Dark matter from new technicolor theories

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

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

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We investigate dark matter candidates emerging in recently proposed technicolor theories. We determine the relic density of the lightest, neutral, stable technibaryon having imposed weak thermal equilibrium conditions and overall electric neutrality of the Universe. In addition we consider sphaleron processes that violate baryon, lepton and technibaryon number. Our analysis is performed in the case of a first order electroweak phase transition as well as a second order one. We argue that, in both cases, the new technibaryon contributes to the dark matter in the Universe. Finally we examine the problem of the constraints on these types of dark matter components from earth based experiments.

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I. INTRODUCTION

The origin of the dark matter is one of the most intriguing open problems of modern particle physics and cosmology. A few decades ago, astronomers realized that it was impossible to explain the motion of galaxies, just by accounting the luminous part of the galaxy. The most important observational evidence comes from the rotation curves of spiral galaxies. Astronomers are able to estimate the velocities of clouds of hydrogen atoms, just by looking at the Doppler shifted 21 cm emission. Naively one would expect by using Newton's law, that the velocity of clouds like these should fall as $\propto r^{-1/2}$, when r becomes larger than the radius of the luminous part of the galaxy (r being the distance to the center of the galaxy). However, astronomers observe an almost constant velocity, independent of r. Unless one proposes a radical change of the laws of gravity or of motion, it is logical to assume that matter that does not interact "too much" and therefore appears dark to us, has to fill the space of the galaxy.

Once we assume that a type of dark matter (DM) is responsible for the discrepancies of the motion of the galaxies, then there are two distinct possibilities for the origin of the DM. The first type of candidates are of baryonic origin and they are called MACHO (Massive Compact Halo Objects). The MACHO can be brown dwarf stars, giant planets and massive black holes. Massive black holes with masses near $100M_{\odot}$ can be remnants at the center of supernovae explosions. The brown dwarf stars are stellar objects with masses less than $0.1M_{\odot}$. Since a proton star needs at least a mass of $0.1M_{\odot}$ to ignite nuclear fusion, brown dwarfs never get to begin the nuclear fusion of hydrogen. Giant planets can be also a component of MACHOs with masses of the order of $0.001M_{\odot}$. However, observations so far showed with a high level of confidence that MACHOs cannot account for more than 20% of the DM [1].

The second possibility is to have matter of nonbaryonic origin. In contrast with MACHOs that are compact objects of baryonic matter, in this case we have particles that are neutral and interact only through gravitational and weak forces. The name WIMP (Weakly Interacting Massive Particles) is frequently used for these particles. The Standard Model does not have the particles with the desired properties. Ordinary neutrinos are the only electrically neutral objects that interact weakly within the SM. They are, however, too light and would compose part of a *hot*-type dark matter, which is usually not considered in a viable cosmological model. This is so since hot dark matter would smear out large scale structure of galaxies. Supersymmetry, for example, provides some natural WIMP candidates for cold dark matter such as neutralinos.

In the search of WIMP candidates, particles related to technicolor theories were also investigated as possible sources of cold DM. This idea was pioneered by Nussinov in [2,3] and further investigated in [4]. He imagined that DM could be accounted by a technibaryon asymmetry which can be ultimately related to the ordinary baryon asymmetry. Although it was clear that technicolor can give "reliable" DM candidates, little interest was shown in the past because of the severe problems that most technicolor theories suffer from such as large flavor changing neutral current processes and/or problems with the Electroweak Precision Measurements. Progress in this direction has been made recently. We have constructed explicit technicolor extensions of the SM passing the precision tests [5–8]. Our results are further supported by ADS/CFT inspired model computations [9–11]. Because of their walking nature they have very much reduced flavor changing neutral current processes. Therefore it does make sense to revisit the possibility of technibaryons as components of DM. We have started this analysis in [8] for the models with technifermions in the adjoint representation of the technicolor gauge group. In this case the technibaryon is one of the would-be Goldstone bosons of the underlying technicolor theory. In Ref. [12] has been also suggested

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

that technicolor theories may lead to dark matter candidates of similar nature.

The goal of the present work is threefold: First, we will provide the basic set up while showing the detailed computations needed to determine the present relic density of our technibaryon DM candidate. Second, we consider two different orders of the electroweak phase transition and compare the results. Third, using our low energy effective theory developed in [8] we compute the relevant cross sections needed to examine the problem of the constraints on these types of DM components from earth based experiments.

