined group and one $SU(2)_{\rm W}$ doublet of vector-like techniquarks and test this model against existing Dark Matter astrophysics data. We show that the spin-independent Dirac technineutron-nucleon cross section is by far too large and ruled out by XENON100 data. We conclude that vector-like techniquark sectors with an odd group of confinement $SU(2n+1)_{\rm TC}$, $n = 1, 2, \dots$ and with ordinary vector-like weak $SU(2)_{\rm W}$ interactions are excluded if the technibaryon number is conserved. We discuss a possible generic TC scenario with technibaryon sector interacting via an extra vector $SU(2)_{\rm V}$ other than the standard weak $SU(2)_{\rm W}$ and consider immediate implications for the cosmological evolution and freeze out of heavy relic technineutrons.

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I. INTRODUCTION

The undoubtful existence of the Dark Matter (DM) comprising about a third (or more precisely, about 27 % [1]) of 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), 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, lightest supersymmetric particles (LSPs) predicted by supersymmetry (SUSY) [3] such as neutralino are often referred to as to the best DM candidates [4], 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. [5, 6] 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 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. Ref. [7]).

In this paper, we consider 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 $SU(3)_{\rm TC}$ 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 mid-eighties from Ref. [8] 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 has been proposed. Generic DM signatures from Technicolor-based models with stable T-baryons were discussed e.g. in Refs. [9–11] (for a review see also Ref. [12] 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 [13, 14]. 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 [15]. Having similar mechanisms for ordinary matter and DM formation in 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 dynamical EWSB mechanism such composite DM candidates may be self-interacting which helps in avoiding problematic cusp-like DM halo profiles [16].

All of the existing composite DM models rely on the basic assumption about New Physics extension of the SM by means of extra confined matter sectors. These ideas were realized in a multitude of Technicolor (TC) models developed so far [17] (for a detailed review on the existing TC models, see e.g. Refs. [18, 19]). 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 $SU(2)_{\rm L} \otimes SU(2)_{\rm R} \to SU(2)_{\rm V}$ 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, $\langle \tilde{Q}\tilde{\bar{Q}} \rangle \neq 0$. 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 [20]. However, many of the existing TC-based models have got severely constrained or often ruled out by the EW precision data [21]. Generally, in these schemes noticeable contributions to strongly constrained Flavor Changing Neutral Current (FCNC) processes appear together with too large contributions to Peskin-Takeuchi (especially, to 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 remain viable scenarios of the dynamical EWSB [22–25].

In this paper, we continue investigation of promising phenomenological implications of the vector-like TC model proposed recently in Ref. [25]. 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. [26–29]). 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 (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 the Higgs mechanism is not the primary source of the EWSB, but effectively induced by an unknown TC dynamics at high scales.

The vector-like TC model [25] is based upon phenomenologically successful gauged linear σ-model (GLσM) initially proposed in Ref. [30] and further elaborated in Refs. [31, 32]. 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 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 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 never at the fundamental level.

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

$$SU(2)_L \otimes SU(2)_R \to SU(2)_{V \equiv L+R} . \tag{1.1}$$

In Ref. [25] it was argued that in the low energy limit the vector-like gauge group $SU(2)_V$ should be *identified* with the weak isospin group $SU(2)_W$ of the SM, i.e.[1]

$$SU(2)_{V \equiv L+R} \simeq SU(2)_W . \tag{1.2}$$

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[1] In addition, in this paper we explore a possible option when gauged $SU(2)_V$ is associated with another

Such a "gauging" of the vector subgroup $SU(2)_{\text{V} \equiv \text{L+R}}$ and its identification with the SM gauge isospin group do not mean that one introduces extra elementary gauge bosons to the existing fundamental theory, e.g. to the SM or its possible high-scale gauge extensions. In our context, procedure (1.2) means the 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 [25]. Note, such an identification is automatic at the fundamental T-quark level – the Dirac T-quarks reside in the weak isospin group $SU(2)_{\text{W}}$ from the beginning. Indeed, in the high energy limit of the theory, the global classification symmetry $SU(2)_{\text{L}} \otimes SU(2)_{\text{R}}$ is restored, while Dirac T-quark fields, along with chiral SM fermion fields, reside in fundamental representations of the SM gauge $SU(2)_{\text{W}}$ and no extra fundamental gauge bosons are needed.

The most critical part of the proposal is that the resulting unbroken local chiral-symmetric subgroup $SU(2)_{\text{V} \equiv \text{L+R}}$ describes the gauge interactions of constituent T-quarks and T-hadrons with local gauge couplings *in the low energy limit of the effective field theory* especially interesting for phenomenology. As of the primary goal of this work, we would like to test against astrophysics DM data if these interactions are ordinary weak or not under an assumption for an odd $SU(3)_{\text{TC}}$ group in confinement.

One should remember that identification of the local vector subgroup of the chiral group with the SM weak isospin group (1.2) is a purely phenomenological procedure which leads to correct results in the low energy limit of the theory. In reality, of course, the global classification T-flavor group $SU_{\text{L}}(2) \otimes SU_{\text{R}}(2)$ has nothing to do with the EW gauge group of the SM. At the first stage, the T-flavor group is used for classification of *composite* T-hadrons and, in particular, predicts the existence of T-pions, T-sigma and T-baryons states. At the second stage, one notices that T-quarks entering the composite T-hadrons besides T-strong interactions participate also in the fundamental EW interactions. One should therefore calculate the *EW form factors* of composite T-hadrons. The corresponding EW interactions must then be also introduced at *the fundamental T-quark level* consistently with those at *the composite level T-hadron level*. At the third stage, in the phenomenologically interesting low-energy limit of the theory the EW form factors approach the renormalized EW constants (since the T-hadron substructure does not emerge at relatively small momentum transfers). The latter should be calculated after *reclassification* of T-hadrons under the EW group representations. This three-fold generic scheme will be used below for description of EW interactions of T-hadrons.

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 $SU(2)_{\text{W}}$ gauge bosons with vector-like couplings, in opposition to ordinary SM fermions, which interact under $SU(2)_{\text{W}}$ by means of their left-handed components only.

Remarkably enough, in this model the oblique (Peskin-Takeuchi) parameters and FCNC corrections turn out to be naturally very small and fully consistent with the current EW constraints as well as with the most recent Higgs couplings measurements at the LHC in the limit of small Higgs-T-sigma mixing. Most importantly, this happens naturally in the

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fundamental gauge group in the T-quark sector different from the SM weak isospin group such that corresponding gauge bosons Z', W' are very heavy and their mixing with SM gauge bosons is strongly suppressed. For example, this group can be the one coming from the LR-symmetric generalization of the SM at high scales.

standard quantum-field theory framework implemented in rigorous quark-meson approaches of hadron physics without attracting any extra holographic or other special arguments from unknown high-scale physics. 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 further considerations. 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 $\tilde{\pi}_a$, $a = 1, 2, 3$, and constituent T-quarks $\tilde{Q}$ which acquire masses via the T-quark condensate as an external source and the T-sigma vev. We are focused on phenomenological studies of such a low energy effective field theory at typical momentum transfers squared $Q^2 \ll \Lambda_{\mathrm{TC}}^2$ in the considered linear σ-model framework without attempting to construct a high energy unifying new strongly coupled dynamics with the SM at the moment.

Note also that the vector-like TC model offers a simple method of phenomenologically consistent construction of the vector-like ultraviolet completion of a strongly-coupled theory which can be further exploited in the composite Higgs models as well as in attempts for Grand-like TC unification with the SM at high scales (see e.g. Refs. [33, 34] and references therein).

The VLTC scenario represents the very first step focussing 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 criterium. 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. Most importantly, 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 criterium, other three important points which are considered to be primary achievements of the VLTC model [25] can be summarized as follows:

- The effective Higgs mechanism of dynamical EW symmetry breaking in the conformal limit of the theory forbidding Higgs μ-terms is naturally emerged in this approach. The Higgs vev is automatically expressed in terms of the T-quark condensate such that the EW symmetry is broken simultaneously with the chiral symmetry breaking. No T-pions are eaten and remain physical, they escape current detection limits due to extremely suppressed loop-induced couplings to two or even three gauge bosons (depending on T-quark hypercharge and TC gauge group) at leading order only, and can remain very light.

- The phenomenologically and theoretically consistent minimal vector-like UV completion based upon the linear σ-model with the global chiral $SU_{\mathrm{L}}(2) \otimes SU_{\mathrm{R}}(2)$ symmetry group is proposed. In the minimal VLTC, the model works perfectly with only two vector-like T-flavors easily passing the EW constraints and (almost) standard Higgs couplings without any extra assumptions. The model is capable of unique predictions of possibly small Higgs couplings deviations from the standard ones.

- 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. [25], does not attempt to resolve the naturalness/hierarchy problem of the SM and does not offer a mechanism for generation of 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).

Making three above points work coherently together in a single phenomenologically motivated model is the first important step towards a consistent high-energy description of vector-like Technicolor dynamics made in Ref. [25]. And here we aim at analysis of immediate implications of this scenario in Dark Matter searches. New heavy vector-like T-baryon states, T-proton P and T-neutron N states, at a TeV mass scale are naturally predicted and introduced in a similar way as in low-energy hadron physics allowing for a possible interpretation of the DM in the considering framework. A thorough analysis of distinct features of the T-neutron DM with generic weak-type $SU(2)_V$ interactions due to a vector-like character of T-neutron gauge interactions and additional T-strong channels (via T-pion/T-sigma) along with existing direct DM detection constraints is the primary goal of our current study.

Even though the EW precision constraints are satisfied for any $SU(n)_{TC}$ group with vector-like weak interactions [25], 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 a 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. In particular, we will discuss for which $SU(n)_{TC}$ groups in confinement it is possible to make the identification of the gauge groups (1.2) 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 II provides a brief overview of vector-like $SU(3)_{TC}$ TC model with Dirac T-baryons with generic weak-type $SU(2)$ interactions (before identification (1.2)). Section III contains a discussion of the T-baryon mass spectrum, in particular, important mass splitting between T-proton and T-neutron. In Section IV, we consider typical T-baryon annihilation processes in the cosmological plasma in two different cases – in the high- and low-symmetry phases. Section V is devoted to a discussion of cosmological evolution of T-neutrons in two cases of symmetric and asymmetric DM. In Section VI, major implications of direct detection limits to the considering 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 [35], which poses an important constraint on the rank of the confined group in the T-quark sector under condition (1.2). Finally, Section VII contains basic concluding remarks.

II. VECTOR-LIKE T-BARYON INTERACTIONS

As one of the basic predictions of the vector-like and other bosonic TC models, the EWSB occurs by means of the effective Higgs mechanism induced by a condensation of confined fermions. The basic hypothesis which should be thoroughly tested against both astrophysical (primarily, DM) and collider (new exotic lightest T-hadron states) data can be formulated

as follows: *the energy scales of the EWSB and T-confinement have a common quantum-topological nature and are determined by a non-perturbative dynamics of the T-quark–T-gluon condensate.* This work aims at testing this hypothesis against DM astrophysics data.

