Towards working technicolor: Effective theories and dark matter

Sven Bjarke Gudnason,* Chris Kouvaris,† and Francesco Sannino‡

The Niels Bohr Institute, Blegdamsvej 17, DK-2100 Copenhagen Ø, Denmark

(Received 29 March 2006; published 12 June 2006)

A fifth force, of technicolor type, responsible for breaking the electroweak theory is an intriguing extension of the standard model. Recently new theories have been shown to feature walking dynamics for a very low number of techniflavors and are not ruled out by electroweak precision measurements. We identify the light degrees of freedom and construct the associated low energy effective theories. These can be used to study signatures and relevant processes in current and future experiments. In our theory the technibaryons are pseudo Goldstone bosons and their masses arise via extended technicolor interactions. There are hypercharge assignments for the techniquarks which renders one of the technibaryons electrically neutral. We investigate the cosmological implications of this scenario and provide a component of dark matter.

DOI: 10.1103/PhysRevD.73.115003 PACS numbers: 12.60.Nz, 95.35.+d

I. INTRODUCTION

A dynamical mechanism behind the breaking of the electroweak theory is a very intriguing possibility. A new strong force is postulated to drive such a mechanism. The large hadron collider (LHC) experiment at CERN is going to directly probe the sector associated with the breaking of the electroweak theory, and hence will be able to shed light on this force. Nature has already shown to privilege such a mechanism which takes place in ordinary superconductivity as well as in the spontaneous breaking of chiral symmetry in quantum chromodynamics (QCD).

Earlier attempts using QCD-like technicolor [1] have been ruled out by precision measurements [2]. Besides, one has also to face the problem of mass generation which typically is provided by extended technicolor (ETC) interactions and thus leads to large flavor changing neutral currents. Recently it has been shown that one can construct viable theories explaining the breaking of the electroweak theory dynamically [3–7] while not being at odds with electroweak precision measurements. In the recently proposed theories, technimatter transforms according to a higher dimensional representation of the new gauge group. By direct comparison with data it turns out that the preferred representation is the two index symmetric [3]. The simplest theory of this kind is a two-technicolor theory. In this case the two index symmetric representation corresponds to the adjoint. Remarkably these theories are already near conformal for a very small number of techniflavors. Further properties of higher dimensional representations have been also explored in [8]. In [4–6] the reader can find a summary of a number of salient properties of the new technicolor theories as well as a comprehensive review of the walking properties with references to the literature. We also note that near the conformal window [9,10] one of the relevant electroweak parameters (S) is smaller than expected in perturbation theory. This observation is further supported by very recent analysis [11,12].

In this paper we examine the phenomenological implications of the technicolor theory with two techniquarks transforming according to the adjoint representation of $SU(2)$. This theory has an $SU(4)$ quantum global symmetry which breaks spontaneously to $SO(4)$. Of the nine Goldstone bosons, three are eaten by the electroweak gauge bosons while the remaining ones carry nonzero technibaryon number, which is associated with one of the diagonal generators of $SU(4)$. These technibaryons must acquire a mass from some, yet unspecified, theory at a higher scale. Since we assume a bottom up approach we postpone the problem of producing the underlying theory providing these masses, but we expect it to be similar to the ETC type theory proposed in [4]. If the technibaryon number is left intact by the ETC interactions the lightest technibaryon (LTB) is stable and the hypercharge assignment can be chosen in a way that the LTB is also electrically neutral.

In the first part of the paper, we provide the associated linear and nonlinear effective theories. The latter can be used to make specific quantitative predictions for the large hadron collider as well as linear collider (LC).

If the technibaryon number is conserved the LTB is stable and it can be made electrically neutral while its mass is expected to be of the order of the electroweak scale. It has, hence, many features required of a dark matter component. Following Refs. [13,14] we have calculated the contribution of this particle to dark matter and we found that it can account for the whole dark matter density. We should emphasize that in our calculation we took under consideration the overall electric neutrality of the matter in the Universe as well as the thermal equilibrium conditions and the sphaleron processes. There are no parameters to tune in order to get the right technibaryon density other than the mass of the neutral particle. Note that if dark

*Electronic address: gudnason@nbi.dk
†Electronic address: kouvaris@nbi.dk
‡Electronic address: sannino@nbi.dk

matter is homogeneously distributed in our galaxy our component may be the dominant contribution to dark matter [15] or can constitute part of it. However, it seems that the dark matter distribution in the Universe is not yet exactly determined [16] and depends crucially on the type and number of components.