In the next section we review the basic properties of the new technicolor theory and present the lightest neutral technibaryon (LTB) of the theory. In Sec. III we present our main calculation of the technibaryon contribution to DM. In Sec. IV we investigate the experimental constraints and comment on the detection of our technibaryon WIMP. Finally we briefly conclude in Sec. V.

II. CONVENTIONS AND NOTATIONS

The minimal walking technicolor model [5–8] has two techniflavors (techni-up U, and techni-down D) transforming according to the adjoint representation of the $SU(2)$ technicolor gauge group. The global flavor symmetry is $SU(4)$ which breaks spontaneously to $SO(4)$. The associated low energy effective theories have been constructed in [8]. There are 9 Goldstone bosons, 3 of which are eaten by the W and Z bosons and are the technicolor equivalent of the ordinary pions. Three of the remnant six Goldstone bosons transform under techniflavor symmetries as:

$$U_L U_L, \qquad D_L D_L, \qquad U_L D_L, \tag{1}$$

with electric charges, respectively

$$y + 1, \qquad y - 1, \qquad y, \tag{2}$$

while the other three are the antiparticles. In the following we drop the subscript L when referring to the above states. The parameter y can take any real value. It is related to the hypercharge of the techniquarks and is such that gauge anomalies cancel out. Additionally in order to cancel Witten's global anomaly we simply add an extra family of leptons

$$\mathcal{L}_L = \begin{pmatrix} \nu' \\ \zeta \end{pmatrix}_L, \qquad (\nu'_R, \zeta_R), \tag{3}$$

with hypercharges:

$$-\frac{3y}{2}, \qquad \left(\frac{-3y+1}{2}, \frac{-3y-1}{2}\right), \tag{4}$$

where we use the convention:

$$Q = T_3 + Y. \tag{5}$$

A typical cold DM candidate must be electrically neutral and at most have weak interactions. For example we can choose $y = 1/3$, i.e. the SM-like hypercharge assignment and in this case ν' (the *new neutrino*) is electrically neutral and can be made stable by requiring no mixing with the lighter neutrino species. This case is similar to the one studied in [13]. In this paper we consider the case with $y = 1$, where the second (technibaryon) Goldstone boson of Eq. (1) is now electrically neutral. One the other hand, the *new leptons* ν' and ζ have electric charges -1 and -2, respectively. The still not directly observed Goldstone bosons will acquire masses via new interactions. If we assume these interactions to preserve the technibaryon number, then the electrically neutral DD, if it is the lightest technibaryon, is stable. We will denote with LTB the lightest neutral technibaryon particle. We conclude that DD is an interesting candidate which can be a cold DM component.

III. COMPUTING THE LTB RELIC DENSITY

We now explicitly compute the DD-type boson relic density in the case it is neutral and stable. We impose thermal equilibrium and overall electric neutrality for the matter in the Universe. Imposing overall electrical neutrality avoids the huge energetic Coulomb costs due to electric fields of the otherwise uncanceled charges in the Universe. In addition to the theoretical reasons, observations confirm an overall neutrality. Thermal equilibrium occurs among different particles as long their rate Γ of interaction is much larger than the the expansion rate of the Universe H, where H is the Hubble constant. If $H > \Gamma$ at a given time the particles decouple from each other and hence can no longer be in thermal equilibrium.

At some energy scale higher than the electroweak one, following the work of Nussinov [2], we assume the existence of a mechanism leading to a technibaryon asymmetry in the Universe. Given that the technibaryon and baryon number have a very similar nature such an asymmetry is very plausible and can have a common origin. Here we will not speculate further on the origin of the (techni)baryon asymmetry but will relate it to the observed baryon asymmetry as done by Nussinov as well as in [4]. Here we provide detailed computations for our specific technicolor model for two types of electroweak phase transitions.

Even if one is able to produce an asymmetry above or around the electroweak scale the (techni)baryon number is spoiled by quantum anomalies. Fortunately although the baryon (B), technibaryon (TB), lepton (L) number and the new lepton number for the new lepton family (L') are not conserved individually, their differences, i.e. $B - L$ and $3TB - L$ are preserved. This fact allows for a nonzero (techni)baryonic asymmetry to survive. The processes leading to such a violation are termed "sphaleron" processes and at the present time are negligible. However these processes were active during the time the Universe had a temperature above or at the scale of the electroweak symmetry breaking (~ 250 GeV). Indeed these processes were rapid enough to thermalize baryons, leptons and technibaryons. At some point as the Universe expands and its temperature falls, the baryon-lepton-technibaryon violation processes cease to be significant. The precise value of this temperature T^* depends on the underlying theory driving electroweak symmetry breaking. Within the SM framework and assuming the validity of the semiclassical calculation of the tunneling effect [14], T^* has been estimated by equating the rate of the sphaleron processes to H. According to [14], T^* satisfies the following equation