The key difference of the vector-like TC approach earlier developed in Ref. [25] from other known dynamical EWSB mechanisms is primarily in the vector-like character of gauge interactions of T-quarks as a natural consequence of the local chiral symmetry breaking (1.1) and a possible identification of the local chiral vector-like subgroup with the weak isospin group of the SM (1.2). The latter requirement strongly reduces extra loop contributions to the EW observables and the new dynamics may slip away from the EW precision tests and ongoing Higgs couplings studies at the LHC even in the case of relative proximity of the T-confinement $\Lambda_{\rm TC}$ and EW $M_{\rm EW} \sim 100$ GeV scales. So, frequent references to "Technicolor" as to a "dead concept" in the literature do not apply to the vector-like TC model, at least, at the current level of experimental precision. Let us remind a few basic features of this scheme relevant to the forthcoming discussion of cosmological consequences, in particular, properties of the DM.

We start with the simplest way to introduce vector-like gauge interactions of elementary T-quarks and composite T-baryons based upon the gauged linear σ-model (GLσM) [30–32]. Recently, this scheme was applied to description of LHC phenomenology of the lowest mass composites – the physical pseudoscalar T-pions $\tilde{\pi}^{\pm,0}$ and scalar $\tilde{\sigma}$-meson, as well as possible modifications of the scalar Higgs boson h couplings, which are relevant for LHC searches for new strongly-coupled dynamics and precision Higgs physics [25].

Consider the local chiral vector-like subgroup $SU(2)_{\rm V\equiv L+R}$ appearing due to the spontaneous chiral symmetry breaking (1.1) and acting on new confined elementary T-quark and simultaneously composite T-baryon sectors. For the moment, we do not assume the condition (1.2), i.e. we do not explicitly introduce ordinary weak interactions into the T-quark/T-baryon sectors. Following to hadron physics analogy, let us extend the fermion sector by incorporating one Dirac T-nucleon vector-like doublet $\tilde{N}$ over $SU(2)_{\rm V}$ (an analog of the nucleon doublet in the SM) in addition to the elementary T-quark vector-like doublet $\tilde{Q}$ (an analog of the first generation of quarks in the SM) such that the initial matter fields content of the vector-like TC model becomes

$$\tilde{Q} = \begin{pmatrix} U \\ D \end{pmatrix}, \qquad \tilde{N} = \begin{pmatrix} P \\ N \end{pmatrix}, \tag{2.1}$$

which are in the fundamental representation of the $SU(2)_{\rm V} \otimes U(1)_{\rm Y}$ group. As usual, in addition we have the initial scalar T-sigma S field which is the singlet representation, and the triplet of initial T-pion fields P_a, $a = 1, 2, 3$ which is the adjoint (vector) representation of $SU(2)_{\rm V}$ (with zeroth $U(1)_{\rm Y}$ hypercharge). Thus, in terms of the fields introduced above the GLσM part of the Lagrangian responsible for Yukawa-type interactions of the T-quarks (2.1) reads

$$\mathcal{L}_Y^{\rm TC} = -g_{\rm TC}^Q \bar{\tilde{Q}}(S + i\gamma_5\tau_a P_a)\tilde{Q} - g_{\rm TC}^N \bar{\tilde{N}}(S + i\gamma_5\tau_a P_a)\tilde{N}, \qquad g_{\rm TC}^Q \neq g_{\rm TC}^N, \quad g_{\rm TC}^{Q,N} > 1, \tag{2.2}$$

where τ_a, $a = 1, 2, 3$ are the Pauli matrices, and T-strong Yukawa couplings $g_{\rm TC}^Q$ and $g_{\rm TC}^N$ are introduced in a complete analogy to low-energy hadron physics, they absorb yet unknown non-perturbative strongly-coupled dynamics and can be chosen to be different. Typically, the perturbativity condition requires them to be bounded, $g_{\rm TC}^{Q,N} < \sqrt{4\pi}$, in order to trust predictions of the linear model. After the EWSB phase, the Yukawa interactions (2.2) will

play an important role determining the strength of T-neutron N self-interactions leading to specific properties of the associated DM which will be studied below.

In the SM, the ordinary gauge boson-hadron interactions are usually introduced by means of gauge bosons hadronisation effects. In the case of a relatively large T-confinement scale $\Lambda_{\text{TC}} \sim 0.1 - 1$ TeV relevant for our study, the effect of T-hadronisation of light W, Z bosons into heavy composite states is strongly suppressed by large constituent masses of T-quarks $M_Q \sim \Lambda_{\text{TC}}$. Following to arguments of Ref. [25], the vector-like interactions of $\tilde{Q}$, $\tilde{N}$ and P_a fields with initial $U(1)_\text{Y}$ and $SU(2)_\text{V}$ gauge fields B_μ, V_μ^a, respectively, can be safely introduced in the local approximation via covariant derivatives over the local $SU(2)_\text{V} \otimes U(1)_\text{Y}$ group in the same way as ordinary SM gauge interactions, i.e.

$$\mathcal{L}_{\text{kin}}^{\text{TC}} = \frac{1}{2}\partial_\mu S\,\partial^\mu S + \frac{1}{2}D_\mu P_a\,D^\mu P_a + i\bar{\tilde{Q}}\hat{D}\tilde{Q} + i\bar{\tilde{N}}\hat{D}\tilde{N}\,. \tag{2.3}$$

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

$$\begin{aligned} \hat{D}\tilde{Q} &= \gamma^\mu \left(\partial_\mu - \frac{iY_Q}{2}\,g_1 B_\mu - \frac{i}{2}\,g_2^V V_\mu^a \tau_a \right) \tilde{Q}\,, \\ \hat{D}\tilde{N} &= \gamma^\mu \left(\partial_\mu - \frac{iY_N}{2}\,g_1 B_\mu - \frac{i}{2}\,g_2^V V_\mu^a \tau_a \right) \tilde{N}\,, \\ D_\mu P_a &= \partial_\mu P_a + g_2^V \epsilon_{abc} V_\mu^b P_c\,, \end{aligned} \tag{2.4}$$

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

$$SU(2)_\text{V} \to SU(2)_\text{W}\,, \qquad V_\mu^a \to W_\mu^a\,, \qquad g_2^V \to g_2 \tag{2.5}$$

replacement rule in Eq. (2.4). A consistency test of the latter scenario against the DM relic abundance and direct DM detection data for rank-2 confined group will enable us to draw important conclusions about properties of TC sectors.

The additional gauge and Yukawa parts (2.2) and (2.3) should be added to the SM Lagrangian written in terms of SM gauge B_μ, W_μ^a and chiral fields as follows

$$\mathcal{L}_{\text{SM}}^{\text{gauge}} = -\frac{1}{2}g_1 B_\mu \bar{l}_L \gamma^\mu l_L - g_1 B_\mu \bar{e}_R \gamma^\mu e_R + \frac{1}{6}g_1 B_\mu \bar{q}_L \gamma^\mu q_L + \frac{2}{3}g_1 B_\mu \bar{u}_R \gamma^\mu u_R - \frac{1}{3}g_1 B_\mu \bar{d}_R \gamma^\mu d_R \tag{2.6}$$

$$+\frac{1}{2}g_2 W_\mu^a \bar{l}_L \gamma^\mu \tau_a l_L + \frac{1}{2}g_2 W_\mu^a \bar{q}_L \gamma^\mu \tau_a q_L\,.$$

Here summation over flavor and family indices is implied. The theory in its simplest formulation discussed here, of course, does not predict particular values for 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 confinement, etc. should be ultimately constrained in extended chiral-gauge or grand-unified theories along with coming 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 nucleon doublet in the SM, i.e. $Y_N = 1$. Thus, the T-baryon states in Eq. (2.1) 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 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 $SU(3)_{\rm TC}$ confined group of QCD-type are generic for other odd $SU(2n + 1)_{\rm TC}, n = 2, 3, \ldots$ groups. In this work we stick to a direct analogy with QCD for simplicity and test it against available DM constraints.

One of the interesting but alternative opportunities would be to consider $Y_Q = 0$ case such that an integer electric charge of T-baryons would only be possible for even TC groups $SU(2n)_{\rm TC}$ with the simplest $SU(2)_{\rm TC}$. Here, T-baryons are two-T-quark systems. In the non-perturbative T-hadron vacuum the UD state with zeroth electric charge is energetically favorable since 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 to be absolutely stable and thus an appealing DM candidate. This case has certain advantages and will be considered elsewhere.

After $SU(2)_{\rm V} \otimes U(1)_{\rm Y}$ symmetry breaking an extra set of heavy gauge $Z', W'^{\pm}$ bosons interacting with T-quark and T-baryons emerges[2]. If one does not imply a straightforward identification (2.5), 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 to not spoil the EW precision tests. An alternative option would be to adopt (2.5) where the EW precision tests are satisfied without any serious tension [25]. Then, in analogy to constituent colored T-quarks [25], the vector-like SM gauge interactions of T-baryons with $Z, W^{\pm}$ bosons are controlled by the following part of the Lagrangian

$$L_{\bar{\tilde{N}}\tilde{N}Z/W} = \delta_W \frac{g_2}{\sqrt{2}} \bar{P}\gamma^\mu N \cdot W^+_\mu + \delta_W \frac{g_2}{\sqrt{2}} \bar{N}\gamma^\mu P \cdot W^-_\mu$$
$$+ \delta_Z \frac{g_2}{c_W} Z_\mu \sum_{f=P,N} \bar{f}\gamma^\mu \big(t_3^f - q_f\, s_W^2\big)f\,. \tag{2.7}$$

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 ($t_3^P = 1/2$, $t_3^N = -1/2$), $q_f = Y_N/2 + t_3^f$ 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:

$$\begin{array}{lll}
\text{I.} & \delta_{W,Z} = 1\,, & SU(2)_{\rm V} \simeq SU(2)_{\rm W}\,, \\
\text{II.} & \delta_{W,Z} \ll 1\,, & SU(2)_{\rm V} \neq SU(2)_{\rm W}\,, \quad m_{Z',W'} \gg 100\,{\rm GeV}\,.
\end{array} \tag{2.8}$$

In the first case, one deals with pure EW vector interactions of T-quarks/T-baryons corresponding to transition (2.5), while in the second case $\delta_{W,Z}$ are related to a very small

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[2] Such extra $Z', W'^{\pm}$ bosons can, in principle, be composite and identified with composite $\rho^{0,\pm}$ mesons or elementary vector bosons from "right isospin" $SU(2)_{\rm R}$ group as a part of chiral-symmetric $SU(2)_{\rm L} \otimes SU(2)_{\rm R}$ extension of the SM. This point, however, is not critical for the current study of the QCD-type T-neutron DM properties and we do not discuss it here.