We have made a preliminary investigation relative to the single step unification problem for the SM couplings and our technicolor coupling and found that they do not unify. This result is perhaps not too surprising since one expects new gauge theories to emerge before the typical unification scale, and besides there could be multiple step unifications.

To date, it is unclear if supersymmetry will ever play a role in nature. If supersymmetry is not discovered at the electroweak scale it can still emerge at much higher energies. The standard model (SM) and the new strong force will then become supersymmetric at this new higher scale. A feature of this scenario, proposed long ago by Dine, Fischler and Srednicki [17], is the fact that these extensions of the standard model lend a natural solution to the so called μ problem of the minimal supersymmetric standard model (MSSM). If we adopt this idea, we find that the our new strong force can be extended to $\mathcal{N} = 4$ super Yang-Mills by adding the missing scalars and suitably adjusting all of the interactions among the matter fields. We define m as the mass-scale above which the Higgs sector (i.e. now the technicolor sector) of the theory becomes $\mathcal{N} = 4$. The corresponding coupling constant freezes above m since this sector of the theory is conformal. This would make our theory an even better candidate for walking technicolor.

II. THE MODEL

The new dynamical sector underlying the Higgs mechanism we consider is an $SU(2)$ technicolor gauge group with two adjoint technifermions. The theory is asymptotically free if the number of flavors $N_f < 2.75$.

To estimate the critical coupling for chiral symmetry breaking we required that the anomalous dimension of the quark mass operator must satisfy the relation $\gamma(2 - \gamma) = 1$ [18]. This yields $\alpha_c \simeq \frac{\pi}{3N}$. The critical value of the number of flavors which gives this fixed point value is $N_f^c \simeq 2.075$ [3,4].

Since we consider adjoint Dirac fermions, the critical number of flavors is independent of the number of colors [4]. We expect that the theory will enter a conformal regime unless the coupling rises above the critical value triggering the formation of a fermion condensate. Hence a $N_f = 2$ theory is sufficiently close to the critical number of flavors N_f^c. This makes it a perfect candidate for a walking technicolor theory.

Although the critical number of flavors is independent of the number of colors the electroweak precision measurements do depend on it. Since the lowest number of colors is privileged by data [5,6] we choose the two-technicolor theory.

Then the two adjoint fermions may be written as

$$T_L^a = \begin{pmatrix} U^a \\ D^a \end{pmatrix}_L, \qquad U_R^a, \qquad D_R^a, \qquad a = 1, 2, 3, \tag{1}$$

with a the adjoint color index of $SU(2)$. The left fields are arranged in three doublets of the $SU(2)_L$ weak interactions in the standard fashion. The condensate is $\langle \bar{U}U + \bar{D}D \rangle$ which breaks correctly the electroweak symmetry.

This model as described so far suffers from the Witten topological anomaly [19]. An $SU(2)$ gauge theory must have an even number of fermion doublets to avoid this anomaly. Here there are three extra electroweak doublets added to the standard model and we need to add one more doublet. Since we do not wish to disturb the walking nature of the technicolor dynamics, the doublet must be a technicolor singlet [6]. Our additional matter content is essentially a copy of a standard model fermion family with quarks (here transforming in the adjoint of $SU(2)$) and the following lepton doublet

$$\mathcal{L}_L = \begin{pmatrix} N \\ E \end{pmatrix}_L, \qquad N_R, \qquad E_R. \tag{2}$$

In general, the gauge anomalies cancel using the following generic hypercharge assignment

$$Y(T_L) = \frac{y}{2}, \qquad Y(U_R, D_R) = \left(\frac{y+1}{2}, \frac{y-1}{2}\right), \tag{3}$$

$$Y(\mathcal{L}_L) = -3\frac{y}{2}, \qquad Y(N_R, E_R) = \left(\frac{-3y+1}{2}, \frac{-3y-1}{2}\right), \tag{4}$$

where the parameter y can take any real value. In our notation the electric charge is $Q = T_3 + Y$, where T_3 is the weak isospin generator. One recovers the SM hypercharge assignment for $y = 1/3$. In [4], the SM hypercharge has been investigated in the context of an extended technicolor theory. Another interesting choice of the hypercharge is $y = 1$, which has been investigated from the point of view of the electroweak precision measurements in [5,6]. In this case

$$\begin{gathered} Q(U) = 1, \qquad Q(D) = 0, \\ Q(N) = -1 \quad \text{and} \quad Q(E) = -2 \quad \text{with } y = 1. \end{gathered} \tag{5}$$