$$T^* = \frac{2M_W(T^*)}{\alpha_W \ln(\frac{M_{Pl}}{T^*})} B\left(\frac{\lambda}{\alpha_W}\right), \tag{6}$$

where M_W is the mass of the W bosons, M_{Pl} is the Planck scale, α_W is the weak coupling constant, λ is the self-coupling of the Higgs boson and $B(\lambda/\alpha_W)$ is a function that takes on values from 1.5 to 2.7 as the ratio λ/α_W goes from zero to infinity [14,15]. As we already mentioned this formula is an approximation and it depends on the not very well known ratio of λ/α_W. According to what is the value of this ratio, T^* can vary within the 150–250 GeV range. In technicolor theories, since the Higgs is a composite object, the self-coupling λ is in principle calculable. An estimation $\lambda = 1/8$ for our specific model was given in [7]. Since $\alpha_W = 1/29$ (or a bit larger at the electroweak scale), the ratio λ/α_W gives a T^* around 200 GeV.

It is time to introduce now the chemical potentials for the relevant particle species. We here follow Ref. [16]

μ_W	for W^-	μ_{dL}	for d_L, s_L, b_L
μ_0	for ϕ^0	μ_{dR}	for d_R, s_R, b_R
μ_-	for ϕ^-	μ_{iL}	for e_L, μ_L, τ_L
μ_{uL}	for u_L, c_L, t_L	μ_{iR}	for e_R, μ_R, τ_R
μ_{uR}	for u_R, c_R, t_R	$\mu_{\nu iR}$	for $\nu_{eR}, \nu_{\mu R}, \nu_{\tau R}$
$\mu_{\nu iL}$	for $\nu_{eL}, \nu_{\mu L}, \nu_{\tau L}$		

where the indices L and R denote chirality. We have a common chemical potential for the up, charm and top quarks, and a different one for the other triplet of down, strange and bottom. A common chemical potential has to do with the fact that at the scale of interest QCD interactions put quarks of the same charge on equal footing. We introduce a different chemical potential for all of the leptons. Also in order to be as general as possible we have assumed the existence of right handed neutrinos and introduced different chemical potentials for the left and the right handed particles. The thermal equilibrium conditions associated to the weak interactions read:

$$\mu_W = \mu_- + \mu_0 \qquad (W^- \leftrightarrow \phi^- + \phi^0), \tag{7}$$

$$\mu_{dL} = \mu_{uL} + \mu_W \qquad (W^- \leftrightarrow \bar{u}_L + d_L), \tag{8}$$

$$\mu_{iL} = \mu_{\nu iL} + \mu_W \qquad (W^- \leftrightarrow \bar{\nu}_{iL} + e_{iL}), \tag{9}$$

$$\mu_{\nu iR} = \mu_{\nu iL} + \mu_0 \qquad (\phi^0 \leftrightarrow \bar{\nu}_{iL} + \nu_{iR}), \tag{10}$$

$$\mu_{uR} = \mu_0 + \mu_{uL} \qquad (\phi^0 \leftrightarrow \bar{u}_L + u_R), \tag{11}$$

$$\mu_{dR} = -\mu_0 + \mu_W + \mu_{uL} \qquad (\phi^0 \leftrightarrow d_L + \bar{d}_R), \tag{12}$$

$$\mu_{iR} = -\mu_0 + \mu_W + \mu_{\nu iL} \qquad (\phi^0 \leftrightarrow e_{iL} + \bar{e}_{iR}), \tag{13}$$

where it is understood that the Higgs is a composite of two techniquarks. The Goldstone bosons of Eq. (1) are gauged under the weak symmetry and hence we introduce the following chemical potential for these Goldstone bosons and the new lepton family of Eq. (3)

$\mu_{\zeta L}$	for ζ_L	μ_{UU}	for UU
$\mu_{\zeta R}$	for ζ_R	μ_{UD}	for UD
$\mu_{\nu' L}$	for $\nu_{\zeta L}$	μ_{DD}	for DD
$\mu_{\nu' R}$	for $\nu_{\zeta R}$		