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 are not shown here. We will test the both options above against 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

$$
\begin{aligned}
L_{\bar{N}\tilde{N}h} + L_{\bar{N}\tilde{N}\tilde{\sigma}} + L_{\bar{N}\tilde{N}\tilde{\pi}} &= -g_{\text{TC}}^N \left(c_\theta \tilde{\sigma} + s_\theta h\right) \cdot (\bar{P}P + \bar{N}N) \\
&- i\sqrt{2}g_{\text{TC}}^N \tilde{\pi}^+ \bar{P}\gamma_5 N - i\sqrt{2}g_{\text{TC}}^N \tilde{\pi}^- \bar{N}\gamma_5 P - ig_{\text{TC}}^N \tilde{\pi}^0(\bar{P}\gamma_5 P - \bar{N}\gamma_5 N) \,.
\end{aligned} \tag{2.9}
$$

The gauge and Yukawa parts of the Lagrangian (2.7) and (2.9) completely determine the T-baryon interactions at relatively low kinetic energies $E_{\text{kin}} \ll M_{B_{\text{T}}}$ 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 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 [25]. The latter is true for both models I and II (2.8).

The interactions of T-pions with Z, $W^{\pm}$ bosons which will be used in further calculations of T-baryon annihilation cross sections are defined as follows

$$
\begin{aligned}
L_{\tilde{\pi}\tilde{\pi}Z/W} &= ig_2^W W^{\mu+} \cdot (\tilde{\pi}^0\tilde{\pi}^-_{,\mu} - \tilde{\pi}^-\tilde{\pi}^0_{,\mu}) + ig_2^W W^{\mu-} \cdot (\tilde{\pi}^+\tilde{\pi}^0_{,\mu} - \tilde{\pi}^0\tilde{\pi}^+_{,\mu}) \\
&+ ig_2^Z c_W Z_\mu \cdot (\tilde{\pi}^-\tilde{\pi}^+_{,\mu} - \tilde{\pi}^+\tilde{\pi}^-_{,\mu}) \,,
\end{aligned} \tag{2.10}
$$

where $g_2^{W,Z} = g_2 \delta_{W,Z}$, $\tilde{\pi}_{,\mu} \equiv \partial_\mu \tilde{\pi}$. The Yukawa interactions $\bar{f}fh + \bar{f}f\tilde{\sigma}$ of the ordinary fermions get modified compared to the SM as follows

$$
L_{\bar{f}fh} + L_{\bar{f}f\tilde{\sigma}} = -g_2(c_\theta h - s_\theta \tilde{\sigma}) \cdot \frac{m_f}{2M_W} \bar{f}f \,. \tag{2.11}
$$

The Lagrangians of the $h\tilde{\pi}\tilde{\pi}$ and $hWW + hZZ$ interactions will also be needed below, so we write them down here as well:

$$
L_{h\tilde{\pi}\tilde{\pi}} = -(\lambda_{\text{TC}} u\, s_\theta - \lambda v c_\theta)\, h(\tilde{\pi}^0\tilde{\pi}^0 + 2\tilde{\pi}^+\tilde{\pi}^-) = -\frac{M_h^2 - m_{\tilde{\pi}}^2}{2M_Q} g_{\text{TC}} s_\theta\, h(\tilde{\pi}^0\tilde{\pi}^0 + 2\,\tilde{\pi}^+\tilde{\pi}^-) \,,
$$

$$
L_{hWW} + L_{hZZ} = g_2 M_W c_\theta\, hW_\mu^+ W^{\mu-} + \frac{1}{2}(g_1^2 + g_2^2)^{1/2} M_Z c_\theta\, hZ_\mu Z^\mu \,. \tag{2.12}
$$

Finally, the interactions $\tilde{\sigma}\tilde{\pi}\tilde{\pi}$ and $\tilde{\sigma}WW + \tilde{\sigma}ZZ$ are determined by

$$
L_{\tilde{\sigma}\tilde{\pi}\tilde{\pi}} = -(\lambda_{\text{TC}} u c_\theta + \lambda v s_\theta)\, \tilde{\sigma}(\tilde{\pi}^0\tilde{\pi}^0 + 2\,\tilde{\pi}^+\tilde{\pi}^-) = -\frac{M_{\tilde{\sigma}}^2 - m_{\tilde{\pi}}^2}{2M_Q} g_{\text{TC}} c_\theta\, \tilde{\sigma}(\tilde{\pi}^0\tilde{\pi}^0 + 2\,\tilde{\pi}^+\tilde{\pi}^-) \,,
$$

$$
L_{\tilde{\sigma}WW} + L_{\tilde{\sigma}ZZ} = -g_2 M_W s_\theta\, \tilde{\sigma}W_\mu^+ W^{\mu-} - \frac{1}{2}(g_1^2 + g_2^2)^{1/2} M_Z s_\theta\, \tilde{\sigma}Z_\mu Z^\mu \,. \tag{2.13}
$$

In the considering vector-like TC model, the T-baryon mass scale $\sqrt{s} \lesssim M_{B_{\text{T}}}$ should be considered as an upper cut-off of the considering model 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 $\sqrt{s} \gtrsim M_{B_{\text{T}}}$ 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.

III. T-BARYON MASS SPECTRUM

Typical (but optional) assumption about a dynamical similarity between color and T-color in the case of confined $SU(3)_{\text{TC}}$ 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 $\zeta = \Lambda_{\text{TC}}/\Lambda_{\text{QCD}} \gtrsim 1000$ following from a relative proximity of EW scale and $\Lambda_{\text{TC}} \sim 0.1 - 1$ TeV, i.e.

$$m_{\tilde{\pi}} \gtrsim 140\,\text{GeV}\,, \quad M_{\tilde{\sigma}} \gtrsim 500\,\text{GeV}\,, \quad M_Q \gtrsim 300\,\text{GeV}\,, \quad M_{B_{\text{T}}} \equiv M_P \simeq M_N \gtrsim 1\,\text{TeV}\,, \tag{3.1}$$

for the T-pion $m_{\tilde{\pi}}$, T-sigma $M_{\tilde{\sigma}}$ and constituent T-quark M_Q and T-baryon $M_{B_{\text{T}}}$ 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

$$M_{B_{\text{T}}} \equiv M_P \simeq M_N \simeq 3M_Q = 3g_{\text{TC}}^Q u\,, \tag{3.2}$$

where $u \sim \Lambda_{\text{TC}} = 0.1 - 1$ 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 (1.1) (for more details on the chiral and EW symmetries breaking in the considering model, see Ref. [25]). In the chiral limit of the theory, T-sigma vev u has the same quantum-topological nature as the SM Higgs vev $v \simeq 246$ 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 $Q^2 \ll l_{\text{TC}}^{-2} \gtrsim 2.3\,\text{TeV}^2$, where l_{TC} is the characteristic length scale of the non-perturbative T-gluon fluctuations estimated by rescaling of that from QCD (3.1), 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 about T-baryon number conservation 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 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 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 the 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 subgroup $SU(2)_{\text{V}}$ 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 T-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 non-zeroth anomalous magnetic moments of 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 $M_{EW} \sim 100$ 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 provides arbitrary but *exactly equal* current T-quark masses in the initial Lagrangian, $m_U = m_D$, such that fundamental T-quark and hence T-baryon spectra are degenerate at tree level. Note in the initial SM Lagrangian current u, d-quark masses are equal to zero due to chiral asymmetry of 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 considering case of 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, $\Delta M_{B_T}^{EW} \lesssim \Delta M_{B_T} \equiv M_P - M_N \ll M_{B_T}$. Corresponding EW ($Z, \gamma$-mediated) one-loop diagrams are shown in Fig. 1.
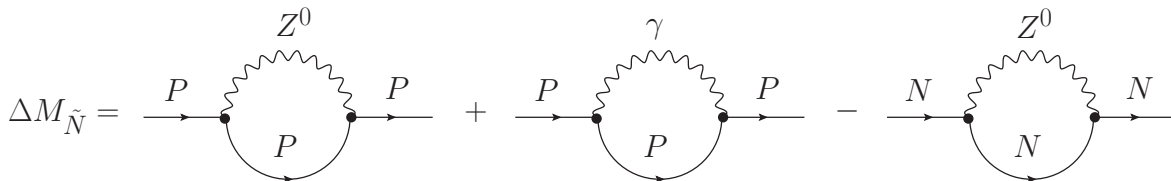

FIG. 1: One-loop EW radiative corrections causing positive mass splitting between T-proton and T-neutron in the chiral limit of the underlined theory, $\Delta M_{B_T}^{EW} > 0$. Other corrections from $W^{\pm}$, T-pion, T-sigma and Higgs boson loops enter symmetrically to P and N self-energies and thus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

In the model I (2.8) 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 a form of the following finite integral

$$\Delta M_{B_T}^{EW} = -\frac{ie^2 M_Z^2}{8\pi^4} \int \frac{(\hat{q} - M_{B_T})dq}{q^2(q^2 - M_Z^2)[(q+p)^2 - M_{B_T}^2]}, \tag{3.3}$$

given by γ and Z corrections shown in Fig. 1 only. Note that logarithmic divergences explicitly cancel out in the difference between P and N mass operators, providing us with the finite result for $\Delta M_{B_T}^{EW}$. Other corrections from $W^{\pm}$, scalar and pseudoscalar loops enter symmetrically into P and N self-energies and thus are canceled out too. In the realistic case

of heavy T-baryons $M_{B_{\mathrm{T}}} \gg M_Z$ we arrive at the following simple relation

$$\Delta M_{B_{\mathrm{T}}}^{\mathrm{EW}} \simeq \frac{\alpha(M_{B_{\mathrm{T}}})M_Z}{2} > 0 \,, \tag{3.4}$$

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

$$\Delta M_{B_{\mathrm{T}}}^{\mathrm{EW}} \simeq 360 \,\mathrm{MeV} \,. \tag{3.5}$$

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, μ, $\nu_{e,\mu}$, 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 such that the latter turns out to be stable and viable as a heavy DM candidate. With the mass splitting value (3.5), we find the following approximate vector-like T-proton lifetime

$$\tau_P \simeq \frac{15\pi^3}{G_F^2}(\Delta M_{B_{\mathrm{T}}}^{\mathrm{EW}})^{-5} \simeq 0.4 \times 10^{-9} \,\mathrm{s} \,. \tag{3.6}$$

In the model II, the Z contributions die out in the limit $\delta_{\mathrm{EW}} \ll 1$, so it can only be induced by extra heavy Z' exchange. The corresponding contribution to the P-N mass splitting $\Delta M_{B_{\mathrm{T}}}^{\mathrm{V}}$ is obtained from Eq. (3.4) by a replacement $m_Z \to m'_Z$,

$$\Delta M_{B_{\mathrm{T}}}^{\mathrm{V}} \simeq \frac{\alpha(M_{B_{\mathrm{T}}})M_{Z'}}{2} > 0 \,, \qquad M_{Z'} < M_{B_{\mathrm{T}}} \,, \tag{3.7}$$

such that the $\Delta M_{B_{\mathrm{T}}}^{\mathrm{V}} \gg \Delta M_{B_{\mathrm{T}}}^{\mathrm{EW}}$ and the T-proton lifetime would even be shorter. Of course, the estimate (3.7) should be taken with care for $m_{Z'} \gtrsim M_{B_{\mathrm{T}}}$ 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, in the most natural and simplest model I the properties of the vector-like T-baryon spectrum are very similar to properties of vector-like Higgsino LSP (e.g. splitting between chargino and neutralino) spectrum due to practically the same structure of EW interactions [36]. The key difference between the lightest Higgsino and T-neutron DM candidates is in capability of T-neutrons to self-interactions (e.g. enhanced self-annihilation and elastic scattering rates) driven essentially by T-strong Yukawa terms (2.9) which make them specifically interesting for DM phenomenology and astrophysics.