Notice that in this particular hypercharge assignment, the D technidown is electrically neutral. Since we have two Dirac fermions in the adjoint representation of the gauge group the global symmetry is $SU(4)$. In practice our technicolor sector has the same fermionic matter content as that of $\mathcal{N} = 4$ super Yang-Mills. To discuss the symmetry properties of the theory it is convenient to use the Weyl base for the fermions and arrange them in the following vector transforming according to the fundamental representation of $SU(4)$

$$Q = \begin{pmatrix} U_L \\ D_L \\ -i\sigma^2 U_R^* \\ -i\sigma^2 D_R^* \end{pmatrix}, \tag{6}$$

where U_L and D_L are the left-handed techniup and techni-down, respectively, and U_R and D_R are the corresponding right-handed particles. Assuming the standard breaking to the maximal diagonal subgroup, the $SU(4)$ symmetry breaks spontaneously down to $SO(4)$. Such a breaking is driven by the following condensate

$$\langle Q_i^\alpha Q_j^\beta \epsilon_{\alpha\beta} E^{ij}\rangle = -2\langle \bar{U}_R U_L + \bar{D}_R D_L\rangle, \tag{7}$$

where the indices $i, j = 1, \ldots, 4$ denote the components of the tetraplet of Q, and the Greek indices indicate the ordinary spin. The matrix E is a 4×4 matrix defined in terms of the 2-dimensional unit matrix as

$$E = \begin{pmatrix} 0 & \mathbb{1} \\ \mathbb{1} & 0 \end{pmatrix}. \tag{8}$$

Following the notation of Wess and Bagger [20] $\epsilon_{\alpha\beta} = -i\sigma^2_{\alpha\beta}$ and $\langle U_L^\alpha U_R^{*\beta} \epsilon_{\alpha\beta}\rangle = -\langle \bar{U}_R U_L\rangle$. A similar expression holds for the D techniquark. The above condensate is invariant under an $SO(4)$ symmetry. The easiest way to check that an $SO(4)$ symmetry remains intact is by going to the following base

$$U_L = \frac{\lambda_1 + i\lambda_2}{\sqrt{2}}, \qquad \epsilon U_R^* = \frac{\lambda_1 - i\lambda_2}{\sqrt{2}},$$
$$D_L = \frac{\lambda_3 + i\lambda_4}{\sqrt{2}}, \qquad \epsilon D_R^* = \frac{\lambda_3 - i\lambda_4}{\sqrt{2}}, \tag{9}$$

where the λs are four independent two component spinors. In this base our condensate becomes simply

$$\langle \lambda_1^2 + \lambda_2^2 + \lambda_3^2 + \lambda_4^2\rangle, \tag{10}$$

which clearly is an $SO(4)$ invariant. Of the original $SU(4)$ global symmetry we are left with nine broken generators with associated Goldstone bosons.

In terms of the underlying degrees of freedom, and focusing only on the techniflavor symmetries, the nine Goldstone bosons transform like

$$\bar{D}_R U_L, \qquad \bar{U}_R D_L, \qquad \frac{1}{\sqrt{2}}(\bar{U}_R U_L - \bar{D}_R D_L), \tag{11}$$

for the three which will be eaten by the longitudinal components of the massive electroweak gauge bosons. The electric charge is, respectively, one, minus one and zero. For the other six Goldstone bosons we have

$$U_L U_L, \qquad D_L D_L, \qquad U_L D_L, \tag{12}$$

with the following electric charges

$$y + 1, \qquad y - 1, \qquad y, \tag{13}$$

together with the associated antiparticles. The last six Goldstone bosons (Eq. (12)) are di-technibaryons with opposite baryonic charge, one and minus one, respectively. The baryon number is a diagonal generator of $SU(4)$. As we already mentioned the choice of $y = 1$ makes one of the Goldstone bosons (namely the D) electrically neutral. We will explore the possibility of a neutral di-technibaryon as a component of cold dark matter in Sec. IV.

III. EFFECTIVE THEORIES

While the leptonic sector can be described within perturbation theory since it interacts only via electroweak interactions, the situation for the techniquarks is more involved since they combine into composite objects interacting strongly among themselves. It is therefore useful to construct low energy effective theories encoding the basic symmetry features of the underlying theory. We construct the linearly and nonlinearly realized low energy effective theories for our underlying theory. The theories we will present can be used to investigate relevant processes of interest at LHC and LC. It would be interesting to perform the analysis in [21] with these specific theories.