IV. ANNIHILATION OF T-BARYONS IN COSMOLOGICAL PLASMA

In the considering vector-like TC model the interaction properties of T-baryons are fixed by gauge and Yukawa interactions determined by Eqs. (2.7) and (2.9), while the radiative splitting in the mass spectrum between P and N states is given by Eq. (3.4) or (3.7). Besides an unknown rank of the confined group, physics of vector-like T-baryon DM, its interaction properties and formation depend quantitatively on four *physical* parameters only: the strong T-hadron coupling g_{TC}^N, T-quark mass scale M_Q, the T-pion $m_{\tilde{\pi}}$ and T-sigma $M_{\tilde{\sigma}}$ masses. The physically interesting parameter space domain corresponds to relatively small

$\tilde{\sigma}h$-mixing angle $s_\theta \lesssim 0.2$ where the oblique corrections (essentially, T-parameter) are small and deviations from the SM Higgs couplings are strongly suppressed and can be well within the current LHC constraints (see Ref. [25] and references therein). 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 on the current (in)direct DM detection measurements and constraints, see e.g. Refs. [37–40]).

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

In practice, the phenomenologically acceptable domain of the vector-like TC parameter space corresponding to $s_\theta \ll 1$ means that the hierarchy of EW and chiral symmetry breaking scales is far from degeneracy, i.e. $u/v \gg 1$. This means that the realistic T-baryon mass scale can be decoupled, although not very strongly, from the EW one, i.e. $M_{B_{\mathrm{T}}} \gg M_{\mathrm{EW}}$ for not too small T-quark/T-baryon Yukawa couplings $g_{\mathrm{TC}}^{N,Q} \gtrsim 1$. The latter means that the T-baryon spectrum, and possibly (pseudo)scalar spectrum, would be way above the Higgs boson mass scale, and the naive QCD scaling (3.1) can be satisfied.

On the other hand, it is also possible that T-pions evade LEP II and current low mass LHC constraints due to very small (T-quark loop induced) production cross sections and narrow widths, and thus they can, in principle, be as light as W boson, $m_{\tilde{\pi}} \gtrsim m_W$. While T-sigma is extremely wide $\Gamma_{\tilde{\sigma}} \sim M_{\tilde{\sigma}}$ and hard to be detected with current collider techniques, it would be possible for it to have a rather low mass down to ~ 150 GeV or even lower without upsetting current EW precision and LHC constraints. Note, current LHC constraints on the T-pion mass from the ordinary TC (e.g. Extended TC) scenarios do not apply to the considering vector-like TC model where T-pions do not couple to ordinary fermions.

The last two possibilities of decoupled and non-decoupled TC sectors in Nature make it reasonable to consider the irreversible annihilation of T-baryons in two different phases of cosmological plasma separately – before and after EW phase transition epoch $T_{\mathrm{EW}} \sim 200$ GeV. Consequently, we will end up with two different scenarios of the DM relic abundance formation which have to be (qualitatively) discussed in detail.

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

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

In the HS case with $SU(2)_{\mathrm{V}} = SU(2)_{\mathrm{W}}$ (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. $\Delta M_{B_{\mathrm{T}}} = 0$ (or more precisely, $T \gg \Delta M_{B_{\mathrm{T}}}$), so the total T-baryon number density is $n_{B_{\mathrm{T}}} \simeq 2(n_N + n_{\bar{N}})$, at least, before the
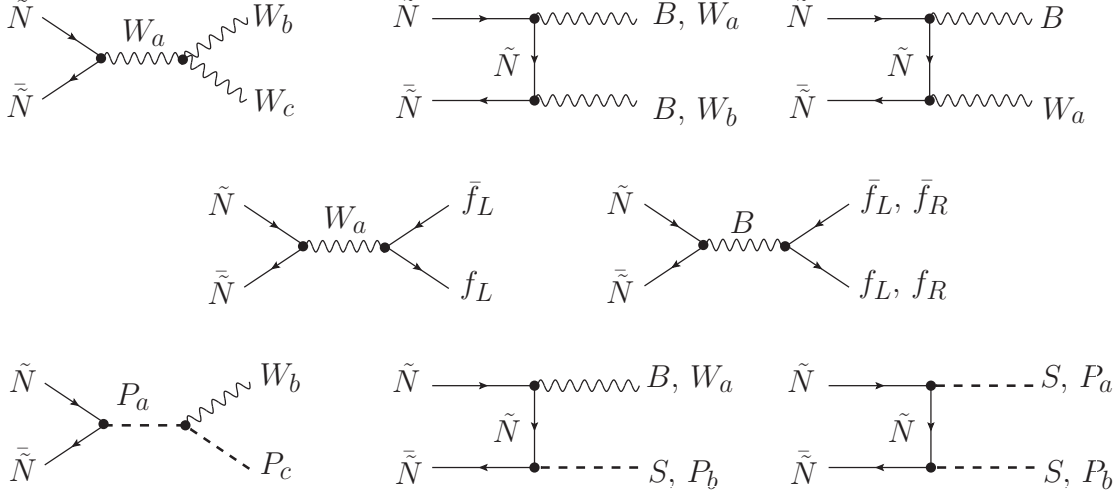
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FIG. 2: Typical diagrams contributing to the T-baryon DM annihilation in the high-symmetry phase of the cosmological plasma corresponding to $T_{\text{EW}} < T \lesssim u$. The model I (2.8) is implied.

T-baryon freeze out epoch. Practically, P and N states are dynamically equivalent in this phase and participate in all reactions as components of the isospin $SU(2)_{\text{V}}$ doublet $\tilde{N}$ (2.1) with Yukawa and gauge interactions determined by Eqs. (2.2) and (2.3), respectively. Consequently, masses of all SM fermions and gauge bosons vanish in this phase (more precisely, m_f, $M_{W,Z} \ll T$), while 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 $M_{B_{\text{T}}} \sim u$ since $u \gg v$, i.e. $M_{\tilde{\sigma}}$, $m_{\tilde{\pi}} \ll M_{B_{\text{T}}}$. Thus, all the masses except for 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 $M_{B_{\text{T}}}$, which has to be constrained together with the strong Yukawa coupling g_{TC}^N from astrophysics data.

Let us evaluate the vector-like T-baryon annihilation cross section $(\sigma v_B)_{\text{ann}}$ in the HS phase of the cosmological plasma in the model I (2.8). All relevant contributions are schematically depicted in Fig. 2. In comparison with the Higgsino LSP scenario in $SU(5)$ Split SUSY Model [36], the T-baryon annihilation in the HS phase is given by essentially the same EW amplitudes due to the same vector-like structure of T-baryon and Higgsino EW interactions, i.e.

$$\tilde{N}\bar{\tilde{N}} \to BB\,, \qquad \tilde{N}\bar{\tilde{N}} \to BW_a\,, \qquad \tilde{N}\bar{\tilde{N}} \to W_aW_b\,,$$
$$\tilde{N}\bar{\tilde{N}} \to B^* \to l_L\bar{l}_L\,, \ q_L\bar{q}_L\,, \ e_R\bar{e}_R\,, \ u_R\bar{u}_R\,, \ d_R\bar{d}_R\,, \qquad \tilde{N}\bar{\tilde{N}} \to W_a^* \to l_L\bar{l}_L\,, \ q_L\bar{q}_L\,, \ (4.1)$$

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 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

$$(\sigma v_B)_{\text{ann}}^{\text{EW}} = \frac{21g_1^4 + 6g_1^2 g_2^2 + 39g_2^4}{512\pi M_{B_{\text{T}}}^2}\,. \tag{4.2}$$

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

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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 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

$$\tilde{N}\bar{\tilde{N}} \to P_a^* \to P_b W_c \,, \qquad \tilde{N}\bar{\tilde{N}} \to P_a B \,, \, SB \,, \, SW_a \,, \qquad (4.3)$$

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

$$(\sigma v_B)_{\mathrm{ann}}^{\mathrm{EW+TC}} \simeq \frac{(g_{\mathrm{TC}}^N)^2 (2g_1^2 + 3g_2^2)}{32\pi M_{B_{\mathrm{T}}}^2} \,. \qquad (4.4)$$

In order to turn to the model II (2.8) in the limit $m_{Z'} \ll M_{B_{\mathrm{T}}}$ corresponding to unbroken $SU(2)_{\mathrm{V}} \neq SU(2)_{\mathrm{W}}$, one has to perform a replacement $g_2 \to g_2^{\mathrm{V}}$ in Eqs. (4.2) and (4.4). This would provide a rough estimate for the annihilation cross sections into B_μ, V_μ^a bosons. For a more precise analysis of the gauge $SU(2)_{\mathrm{V}}$ part of the cross sections one should consider details of the broken phase of $SU(2)_{\mathrm{V}}$ and evaluate them for massive $Z', W'^\pm$ bosons for kinematically allowed channels, i.e. for $m_{W',Z'} \lesssim M_{B_{\mathrm{T}}}/2$ (for more details, see calculations in the low-symmetry phase below). The latter, however, do not affect our conclusions here since corresponding cross sections are relatively small compared to those in the T-strong channels.

For pure T-strong channels

$$\tilde{N}\bar{\tilde{N}} \to P_a P_b \,, \, S P_a \,, \, S S \,, \qquad (4.5)$$

we have the total cross section

$$(\sigma v_B)_{\mathrm{ann}}^{\mathrm{TC}} \simeq \frac{9(g_{\mathrm{TC}}^N)^4}{32\pi M_{B_{\mathrm{T}}}^2} \,. \qquad (4.6)$$

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

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

$$(\sigma v_B)_{\mathrm{ann}} \simeq (\sigma v_B)_{\mathrm{ann}}^{\mathrm{TC}} \,, \qquad (4.7)$$

which makes the considering T-baryon DM model specific compared to other standard SUSY-based DM models where $(\sigma v_B)_{\mathrm{ann}} \sim \alpha_{\mathrm{W}}^2/M_\chi^2$, $\alpha_{\mathrm{W}} \simeq 1/30$ given by weak interactions only. Thus, more intense T-baryons annihilation with extremely weak interactions with ordinary matter (see below) makes them promising DM candidates alternative to standard WIMPs. Note that the value (4.6) behaves as forth power of g_{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).

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

Again, consider first the phenomenologically appealing model I (2.8) in detail. At lower temperatures, $T < T_{\mathrm{EW}}$ 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-zeroth 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. neutralino LSP [41]. This type of DM annihilation dynamics would happen entirely after the EW phase transition epoch for rather low mass T-baryons, i.e. $M_{B_\mathrm{T}} < 3\,T_\mathrm{EW} \sim 600$ 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_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 $(\sigma v_B)_\mathrm{ann}$ in this case are listed in Fig. 3.

Here, only T-neutrons participate in the annihilation processes. Indeed, T-protons have very small mean life-time (3.6) so they rapidly decay soon after the EW phase transition epoch and can not 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. 3 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.
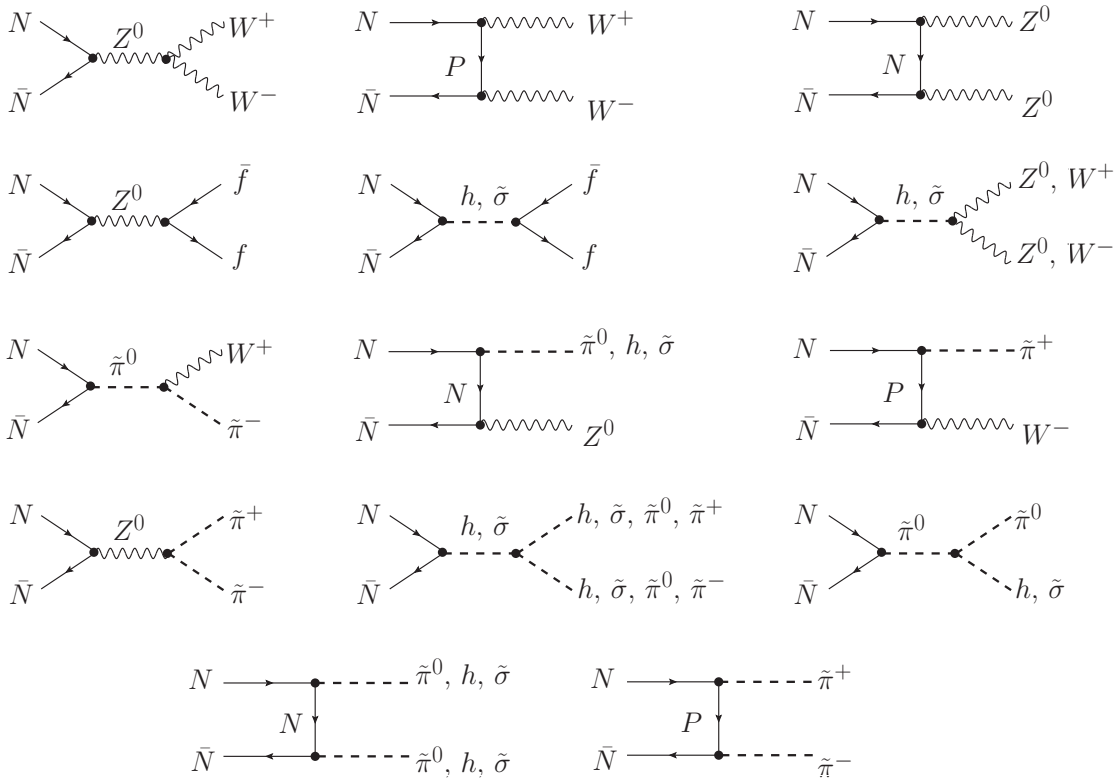

FIG. 3: Typical diagrams contributing to the T-neutron DM annihilation in the low-symmetry phase of the cosmological plasma $T < T_\mathrm{EW}$ including much later stages of structure formation and the present epoch. The model I (2.8) is implied.

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 T-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 v_B^2 order including interference terms and can be neglected.

The EW channels with SM fermion and gauge boson final states shown in the first two lines in Fig. 3 are

$$N\bar{N} \;\to\; Z^{0*}, h^*, \tilde{\sigma}^* \;\to\; l\bar{l}, \; q\bar{q}, \qquad N\bar{N} \;\to\; W^+W^-, \; Z^0Z^0 \,, \qquad (4.8)$$

respectively. The annihilation cross section into fermions is

$$(\sigma v_B)_{\mathrm{ann}}^{f\bar{f}} \;\simeq\; \frac{1}{192\pi M_{B_{\mathrm{T}}}(4M_{B_{\mathrm{T}}}^2 - m_Z^2)^2}\Big\{2M_{B_{\mathrm{T}}}^3(103g_1^4 + 6g_1^2g_2^2 + 63g_2^4) + $$
$$(17g_1^4 - 6g_1^2g_2^2 + 9g_2^4)(2M_{B_{\mathrm{T}}}^2 + m_t^2)\sqrt{M_{B_{\mathrm{T}}}^2 - m_t^2}\Big\} \,, \qquad (4.9)$$

into W^+W^- pair

$$(\sigma v_B)_{\mathrm{ann}}^{W^+W^-} \;\simeq\; \frac{g_2^4}{64\pi M_{B_{\mathrm{T}}} m_W^4} \frac{(M_{B_{\mathrm{T}}}^2 - m_W^2)^{3/2}}{(2M_{B_{\mathrm{T}}}^2 - m_W^2)^2(4M_{B_{\mathrm{T}}}^2 - m_Z^2)^2} \times$$
$$\Big\{4M_{B_{\mathrm{T}}}^4(12m_W^4 - 4m_W^2m_Z^2 + m_Z^4) + $$
$$4M_{B_{\mathrm{T}}}^2 m_W^2(20m_W^4 - 24m_W^2m_Z^2 + 5m_Z^4) + $$
$$m_W^4(12m_W^4 - 12m_W^2m_Z^2 + 5m_Z^4)\Big\} \,, \qquad (4.10)$$

and into Z^0Z^0 pair

$$(\sigma v_B)_{\mathrm{ann}}^{Z^0Z^0} \;\simeq\; \frac{g_2^4(M_{B_{\mathrm{T}}}^2 - m_Z^2)^{3/2}}{64\pi c_W^4 M_{B_{\mathrm{T}}}(2M_{B_{\mathrm{T}}}^2 - m_Z^2)^2} \qquad (4.11)$$

Above we have neglected all the fermion masses assuming for simplicity $m_{l,q} \ll M_{B_{\mathrm{T}}}$ except for top-quark mass $m_t \simeq 173$ GeV. Also, 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 have an order of $\sim v_B^2$. So the $\tilde{\sigma}$ and h-mediated diagrams were neglected in the integrated cross sections (4.9), (4.10) and (4.11).

The channels with mixed (gauge and (pseudo)scalar) final states (third line in Fig. 3) are

$$N\bar{N} \;\to\; W^{\pm}\tilde{\pi}^{\mp} \,, \qquad N\bar{N} \;\to\; Z^0\tilde{\pi}^0 \,, \qquad N\bar{N} \;\to\; Z^0\tilde{\sigma} \,, \qquad N\bar{N} \;\to\; Z^0h \,, \qquad (4.12)$$

leading to the following contributions

$$(\sigma v_B)_{\mathrm{ann}}^{W^{\pm}\tilde{\pi}^{\mp}} \;\simeq\; \frac{g_2^2(g_{\mathrm{TC}}^N)^2\Big[16M_{B_{\mathrm{T}}}^4 - 8M_{B_{\mathrm{T}}}^2(m_{\tilde{\pi}}^2 + m_W^2) + (m_{\tilde{\pi}}^2 - m_W^2)^2\Big]^{3/2}}{32\pi M_{B_{\mathrm{T}}}^4(4M_{B_{\mathrm{T}}}^2 - m_{\tilde{\pi}}^2)^2(4M_{B_{\mathrm{T}}}^2 - m_{\tilde{\pi}}^2 - m_W^2)^2} \times \qquad (4.13)$$
$$\Big\{(4M_{B_{\mathrm{T}}}^2 - m_{\tilde{\pi}}^2)^2 + 2M_{B_{\mathrm{T}}}^2 m_W^2\Big\} \,,$$

$$(\sigma v_B)_{\mathrm{ann}}^{Z^0\tilde{\pi}^0} \;\simeq\; \frac{g_2^2(g_{\mathrm{TC}}^N)^2\Big[16M_{B_{\mathrm{T}}}^4 - 8M_{B_{\mathrm{T}}}^2(m_{\tilde{\pi}}^2 + m_Z^2) + (m_{\tilde{\pi}}^2 - m_Z^2)^2\Big]^{3/2}}{128\pi c_W^2 M_{B_{\mathrm{T}}}^4(4M_{B_{\mathrm{T}}}^2 - m_{\tilde{\pi}}^2 - m_Z^2)^2} \,, \qquad (4.14)$$

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$$(\sigma v_B)^{Z^0\tilde{\sigma}}_{\text{ann}} \simeq \frac{g_2^2(g_{\text{TC}}^N)^2\, c_\theta^2 \sqrt{16M_{B_{\text{T}}}^4 - 8M_{B_{\text{T}}}^2(M_{\tilde{\sigma}}^2 + m_Z^2) + (M_{\tilde{\sigma}}^2 - m_Z^2)^2}}{128\pi c_W^2 M_{B_{\text{T}}}^4 (4M_{B_{\text{T}}}^2 - M_{\tilde{\sigma}}^2 - m_Z^2)^2} \times$$
$$\left\{ 16M_{B_{\text{T}}}^4 - 8M_{B_{\text{T}}}^2(M_{\tilde{\sigma}}^2 - 2m_Z^2) + (M_{\tilde{\sigma}}^2 - m_Z^2)^2 \right\}, \qquad (4.15)$$

$$(\sigma v_B)^{Z^0 h}_{\text{ann}} \simeq \frac{g_2^2(g_{\text{TC}}^N)^2\, s_\theta^2 \sqrt{16M_{B_{\text{T}}}^4 - 8M_{B_{\text{T}}}^2(m_h^2 + m_Z^2) + (m_h^2 - m_Z^2)^2}}{128\pi c_W^2 M_{B_{\text{T}}}^4 (4M_{B_{\text{T}}}^2 - m_h^2 - m_Z^2)^2} \times$$
$$\left\{ 16M_{B_{\text{T}}}^4 - 8M_{B_{\text{T}}}^2(m_h^2 - 2m_Z^2) + (m_h^2 - m_Z^2)^2 \right\}, \qquad (4.16)$$

where $m_h \simeq 126$ GeV is the Higgs boson mass [42].