A. The linear realization

The relevant effective theory for the Higgs sector at the electroweak scale consists, in our model, of a light composite Higgs and nine Goldstone bosons. These can be assembled in the matrix

$$M = \left(\frac{\sigma}{2} + i\sqrt{2}\Pi^a X^a\right)E, \tag{14}$$

which transforms under the full $SU(4)$ group according to

$$M \to uMu^T, \quad \text{with} \quad u \in SU(4), \tag{15}$$

and X^a are the generators of the $SU(4)$ group which do not leave invariant the vacuum expectation value of M

$$\langle M\rangle = \frac{v}{2}E. \tag{16}$$

It is convenient to separate the 15 generators of $SU(4)$ into the six that leave the vacuum invariant (S^a) and the other nine that do not (X^a). One can show that the S^a generators of the $SO(4)$ subgroup satisfy the following relation

$$S^a E + E S^{aT} = 0 \quad \text{with} \quad a = 1, \ldots, 6. \tag{17}$$

The explicit realization of the generators is shown in Appendix A.

The electroweak subgroup can be embedded in $SU(4)$, as explained in detail in [22]. The main difference here is that we have a more general definition of the hypercharge. The electroweak covariant derivative is

$$D_\mu M = \partial_\mu M - ig[G_\mu M + M G_\mu^T], \tag{18}$$

with

$$G_\mu = \begin{pmatrix} W_\mu & 0 \\ 0 & -\frac{g'}{g}B_\mu^T \end{pmatrix} + \frac{y}{2}\frac{g'}{g}B_\mu \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (19)$$

We also have

$$W_\mu = W_\mu^a \frac{\tau^a}{2}, \qquad B_\mu^T = B_\mu \frac{\tau^{3^T}}{2} = B_\mu \frac{\tau^3}{2}, \quad (20)$$

where τ^a are the Pauli matrices. It is convenient to rewrite the gauge bosons in a more compact form

$$G = W^a L^a - \frac{g'}{g}B_\mu R^{3^T} + \sqrt{2}y\frac{g'}{g}B_\mu S^4, \quad (21)$$

with

$$L^a = \frac{S^a + X^a}{\sqrt{2}}, \qquad R^{aT} = \frac{X^a - S^a}{\sqrt{2}}, \quad \text{and} \quad a = 1, 2, 3. \quad (22)$$

With this gauging we are ensuring the correct pattern of electroweak symmetry breaking. In fact we can rewrite

$$G = G_S + G_X, \quad (23)$$

with

$$G_S = \frac{1}{\sqrt{2}}\sum_{a=1}^{3} S^a\left[W^a + \frac{g'}{g}B\delta_a^3\right] + \sqrt{2}y\frac{g'}{g}BS^4,$$
$$G_X = \frac{1}{\sqrt{2}}\sum_{a=1}^{3} X^a\left[W^a - \frac{g'}{g}B\delta_a^3\right]. \quad (24)$$

The generators satisfy the normalization conditions $\mathrm{Tr}[X^aX^b] = \delta^{ab}/2$, $\mathrm{Tr}[S^aS^b] = \delta^{ab}/2$ and $\mathrm{Tr}[SX] = 0$. Three of the Goldstone bosons, in the unitary gauge, are absorbed in the longitudinal degrees of freedom of the massive weak gauge bosons while the extra six Goldstone bosons will acquire a mass due to extended technicolor interactions as well as the electroweak interactions *per se*. Assuming a bottom up approach we will introduce by hand a mass term for the Goldstone bosons. The new Higgs Lagrangian is then

$$L = \frac{1}{2}\mathrm{Tr}[D_\mu M D^\mu M^\dagger] + \frac{m^2}{2}\mathrm{Tr}[MM^\dagger] - \frac{\lambda}{4}\mathrm{Tr}[MM^\dagger]^2 - \tilde{\lambda}\,\mathrm{Tr}[MM^\dagger MM^\dagger] - \frac{1}{2}\Pi_a(M_{\mathrm{ETC}}^2)^{ab}\Pi_b, \quad (25)$$

with $m^2 > 0$ and a and b running over the six uneaten Goldstone bosons. The matrix M_{ETC}^2 is dynamically generated and parametrizes our ignorance about the underlying extended technicolor model yielding the specific mass texture. The pseudo Goldstone bosons are expected to acquire a mass of the order of a TeV. Direct and computable contributions from the electroweak corrections break $SU(4)$ explicitly down to $SU(2)_L \times SU(2)_R$ yielding an extra contribution to the uneaten Goldstone bosons. However the main contribution comes from the ETC interactions.