Finally, last two lines in Fig. 3 represent diagrams with T-strong final states (T-pions and T-sigma) as well as the Higgs boson:

$$N\bar{N} \rightarrow \tilde{\pi}^+\tilde{\pi}^-, \qquad N\bar{N} \rightarrow \tilde{\pi}^0\tilde{\pi}^0, \qquad N\bar{N} \rightarrow \tilde{\pi}^0\tilde{\sigma}, \qquad N\bar{N} \rightarrow \tilde{\pi}^0 h,$$
$$N\bar{N} \rightarrow h\tilde{\sigma}, \qquad N\bar{N} \rightarrow \tilde{\sigma}\tilde{\sigma}, \qquad N\bar{N} \rightarrow hh. \qquad (4.17)$$

The relevant contributions to the total cross section are

$$(\sigma v_B)^{\tilde{\pi}^+\tilde{\pi}^-}_{\text{ann}} \simeq \frac{(M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2)^{3/2}}{16\pi M_{B_{\text{T}}}} \left(\frac{g_2^4}{(4M_{B_{\text{T}}}^2 - m_Z^2)^2} + \frac{4(g_{\text{TC}}^N)^4}{(2M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2)^2} \right), \qquad (4.18)$$

$$(\sigma v_B)^{\tilde{\pi}^0 h}_{\text{ann}} \simeq \frac{(g_{\text{TC}}^N)^2\, s_\theta^2 \sqrt{16M_{B_{\text{T}}}^4 - 8M_{B_{\text{T}}}^2(m_h^2 + m_{\tilde{\pi}}^2) + (m_h^2 - m_{\tilde{\pi}}^2)^2}}{256\pi M_{B_{\text{T}}}^4 M_Q^2 (4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2)^2(4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2 - m_h^2)^2} \times$$
$$\left(g_{\text{TC}}^Q M_{B_{\text{T}}}(m_{\tilde{\pi}}^2 - m_h^2)(4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2 - m_h^2) + \right.$$
$$\left. 2g_{\text{TC}}^N M_Q(4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2)(4M_{B_{\text{T}}}^2 + m_{\tilde{\pi}}^2 - m_h^2) \right)^2, \qquad (4.19)$$

$$(\sigma v_B)^{\tilde{\pi}^0\tilde{\sigma}}_{\text{ann}} \simeq \frac{(g_{\text{TC}}^N)^2\, c_\theta^2 \sqrt{16M_{B_{\text{T}}}^4 - 8M_{B_{\text{T}}}^2(M_{\tilde{\sigma}}^2 + m_{\tilde{\pi}}^2) + (M_{\tilde{\sigma}}^2 - m_{\tilde{\pi}}^2)^2}}{256\pi M_{B_{\text{T}}}^4 M_Q^2 (4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2)^2(4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2 - M_{\tilde{\sigma}}^2)^2} \times$$
$$\left(g_{\text{TC}}^Q M_{B_{\text{T}}}(m_{\tilde{\pi}}^2 - M_{\tilde{\sigma}}^2)(4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2 - M_{\tilde{\sigma}}^2) + \right.$$
$$\left. 2g_{\text{TC}}^N M_Q(4M_{B_{\text{T}}}^2 - m_{\tilde{\pi}}^2)(4M_{B_{\text{T}}}^2 + m_{\tilde{\pi}}^2 - M_{\tilde{\sigma}}^2) \right)^2, \qquad (4.20)$$

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 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 $SU(2)_{\text{V}} \neq SU(2)_{\text{W}}$ symmetry, one has to make the following replacements $m_Z \rightarrow m_{Z'}$, $m_W \rightarrow m_{W'}$, $g_2 \rightarrow g_2^{\text{V}}$ in the above formulae (4.9)–(4.20).

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

It is worth to mention here that in practice depending on T-baryon mass $M_{B_{\mathrm{T}}}$ and intensity of annihilation processes an intermediate scenario when the annihilation epoch starts in the HS phase and terminates in the LS phase is possible. In this case, the residual T-baryon density may be strongly correlated with details of EW phase transition epoch, as well as on relative annihilation intensities and time periods during which the plasma "lives" in each of the phases. Such an analysis is certainly much more complicated and invokes a larger amount of unknowns into the problem. However, if the considering vector-like TC model is phenomenologically justified, a more detailed analysis of cosmological evolution of the T-baryon component along these lines would be necessary. Instead, we are focused here on two simplistic toy-scenarios capturing basic dynamics of vector-like T-baryons in cosmological environment when their annihilation happens either in the HS or LS phase entirely.

V.   COSMOLOGICAL EVOLUTION OF VECTOR-LIKE T-BARYONS

The T-proton lifetime (3.6) has an order of Universe age at the EW phase transition epoch t_{EW}. So after this epoch the T-proton remnants quickly decay into T-neutrons and light SM fermions without leaving significant traces on cosmological evolution of cold DM and ordinary matter. Indeed, due to a rather small T-nucleon mass splitting (3.5) 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 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 the first approximation. Let us consider two possible scenarios of T-baryon DM relic abundance formation – symmetric and asymmetric DM cases.

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

Consider thermal evolution of 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 to each other at all stages of Universe evolution until the present epoch, i.e.

$$n_{B_{\mathrm{T}}} \simeq n_{\bar{B}_{\mathrm{T}}}\,, \qquad n_{B_{\mathrm{T}}} - n_{\bar{B}_{\mathrm{T}}} \ll n_{B_{\mathrm{T}}}\,, \tag{5.1}$$

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

For details of the thermal evolution and freeze out of symmetric heavy relic, see e.g. Ref. [43] 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 $M_{B_{\mathrm{T}}}$ in the radiation-dominated epoch, i.e. $\bar{\varepsilon}_f \simeq 3T_0 \simeq M_{B_{\mathrm{T}}}$ 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 $n_{B_{\mathrm{T}}} = n_{B_{\mathrm{T}}}(t)$ at later times $t \gg t_0$ is then described by a standard solution of the Boltzmann evolution equations [43].

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

$$t_0 \simeq \frac{1}{4T_0^2} \left(\frac{3}{2\pi G w} \right)^{1/2} \ll t_{\mathrm{EW}}, \qquad T_0 \simeq \frac{M_{B_{\mathrm{T}}}}{3}, \qquad (5.2)$$

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 $T > T_{\mathrm{EW}} \sim 200$ 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_* happens around the QCD phase transition $T_{\mathrm{QCD}} \sim 100$ 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 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 the chemical equilibrium. Then the freeze out of heavy T-baryons happens at $t = t_1$ when temperature of the Universe $T_1 = T_1(t_1) \gtrsim T_{\mathrm{EW}}$ given by the standard formula

$$T_1 \simeq \frac{M_{B_{\mathrm{T}}}}{\log \left(\frac{g_{B_{\mathrm{T}}} M_{B_{\mathrm{T}}} M_{\mathrm{PL}}^*(\sigma v_B)_{\mathrm{ann}}}{(2\pi)^{3/2}} \right)} \ll M_{B_{\mathrm{T}}}, \qquad M_{B_{\mathrm{T}}} \ll M_{\mathrm{PL}}^*, \qquad (5.3)$$

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

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

$$\rho_{B_{\mathrm{T}}}(t_{\mathrm{U}}) = M_{B_{\mathrm{T}}} n_{B_{\mathrm{T}}}(t_{\mathrm{U}}) \simeq \rho_{\mathrm{DM}}(t_{\mathrm{U}}), \qquad (5.4)$$

provides the canonical constraint on thermally averaged kinetic annihilation cross section known from Ref. [43]

$$(\sigma v_B)_{\mathrm{ann}}^{\mathrm{DM}} \simeq 2.0 \times 10^{-9} \text{ GeV}^{-2}. \qquad (5.5)$$

In terms of the latter, the relic DM abundance is

$$\Omega_{\mathrm{DM}} \simeq 0.2 \left[(\sigma v_B)_{\mathrm{ann}}^{\mathrm{DM}} / (\sigma v_B)_{\mathrm{ann}}^{\mathrm{th}} \right]. \qquad (5.6)$$

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

By comparing the above astrophysical constraint with the theoretical cross section given by Eq. (4.6) one extracts the lower bound for the T-baryon mass scale

$$M_{B_{\mathrm{T}}} \gtrsim 5 \text{ TeV}, \qquad g_{\mathrm{TC}}^N \gtrsim 1.0 \tag{5.7}$$

Clearly, this bound is consistent with naive QCD scaling hypothesis (3.1). 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 T-strong Yukawa coupling, $(\sigma v_B)_{\mathrm{ann}} \sim (g_{\mathrm{TC}}^N)^4$. 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 the model I (2.8) demonstrating consistency of the considering scenario of T-baryon relic formation. For the model II, a high-scale $SU(2)_{\mathrm{V}}$ symmetry is likely to be broken at temperatures $T < M_{B_{\mathrm{T}}}/3$ so in this case the annihilation occurs in the low-symmetry (w.r.t. $SU(2)_{\mathrm{V}}$) phase. Note, somewhat lower $M_{B_{\mathrm{T}}}$ estimates are also possible for smaller g_{TC}^N values in both models I and II, but then mixed EW+TC and pure EW channels become important, such that the mean value $M_{B_{\mathrm{T}}}$ never goes below 3 TeV corresponding to "switched off" Yukawa TC interactions, i.e. $g_{\mathrm{TC}}^N \to 0$.

Note, the estimate (5.7) should be taken with care since they have 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 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;

and 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 $M_{B_{\mathrm{T}}}$ and M_{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 relic T-baryon abundance evolution lifting out one or more of the specific assumptions would be necessary.

Finally, consider formation of the relic symmetric DM density in the LS phase of the cosmological plasma for the model I set-up. It appears to be very hard to realize such a scenario with $M_{B_{\mathrm{T}}} \lesssim 600$ GeV due to a large $(\sigma v_B)_{\mathrm{ann}}^{\tilde{\pi}^+ \tilde{\pi}^-}$ contribution for $g_{\mathrm{TC}}^N \gtrsim 1$ and $M_{B_{\mathrm{T}}} \gg m_{\tilde{\pi}}$ (see also Fig. 4). Formally, it may still be possible to tune $m_{\tilde{\pi}}$ and g_{TC}^N in a special way by minimizing the total LS cross section given by

$$(\sigma v_B)_{\mathrm{ann}}^{\mathrm{LS}} \simeq (\sigma v_B)_{\mathrm{ann}}^{\tilde{\pi}^+ \tilde{\pi}^-} + (\sigma v_B)_{\mathrm{ann}}^{\tilde{\pi}^0 \tilde{\sigma}} + (\sigma v_B)_{\mathrm{ann}}^{\tilde{\pi}^0 h}, \tag{5.8}$$

but the corresponding TC model would be unnatural. In the model II, there is a larger freedom in relative (rather fine) tuning between large $m_{W'}$, $m_{Z'}$ and M_{B_T} parameters capable of reducing the annihilation cross section. This, however, cannot bring M_{B_T} down significantly without spoiling EW precision constraints and keeps it roughly at the same level (5.7). Other possibility would be to have the annihilation epoch in the mixed HS/LS phase by relaxing the constraint on the T-baryon mass scale $M_{B_T} \lesssim 600$ GeV. The LS and HS annihilation cross sections become comparable for $M_{B_T} \gtrsim 3$ TeV, such that in this case most of the T-baryon annihilation epoch happens in the HS phase again justifying the above estimate (5.7).
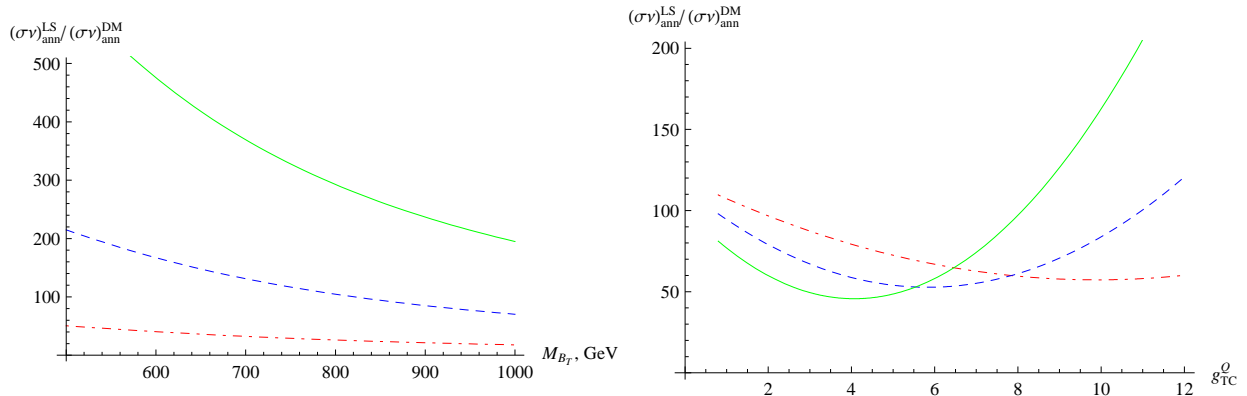

FIG. 4: Dependence of the ratio of kinetic T-baryon cross section in the LS phase for the model I $(\sigma v_B)_{\mathrm{ann}}^{\mathrm{LS}}$ (5.8) to the kinetic cross section of the DM $(\sigma v_B)_{\mathrm{ann}}^{\mathrm{DM}} \simeq 2.0 \times 10^{-9}$ GeV^{-2} required by observations on T-neutron mass M_{B_T} (left) and on T-quark Yukawa coupling g_{TC}^Q (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, three different $g_{\mathrm{TC}}^N = 1.0, 1.4, 1.8$ values correspond to dash-dotted, dashed and solid lines, respectively, and fixed $m_{\tilde{\pi}} = 250$ GeV and $g_{\mathrm{TC}}^Q = 2.0$. In the right panel, $m_{\tilde{\pi}} = 150, 200, 250$ GeV correspond to dash-dotted, dashed and solid lines, respectively, and fixed $M_{B_T} = 400$ GeV and $g_{\mathrm{TC}}^N = 1.0$.

Thus, we come to a 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.

B. Scenario II: Asymmetric vector-like T-baryon Dark Matter

As was argued above the LS annihilation cross section (5.8) (model I) becomes too high for low mass T-baryons. Indeed, in Fig. 4 we show typical parameter dependencies of the ratio $(\sigma v_B)_{\mathrm{ann}}^{\mathrm{LS}}/(\sigma v_B)_{\mathrm{ann}}^{\mathrm{DM}}$. It turns out that the LS cross section is by 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 $M_{B_T} \lesssim 600$ 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. So, the bulk of observable DM density is essentially given in terms of

the T-baryon asymmetry $\Delta n_{B_{\mathrm{T}}}$, i.e.

$$\rho_{\mathrm{DM}}(t_{\mathrm{U}}) \simeq \Delta n_{B_{\mathrm{T}}} M_{B_{\mathrm{T}}}, \qquad \Delta n_{B_{\mathrm{T}}} \equiv n_{B_{\mathrm{T}}} - n_{\bar{B}_{\mathrm{T}}}, \tag{5.9}$$

even though it has been negligible in the beginning of the T-baryon annihilation epoch, $\Delta n_{B_{\mathrm{T}}} \ll n_{B_{\mathrm{T}}}(t_0)$. This is the so-called asymmetric DM (ADM) model which has been previously studied in other TC/compositeness scenarios (see e.g. Refs. [44–46]), and it can be realized in the considering vector-like TC model as well.

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

$$r \equiv \frac{n(\bar{N})}{n(N)}, \qquad 0 < r < 1, \tag{5.10}$$

one could estimate its time evolution via the detailed analysis of the system of coupled Boltzmann equations performed in Refs. [47, 48]. 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.

$$r_\infty \sim \exp\left\{ -2\frac{(\sigma v_B)_{\mathrm{ann}}^{\mathrm{LS}}}{(\sigma v_B)_{\mathrm{ann}}^{\mathrm{DM}}} \right\} \ll 1 \tag{5.11}$$

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. 4 one concludes that there practically no light anti-T-neutrons remain in 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 [49].

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 $M_{B_{\mathrm{T}}} \gtrsim 200$ GeV such 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 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 details of the chemical equilibrium the T-baryon mass scale is expected to be rather high $M_{B_{\mathrm{T}}} \sim$ TeV as was advocated in Ref. [50] 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 $g_{\mathrm{TC}}^N \sim 1$ and high T-neutron masses $M_{B_{\mathrm{T}}} \gtrsim 1$ TeV the ratio of the cross sections goes down

$$1 < \frac{(\sigma v_B)_{\mathrm{ann}}^{\mathrm{LS}}}{(\sigma v_B)_{\mathrm{ann}}^{\mathrm{DM}}} \lesssim 10, \tag{5.12}$$

and 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 the model I and LS annihilation in the model II. In particular, having a small but non-zeroth 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 self-interacting DM model which allows

to avoid problematic cusp-like DM density profile in the central regions of galactic haloes leading to a core-like DM distribution favored by astrophysical observations [51]. In the considering scenario, the intensive annihilation together with elastic NN scattering play an important role of self-interactions of DM particles allowing for necessary adjustments of the DM density profile. Indeed, in the early structure formation epoch an overdensity of DM in cusp-like regions has been eliminated by intensive annihilation processes such that they do not exist today.

Another interesting astrophysical implication of the partially ADM scenario in general is a 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 (5.8) 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-neutrons 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 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 to explain why the DM in the Galactic halo is much more tepid than ordinary CDM models predict. 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 then CDM WIMP-based simulations predict [52–54]. In the considering 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.

VI. DIRECT T-NEUTRON DETECTION CONSTRAINTS

At last, 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 spin-independent (SI) component of the elastic WIMP-nucleon cross section from CDMS II [55] and, especially, XENON100 [35] 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. (2.7). This leads to a sizable SI scattering cross section off nuclei which needs to be compared to the data.

Let us consider the 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. [36, 56, 57]), so we do not go into details of explicit calculations here. Following to Ref. [57] the SI spin-averaged T-neutron-nucleus cross section reads

$$\sigma_{\mathrm{SI}} = \frac{\mu^2}{16\pi M_{B_{\mathrm{T}}}^2 m_A^2} \left(\frac{1}{4} \sum_{\mathrm{spins}} |\mathcal{M}_{\mathrm{SI}}|^2 \right) , \qquad (6.1)$$

where μ is the reduced mass and m_A is the mass of the target nucleus. In the considering 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

$$\mathcal{O}_Z^q = \delta_Z \, \lambda_{qV} \, \frac{ig_2^2}{2c_W^2 m_Z^2} \, \bar{N}\gamma^\mu N \cdot \bar{q}\gamma_\mu q \,, \qquad (6.2)$$

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

$$\lambda_{qV} = \frac{g_2}{c_W} \left[t_{qL}^3 - 2Q_q s_W^2 \right] . \qquad (6.3)$$

This leads to a squared matrix element (cf. Ref. [57])

$$\frac{1}{4} \sum_{\mathrm{spins}} |\mathcal{M}_{\mathrm{SI}}|^2 \simeq 16 \, \frac{M_{B_{\mathrm{T}}}^2 m_A^2}{m_Z^4} \, \frac{g_2^4}{4c_W^4} \, \delta_Z^2 \left[\sum_{q=u,d} \lambda_{qV} [ZB_{qV}^p + (A-Z)B_{qV}^n] \right]^2 , \qquad (6.4)$$

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 individual elastic N-p and N-n cross sections one obtains

$$\sigma_{\mathrm{SI}}^{N-p,n} = \delta_Z^2 \, \frac{g_2^4}{16\pi c_W^4} \, \frac{\mu^2}{m_Z^4} \, (c_V^{p,n})^2 \,, \qquad c_V^p = 1 - 4c_W^2 \,, \quad c_V^n = -1 \,, \qquad (6.5)$$

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

$$\sigma_{\mathrm{SI}}^{N-A} = \delta_Z^2 \, \frac{g_2^4}{16\pi c_W^4} \, \frac{\mu^2}{m_Z^4} \left[Z \, c_V^p + (A-Z) \, c_V^n \right]^2 . \qquad (6.6)$$

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

$$\sigma_{\mathrm{SI}}^{N-p} = 1.5 \times 10^{-40} \, \mathrm{cm}^2 \times \delta_Z^2 \left(\frac{\mu}{m_p} \right)^2 , \qquad \sigma_{\mathrm{SI}}^{N-n} = 2.5 \times 10^{-38} \, \mathrm{cm}^2 \times \delta_Z^2 \left(\frac{\mu}{m_n} \right)^2 . (6.7)$$

These cross sections are rather large and strongly constrained by direct detection experiments. At the moment, XENON100 [35] provides the most stringent limit on $\sigma_{\mathrm{SI}}^{\mathrm{nucleon}}$ (per nucleon in the case of a xenon target) for high mass Dirac T-neutrons which is roughly

$$- \log_{10} \left(\frac{\sigma_{\mathrm{SI}}^{\mathrm{nucleon}}}{\mathrm{cm}^2} \right) \simeq 44.6 - 43.4 \,, \qquad (6.8)$$

corresponding to the T-neutron mass range of about

$$M_{B_{\mathrm{T}}} \simeq 0.1 - 2\,\mathrm{TeV}\,, \qquad (6.9)$$

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

$$\delta_Z \lesssim 2 \times 10^{-3}\,, \qquad M_{B_{\mathrm{T}}} \lesssim 2\,\mathrm{TeV}\,, \qquad (6.10)$$

and a little bit weaker constraint for $M_{B_{\mathrm{T}}} > 2\,\mathrm{TeV}$. 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 $\delta_Z = 1$ introduced above in Eq. (2.8), are firmly ruled out. The only possibility to accommodate the Dirac T-baryons is in the model II with $\delta_Z \ll 1$ which, however, requires an extra assumption about the existence of an extra vector $SU(2)_{\mathrm{V}} \neq SU(2)_{\mathrm{W}}$ gauge symmetry in the T-quark/T-baryon sector. The constraint (6.10) is then the upper limit on the mixing parameter (e.g. sine of a mixing angle) between Z and new Z' bosons from an unspecified high-scale $SU(2)_{\mathrm{V}}$ which should also be additionally constrained by EW precision tests and extra gauge bosons (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.

VII. 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 [25] and have found important limitations on the structure of T-strong dynamics from the direct DM detection data (in particular, by the XENON100 experiment [35]). This has been done in the simplest QCD-like setting of the T-confinement $SU(3)_{\mathrm{TC}}$ group and one generation of T-quarks. We have shown that under a natural assumption about the T-baryon number conservation, the local chiral symmetry breaking gives rise to the vector $SU(2)_{\mathrm{V}}$ gauge symmetry, which acts on T-quark and T-baryons sectors additional to the SM sectors. Whether the $SU(2)_{\mathrm{V}}$ group is identified with the weak isospin $SU(2)_{\mathrm{W}}$ symmetry of the SM or not provides us with the two possible models I and II, respectively, for the structure of (vector) weak interactions in the T-quark/T-baryon sectors which have been studied above in detail.