The relation between the vacuum expectation value of the Higgs and the parameters of the present theory is

$$v^2 = \langle\sigma\rangle^2 = \frac{m^2}{\lambda + \tilde{\lambda}}. \quad (26)$$

Since in our theory we expect a light composite Higgs whose mass (in the broken phase) is $2m^2$ [23] this corresponds to a small overall self coupling. We have predicted in [6] a Higgs mass in the range $M_H \simeq 90$–150 GeV. By choosing the fiducial value 125 GeV and recalling that in our conventions we have $M_W = \frac{vg}{2}$, we then find

$$\lambda + \tilde{\lambda} \approx \frac{1}{8}, \quad \text{with} \quad v \approx 250 \text{ GeV}. \quad (27)$$

$\lambda + \tilde{\lambda}$ corresponds to the Higgs self coupling in the SM. It turns out that due to the presence of a light Higgs the associated sector can be treated perturbatively. We stress that the expectation of a light composite Higgs relies on the assumption that the quantum chiral phase transition as function of number of flavors near the nontrivial infrared fixed point is smooth and possibly of second order [24]. The composite Higgs Lagrangian is a low energy effective theory and higher dimensional operators will also be phenomenologically relevant.

B. The nonlinearly realized effective theory

One can always organize the low energy effective theory in a derivative expansion. The best way is to make use of the exponential map

$$U = \exp\left(i\frac{\Pi^a X^a}{F}\right)E, \quad (28)$$

where Π^a represents the 9 Goldstone bosons and X^a are the 9 generators of $SU(4)$ that do not leave the vacuum invariant (see Appendix A for an explicit realization of the group generators). To introduce the electroweak interactions one simply adopts the same covariant derivative used for the linearly realized effective theory, see Eqs. (18)–(24).

The associated nonlinear effective Lagrangian reads

$$L = \frac{F^2}{2}\mathrm{Tr}[D_\mu U D^\mu U^\dagger] - \frac{1}{2}\Pi_a(M_{\mathrm{ETC}}^2)^{ab}\Pi_b. \quad (29)$$

Still the mass squared matrix parametrizes our ignorance about the underlying ETC dynamics.

A common ETC mass for all the pseudo Goldstone bosons carrying baryon number can be provided by adding the following term to the previous Lagrangian

$$2C\,\mathrm{Tr}[UBU^\dagger B] + C = \frac{C}{4F^2}\sum_{i=1}^{6}\Pi_B^i\Pi_B^i, \quad (30)$$

with

$$B = \frac{1}{2\sqrt{2}}\begin{pmatrix} \mathbb{1} & 0 \\ 0 & -\mathbb{1} \end{pmatrix}. \quad (31)$$

Dimensional analysis requires $C \propto \Lambda_{\rm TC}^6/\Lambda_{\rm ETC}^2$. A similar term can be added to the linearly realized version of our theory. The previous naive estimate for C leads to phenomenologically too light new particles and we hence expect new ingredients to also contribute to these masses.

It is straightforward to add the vector meson sector to these theories, which would then allow to repeat the analysis performed in [21].

IV. THE DARK SIDE OF THE 5$^{\rm th}$ FORCE

We now provide a component for cold dark matter within our model. Such a candidate must be electrically and color neutral and have a mass above the current experimental exclusion limits.

According to the choice of the hypercharge there are two distinct possibilities. If we assume the SM-like hypercharge assignment for the techniquarks and the new lepton family, the new heavy neutrino can be an interesting dark matter candidate. For that, it must be made sufficiently stable by requiring no flavor mixing with the lightest generations and be lighter than the unstable charged lepton [6]. This possibility is currently under investigation [25]. However, we can also consider another possibility. We can choose the hypercharge assignment in such a way that one of the pseudo Goldstone bosons does not carry electric charge. The dynamics providing masses for the pseudo Goldstone bosons may be arranged in a way that the neutral pseudo Goldstone boson is the LTB. If conserved by ETC interactions the technibaryon number protects the lightest baryon from decaying. Since the mass of the technibaryons are of the order of the electroweak scale they may constitute interesting sources of dark matter. Some time ago in a pioneering work Nussinov [13] suggested that, in analogy with the ordinary baryon asymmetry in the Universe, a technibaryon asymmetry is a natural possibility. A new contribution to the mass of the Universe then emerges due to the presence of the LTB. It is useful to compare the fraction of technibaryon mass Ω_{TB} to baryon mass Ω_B in the Universe

$$\frac{\Omega_{TB}}{\Omega_B} = \frac{TB}{B}\frac{m_{TB}}{m_p}, \tag{32}$$

where m_p is the proton mass, m_{TB} is the mass of the LTB. TB and B are the technibaryon and baryon number densities, respectively. The net baryon number at a given Universe temperature T, in our conventions is:

$$B = \frac{n_B - n_{\bar{B}}}{gT^2/6}, \tag{33}$$

with $n_B(n_{\bar{B}})$ the (anti)baryon number density and g the degrees of freedom multiplicity. A similar expression holds for the net technibaryon number TB.