As a crucial specific prediction of the vector-like TC model, the T-quark and hence the T-baryon mass spectrum is degenerated at tree level. The EW radiative corrections with vector T-baryon-Z, γ couplings effectively split the T-baryon sector making the T-proton $P = (UUD)$ (slightly) heavier than the T-neutron $N = (UDD)$ irrespectively of the nature of $SU(2)_{\mathrm{V}}$ gauge 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 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.

As was discussed thoroughly in Ref. [25], the EW precision constraints at the fundamental level can only be satisfied for the vector-like T-quarks under the weak isospin $SU(2)_{\mathrm{W}}$. Alternatively, one could introduce the vector-like weak interactions via a small mixing $\delta_Z \ll 1$ which can be limited to not 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 can be consistent with the DM astrophysics constraints. This provides an extra very important limit on the structure of TC sectors and new strongly-coupled dynamics.

Indeed, the Dirac T-baryons were originated under the simplest assignment of an odd T-confined group in the T-quark sector having rank three, i.e. $SU(3)_{\text{TC}}$. Thus, we conclude that in this case or, more generally, in the case of any odd group $SU(2n+1)_{\text{TC}}$, $n = 1, 2, \ldots$, it is not possible to introduce the standard EW interactions over the SM $SU(2)_{\text{W}}$ gauge group in a phenomenologically consistent way as suggested by the XENON100 constraint (6.10). The only way to satisfy the existing phenomenological (EW, collider and astrophysics) constraints is to consider an even T-confinement $SU(2n)_{\text{TC}}$, $n = 1, 2, \ldots$ group, for example, the simplest $SU(2)_{\text{TC}}$ where the lightest stable neutral T-baryon state $B_0 = (UD)$ is scalar and does not interact with the Z-boson thus evading the direct detection limits. Corresponding analysis is ongoing.

[1] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1303.5076 [astro-ph.CO].

[2] V. Trimble, Ann. Rev. Astron. Astrophys. **25**, 425 (1987).

[3] H. P. Nilles, Phys. Rept. **110**, 1 (1984);
    H. E. Haber and G. L. Kane, Phys. Rept. **117**, 75 (1985).

[4] G. Jungman, M. Kamionkowski and K. Griest, Phys. Rept. **267**, 195 (1996) [hep-ph/9506380].

[5] G. Bertone, D. Hooper and J. Silk, Phys. Rept. **405**, 279 (2005) [hep-ph/0404175].

[6] L. Roszkowski, Pramana **62**, 389 (2004) [hep-ph/0404052].

[7] C. Boehm, P. S. B. Dev, A. Mazumdar and E. Pukartas, JHEP **1306**, 113 (2013) [arXiv:1303.5386 [hep-ph]];
    S. Scopel, N. Fornengo and A. Bottino, arXiv:1304.5353 [hep-ph];
    A. Choudhury and A. Datta, arXiv:1305.0928 [hep-ph];
    A. Fowlie, K. Kowalska, L. Roszkowski, E. M. Sessolo and Y. -L. S. Tsai, arXiv:1306.1567 [hep-ph].

[8] S. Nussinov, Phys. Lett. B **165**, 55 (1985);
    S. M. Barr, R. S. Chivukula, and E. Farhi, Phys. Lett. B **241**, 387 (1990).

[9] S. B. Gudnason, C. Kouvaris and F. Sannino, Phys. Rev. D **74**, 095008 (2006) [hep-ph/0608055];
    S. B. Gudnason, C. Kouvaris and F. Sannino, Phys. Rev. D **73**, 115003 (2006) [hep-ph/0603014].

[10] M. Y. Khlopov and C. Kouvaris, Phys. Rev. D **78**, 065040 (2008) [arXiv:0806.1191 [astro-ph]];
    M. Y. Khlopov, A. G. Mayorov and E. Y. Soldatov, Int. J. Mod. Phys. D **19**, 1385 (2010)

[arXiv:1003.1144 [astro-ph.CO]].

[11] E. Del Nobile and F. Sannino, Int. J. Mod. Phys. A **27**, 1250065 (2012) [arXiv:1102.3116 [hep-ph]].

[12] F. Sannino, Acta Phys. Polon. B **40**, 3533 (2009) [arXiv:0911.0931 [hep-ph]].

[13] T. A. Ryttov and F. Sannino, Phys. Rev. D **78**, 115010 (2008) [arXiv:0809.0713 [hep-ph]].

[14] A. Belyaev, M. T. Frandsen, S. Sarkar and F. Sannino, Phys. Rev. D **83**, 015007 (2011) [arXiv:1007.4839 [hep-ph]].

[15] K. Petraki and R. R. Volkas, arXiv:1305.4939 [hep-ph].

[16] C. Kouvaris, arXiv:1304.7476 [hep-ph].

[17] S. Weinberg, Phys. Rev. D **13**, 974 (1976);
L. Susskind, Phys. Rev. D **20**, 2619 (1979).

[18] C. T. Hill and E. H. Simmons, Phys. Rept. **381**, 235 (2003) [Erratum-ibid. **390**, 553 (2004)] [hep-ph/0203079].

[19] F. Sannino, Acta Phys. Polon. B **40**, 3533 (2009) [arXiv:0911.0931 [hep-ph]].

[20] S. Dimopoulos and L. Susskind, Nucl. Phys. B **155**, 237 (1979);
E. Eichten and K. D. Lane, Phys. Lett. B **90**, 125 (1980).

[21] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990);
M. E. Peskin, T. Takeuchi, Phys. Rev. D **46**, 381 (1992).

[22] T. W. Appelquist, D. Karabali, L. C. R. Wijewardhana, Phys. Rev. Lett. **57**, 957 (1986).

[23] F. Sannino and K. Tuominen, Phys. Rev. D **71**, 051901 (2005) [hep-ph/0405209].

[24] R. Foadi, M. T. Frandsen, T. A. Ryttov, F. Sannino, Phys. Rev. D **76**, 055005 (2007) [arXiv:0706.1696 [hep-ph]].

[25] R. Pasechnik, V. Beylin, V. Kuksa and G. Vereshkov, arXiv:1304.2081 [hep-ph].

[26] E. H. Simmons, Nucl. Phys. B **312**, 253 (1989).

[27] S. Samuel, Nucl. Phys. B **347**, 625 (1990).

[28] A. Kagan and S. Samuel, Phys. Lett. B **270**, 37 (1991).

[29] S. Samuel, Nucl. Phys. B **347**, 625 (1990);
C. D. Carone, Phys. Rev. D **86**, 055011 (2012) [arXiv:1206.4324 [hep-ph]].

[30] B.W. Lee and H.T. Nieh, Phys. Rev. **166**, 1507 (1968).

[31] S. Gasiorowicz and D. Geffen, Rev. Mod. Phys. 41, 531 (1969);
P. Ko and S. Rudaz, Phys. Rev. D 50, 6877 (1994);
M. Urban, M. Buballa, and J. Wambach, Nucl. Phys. A697, 338 (2002).

[32] B.D. Serot and J.D. Walecka, Acta Phys. Pol. B **21**, 655 (1992).

[33] D. Marzocca, M. Serone and J. Shu, JHEP **1208**, 013 (2012) [arXiv:1205.0770 [hep-ph]];
L. Vecchi, arXiv:1304.4579 [hep-ph];
M. Montull, F. Riva, E. Salvioni and R. Torre, arXiv:1308.0559 [hep-ph].

[34] F. Caracciolo, A. Parolini and M. Serone, JHEP **1302**, 066 (2013) [arXiv:1211.7290 [hep-ph]].

[35] E. Aprile *et al.* [XENON100 Collaboration], Phys. Rev. Lett. **109**, 181301 (2012) [arXiv:1207.5988 [astro-ph.CO]].

[36] V. A. Beylin, V. I. Kuksa, R. S. Pasechnik and G. M. Vereshkov, Eur. Phys. J. C **56**, 395 (2008) [hep-ph/0702148 [hep-ph]];
V. A. Beylin, V. I. Kuksa, G. M. Vereshkov and R. S. Pasechnik, Int. J. Mod. Phys. A **24**, 6051 (2009) [arXiv:0903.4201 [hep-ph]].

[37] R. W. Schnee, arXiv:1101.5205 [astro-ph.CO].

[38] M. Cirelli, Pramana **79**, 1021 (2012) [arXiv:1202.1454 [hep-ph]].

[39] M. Drees and G. Gerbier, arXiv:1204.2373 [hep-ph].

[40] L. Bergstrom, Annalen Phys. **524**, 479 (2012) [arXiv:1205.4882 [astro-ph.HE]].

[41] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, arXiv:1305.0237 [hep-ph].

[42] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716**, 1 (2012) [arXiv:1207.7214 [hep-ex]];
S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716**, 30 (2012) [arXiv:1207.7235 [hep-ex]].

[43] G. Steigman, B. Dasgupta and J. F. Beacom, Phys. Rev. D **86**, 023506 (2012) [arXiv:1204.3622 [hep-ph]].

[44] I. Masina, P. Panci and F. Sannino, JCAP **1212**, 002 (2012) [arXiv:1205.5918 [astro-ph.CO]].

[45] R. Lewis, C. Pica and F. Sannino, Phys. Rev. D **85**, 014504 (2012) [arXiv:1109.3513 [hep-ph]].

[46] E. Del Nobile, C. Kouvaris and F. Sannino, Phys. Rev. D **84**, 027301 (2011) [arXiv:1105.5431 [hep-ph]].

[47] M. L. Graesser, I. M. Shoemaker and L. Vecchi, JHEP **1110**, 110 (2011) [arXiv:1103.2771 [hep-ph]].

[48] H. Iminniyaz, M. Drees and X. Chen, JCAP **1107**, 003 (2011) [arXiv:1104.5548 [hep-ph]].

[49] M. W. Goodman and E. Witten, Phys. Rev. D **31**, 3059 (1985).

[50] M. R. Buckley and L. Randall, JHEP **1109**, 009 (2011) [arXiv:1009.0270 [hep-ph]].

[51] B. D. Wandelt, R. Dave, G. R. Farrar, P. C. McGuire, D. N. Spergel and P. J. Steinhardt, astro-ph/0006344.

[52] M. Mateo, Ann. Rev. Astron. Astrophys. **36**, 435 (1998) [astro-ph/9810070].

[53] B. Moore, S. Ghigna, F. Governato, G. Lake, T. R. Quinn, J. Stadel and P. Tozzi, Astrophys. J. **524**, L19 (1999) [astro-ph/9907411].

[54] J. Zavala, M. Vogelsberger and M. G. Walker, arXiv:1211.6426 [astro-ph.CO].

[55] Z. Ahmed *et al.* [CDMS-II Collaboration], Science **327**, 1619 (2010) [arXiv:0912.3592 [astro-ph.CO]].

[56] E. J. Chun, J. -C. Park and S. Scopel, JCAP **1002**, 015 (2010) [arXiv:0911.5273 [hep-ph]].

[57] M. R. Buckley, D. Hooper and J. Kumar, arXiv:1307.3561 [hep-ph].