Knowing the distribution of dark matter in the galaxy earth based experiments can set stringent limits on the physical features of the dominant component of dark matter [15]. Such a distribution, however, is not known exactly [16] and it depends on the number of components and type of dark matter. In order to determine few features of our LTB particle we make the oversimplified approximation in which our LTB constitutes the whole dark matter contribution to the mass of the Universe. In this limit the previous ratio should be around 5 [26]. By choosing in our model the hypercharge assignment $y = 1$ the lightest neutral Goldstone boson is the state consisting of the DD techniquarks. The fact that it is charged under $SU(2)_L$ makes it detectable in Ge detectors [27].

It is well known that weak anomalies violate the baryon and the lepton number. More precisely, weak processes violate $B + L$, while they preserve $B - L$. Similarly, the weak anomalies violate also the technibaryon number, since technibaryons couple weakly. The weak technibaryon-, lepton- and baryon- number violating effects are highly suppressed at low temperatures while they are enhanced at temperatures comparable to the critical temperature of the electroweak phase transition where sphaleron processes are active (though sphaleron processes only occur below the scale of the electroweak phase transition) [28]. With T^* we define the temperature below which the sphaleron processes cease to be important. This temperature is not exactly known but it is expected to be in the range between 150–250 GeV [28].

Following early analysis [14,29] we have performed a careful computation of Ω_{TB}/Ω_B within our model. There are few differences with respect to the work in [14]. For example, our dark matter candidate is not a typical technibaryon whose mass is uniquely fixed by the underlying technicolor dynamics but a pseudo Goldstone boson whose mass is set by yet unspecified dynamics. The fact that the mass of our LTB is essentially a free parameter in the present model further differentiates it with respect to the original models in which the tehnibaryon mass is fixed by the underlying technicolor dynamics alone [14].

Imposing thermal and chemical equilibrium, the electric neutrality condition as well as the presence of a continuous electroweak phase transition (T^* now is below the critical temperature) we find:

$$\frac{TB}{B} = \frac{11}{36}\sigma_{TB}\left(\frac{m_{TB}}{T^*}\right), \tag{34}$$

with σ_{TB} the statistical weight function

$$\sigma_{TB}\left(\frac{m_{TB}}{T^*}\right) = \frac{3}{2\pi^2}\int_0^\infty dx x^2 \sinh^{-2}\left(\frac{1}{2}\sqrt{x^2 + \left(\frac{m_{TB}}{T^*}\right)^2}\right). \tag{35}$$

In the previous estimate the LTB is taken to be lighter then the other technibaryons and the new lepton number is violated. We have, however, considered different scenarios and various limits which will be reported in [30].

In order for the reader to have a better idea of how we arrived at the Eq. (34) we will now sketch the basic steps.

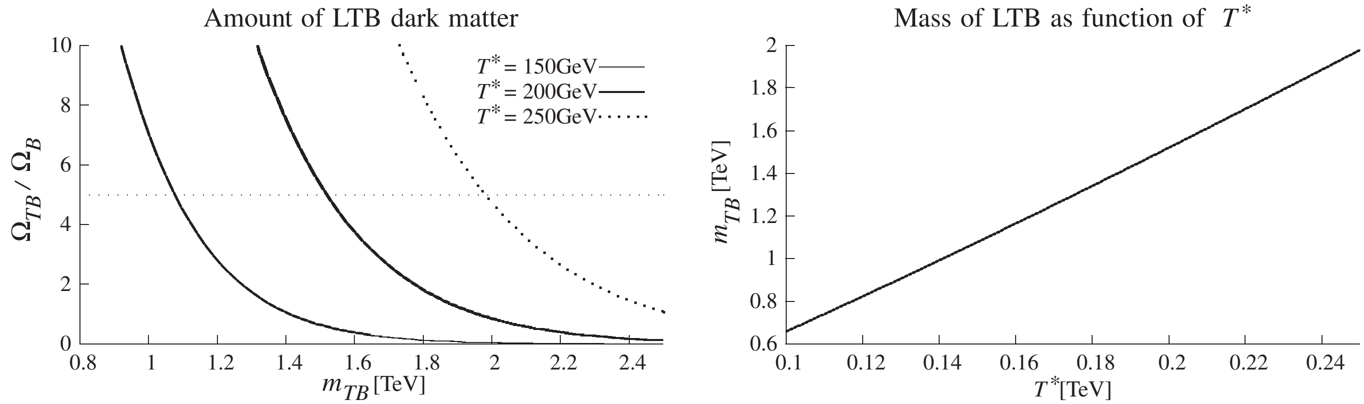

FIG. 1. *Left Panel*: The fraction of technibaryon matter density over the baryonic one as function of the technibaryon mass. The desired value of $\Omega_{TB}/\Omega_B \sim 5$ depends on the lightest technibaryon mass and the value of T^*. *Right Panel*: By requiring the correct amount ($\Omega_{TB}/\Omega_B \sim 5$) of dark matter we show the relation between the technibaryon mass and T^*.

First one introduces a chemical potential for each particle species. These chemical potentials depend on each other as long as the particles are in thermal equilibrium. This step, opportunely adapted for our model, is explained in [29]. The processes violating the technibaryon, baryon, and lepton numbers (i.e. the sphaleron processes) are also relevant, yielding in practice extra relations among the various chemical potentials. These processes are active also just below the electroweak phase transition. In the end one finds a sufficient number of relations among the various chemical potentials which leads to the expression (34). Another important point is that we have not specified the common origin of both asymmetries, i.e. the baryonic and technibaryonic one, this is reflected in the fact that we can only relate the two after electroweak symmetry breaking. But since the baryon number is *experimentally* known this is sufficient to determine the other. Our basic results are shown in Fig. 1. The desired value of the dark matter fraction in the Universe can be obtained for a LTB mass of the order of a TeV for quite a wide range of values of T^*. The only free parameter in our analysis is essentially the mass of the LTB which is ultimately provided by ETC interactions.

If we now consider the case of a different number of technicolors always in the adjoint representation of the gauge group, except for the change of nonuniversal quantities such as T^*, our analysis is not modified. The present idea and model computations can be used to the physical situations in which a new strongly interacting theory emerges at the electroweak scale whose underlying new fermions couple to the weak interactions, and there exist new composite baryonic like objects arising as Goldstone bosons of a new global symmetry group broken by other nonstandard model like interactions.

V. CONCLUSIONS

Imminent experiments will shed light on the electroweak breaking sector of the standard model. We have constructed the effective theories associated to a strong fifth force, of technicolor type, responsible to the breaking of the electroweak theory. These can be used for studying signatures and processes which are relevant for the upcoming experiments such as LHC. Interestingly one of the neutral pseudo Goldstone bosons is a natural candidate for a sizable component of cold dark matter, and we have shown that it is possible to ascribe the whole dark matter in the Universe to the LTB of our theory, given that it has a mass of ~1 TeV.

ACKNOWLEDGMENTS

We thank S. Bolognesi, D. D. Dietrich, A. Jokinen, K. Petrov, K. Rajagopal and K. Tuominen for careful reading of the manuscript, discussions and useful comments. The work of C. K. and F. S. is supported by the Marie Curie Excellence Grant under contract No. MEXT-CT-2004-013510. F. S. is also supported by the Danish Research Agency.

APPENDIX A: GENERATORS

It is convenient to use the following representation of $SU(4)$

$$S^a = \begin{pmatrix} \mathbf{A} & \mathbf{B} \\ \mathbf{B}^\dagger & -\mathbf{A}^T \end{pmatrix}, \qquad X^i = \begin{pmatrix} \mathbf{C} & \mathbf{D} \\ \mathbf{D}^\dagger & \mathbf{C}^T \end{pmatrix}, \tag{A1}$$

where A is Hermitian, C is Hermitian and traceless, $B = -B^T$ and $D = D^T$. The S are also a representation of the $SO(4)$ generators, and thus leave the vacuum invariant $S^a E + E S^T = 0$. Explicitly, the generators read

$$S^a = \frac{1}{2\sqrt{2}} \begin{pmatrix} \tau^a & \mathbf{0} \\ \mathbf{0} & \tau^{aT} \end{pmatrix}, \qquad a = 1, \ldots, 4, \tag{A2}$$

where $a = 1, 2, 3$ are the Pauli matrices and $\tau^4 = \mathbb{1}$. These are the generators for $SU_V(2) \times U_V(1)$.

$$S^a = \frac{1}{2\sqrt{2}}\begin{pmatrix} \mathbf{0} & \mathbf{B}^a \\ \mathbf{B}^{a\dagger} & \mathbf{0} \end{pmatrix}, \qquad a = 5, 6, \tag{A3}$$

with

$$B^5 = \tau^2, \qquad B^6 = i\tau^2. \tag{A4}$$

The rest of the generators which do not leave the vacuum invariant are

$$X^i = \frac{1}{2\sqrt{2}}\begin{pmatrix} \tau^i & \mathbf{0} \\ \mathbf{0} & \tau^{iT} \end{pmatrix}, \qquad i = 1, 2, 3, \tag{A5}$$

and

$$X^i = \frac{1}{2\sqrt{2}}\begin{pmatrix} \mathbf{0} & \mathbf{D}^i \\ \mathbf{D}^{i\dagger} & \mathbf{0} \end{pmatrix}, \qquad i = 4, \ldots, 9, \tag{A6}$$

with

$$\begin{aligned} &D^4 = \mathbb{1}, \qquad D^6 = \tau^3, \qquad D^8 = \tau^1, \\ &D^5 = i\mathbb{1}, \qquad D^7 = i\tau^3, \qquad D^9 = i\tau^1. \end{aligned} \tag{A7}$$

The generators are normalized as follows

$$\mathrm{Tr}[S^a S^b] = \mathrm{Tr}[X^a X^b] = \frac{1}{2}\delta^{ab}, \qquad \mathrm{Tr}[X^i S^a] = 0. \tag{A8}$$

---

[1] S. Weinberg, Phys. Rev. D **19**, 1277 (1979); L. Susskind, Phys. Rev. D **20**, 2619 (1979).
[2] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990).
[3] F. Sannino and K. Tuominen, Phys. Rev. D **71**, 051901 (2005).
[4] N. Evans and F. Sannino, hep-ph/0512080.
[5] D. D. Dietrich, F. Sannino, and K. Tuominen, Phys. Rev. D **73**, 037701 (2006)
[6] D. D. Dietrich, F. Sannino, and K. Tuominen, Phys. Rev. D **72**, 055001 (2005).
[7] F. Sannino, Int. J. Mod. Phys. A **20**, 6133 (2005).
[8] N. D. Christensen and R. Shrock, Phys. Lett. B **632**, 92 (2006).
[9] T. Appelquist and F. Sannino, Phys. Rev. D **59**, 067702 (1999).
[10] R. Sundrum and S. D. H. Hsu, Nucl. Phys. **B391**, 127 (1993).
[11] D. K. Hong and H. U. Yee, hep-ph/0602177.
[12] M. Harada, M. Kurachi, and K. Yamawaki, hep-ph/0509193.
[13] S. Nussinov, Phys. Lett. **165B**, 55 (1985).
[14] S. M. Barr, R. S. Chivukula, and E. Farhi, Phys. Lett. B **241**, 387 (1990).
[15] D. S. Akerib *et al.* (CDMS Collaboration), Phys. Rev. Lett. **96**, 011302 (2006).
[16] C. P. Ma and E. Bertschinger, Astrophys. J. **612**, 28 (2004).
[17] M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189**, 575 (1981).
[18] A. G. Cohen and H. Georgi, Nucl. Phys. **B314**, 7 (1989).
[19] E. Witten, Phys. Lett. **117B**, 324 (1982).
[20] J. Wess and J. Bagger, *Supersymmetry and Supergravity* (Princeton Series in Physics, Princeton, 1992).
[21] A. R. Zerwekh, hep-ph/0512261.
[22] T. Appelquist, P. S. Rodrigues da Silva, and F. Sannino, Phys. Rev. D **60**, 116007 (1999); Z. y. Duan, P. S. Rodrigues da Silva, and F. Sannino, Nucl. Phys. **B592**, 371 (2001).
[23] Note that if one assumes a strongly coupled linear sigma model the relation between the physical mass and the mass parameter in the theory is no longer linear and important modifications are expected [6].
[24] We have provided supporting arguments for this picture in [6] where the reader will find also a more general discussion of this issue and possible pitfalls.
[25] K. Tuominen (private communication).
[26] O. Lahav and A. R. Liddle, astro-ph/0601168.
[27] J. Bagnasco, M. Dine, and S. D. Thomas, Phys. Lett. B **320**, 99 (1994).
[28] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).
[29] J. A. Harvey and M. S. Turner, Phys. Rev. D **42**, 3344 (1990).
[30] S. B. Gudnason, C. Kouvaris, and F. Sannino (unpublished).