The corresponding thermal equilibrium equations for the extra particles introduced by the technicolor theory *per se* are

$$\mu_{\zeta L} = \mu_W + \mu_{\nu' L} \qquad (\zeta_L \leftrightarrow W^- + \nu_{\zeta L}), \tag{14}$$

$$\mu_{UD} = \mu_{DD} - \mu_W \qquad (DD \leftrightarrow UD + W^-), \tag{15}$$

$$\mu_{UU} = \mu_{UD} - \mu_W = \mu_{DD} - 2\mu_W \qquad (UD \leftrightarrow UU + W^-), \tag{16}$$

$$\mu_{\zeta R} = -\mu_0 + \mu_{\zeta L} \qquad (\phi^0 \leftrightarrow \zeta_L + \bar{\zeta}_R), \tag{17}$$

$$\mu_{\nu' R} = \mu_0 + \mu_{\nu' L} \qquad (\phi^0 \leftrightarrow \bar{\nu}_{\zeta L} + \nu_{\zeta R}), \tag{18}$$

where Eq. (15) has been used in Eq. (16).

Each classical gauge and scalar field configuration with a given topological number leads to a simultaneous jump for *all* of the anomalous charges. Hence each quark-doublet generation, lepton-doublet generation, the new lepton family number as well as techniquark number are violated by the same classical field configuration. The one loop anomalous coefficient dictates the relative amount of the jump for each anomalous charge when turning on a given classical field configuration.

With the normalization of $1/3$ for the technibaryonic charge for our techniquarks and $1/3$ for the ordinary quark-baryonic charge of the quarks, 1 for all of the leptons, the simplest classical configuration with one unit of topological charge will induce a transition from the the vacuum of the theory to a state containing one baryon (per each generation), one lepton (for each generation), a technibaryonlike object with 3 technibaryons and one new lepton. The relation among the chemical potentials emerging from the above is:

$$3(\mu_{u_L} + 2\mu_{d_L}) + \mu + \frac{1}{2}\mu_{UU} + \mu_{DD} + \mu_{\nu'} = 0. \quad (19)$$

The parameter μ is defined as $\sum_i \mu_{\nu iL} \equiv \mu$. We have assumed that the difference in the baryon number between two different quark-doublet generations is created identical before the electroweak phase transition. A similar relation will be assumed for the lepton charges. Note that the difference is not affected by the weak anomaly and hence will not be generated later on.

We can now turn to the calculation of number densities. The difference between the number densities of particles and their corresponding antiparticles is given by

$$\begin{aligned} n &= n_+ - n_- \\ &= g\int \frac{d^3k}{(2\pi)^3}\frac{1}{z^{-1}e^{E\beta} - \eta} - g\int \frac{d^3k}{(2\pi)^3}\frac{1}{ze^{E\beta} - \eta}, \end{aligned} \quad (20)$$

where n_+ and n_- are the number densities of the particles and antiparticles, respectively. The constant g is the multiplicity of the degrees of freedom (spin, for example), $\beta = 1/T$ in units $k_B = 1$, and η takes on the values 1 and -1 for bosons and fermions, respectively. The fugacity $z = e^{\mu\beta}$ and E is the energy. The ratio μ/T in the Universe is sufficiently small that we can Taylor expand the above relation. The number density now can be written as

$$n = \begin{cases} gT^3\frac{\mu}{T}f(\frac{m}{T}) & \text{for fermions,} \\ gT^3\frac{\mu}{T}g(\frac{m}{T}) & \text{for bosons,} \end{cases} \quad (21)$$

where the functions f and g are defined as follows

$$f(z) = \frac{1}{4\pi^2}\int_0^\infty dx x^2 \cosh^{-2}\left(\frac{1}{2}\sqrt{x^2 + z^2}\right), \quad (22)$$

$$g(z) = \frac{1}{4\pi^2}\int_0^\infty dx x^2 \sinh^{-2}\left(\frac{1}{2}\sqrt{x^2 + z^2}\right). \quad (23)$$

We now differentiate two cases according to the order of the electroweak phase transition. In the case of a second order or weak first order electroweak phase transition we expect that the temperature T^* is below the temperature of the phase transition. This means that baryon, lepton and technibaryon violating processes persist after the phase transition. The second possibility is to have a strong first order phase transition where the violating processes freeze right at the phase transition. We are going to examine separately the two different cases.

Assuming that the violating processes persist even after the phase transition, we need to impose two conditions here: Electric neutrality and set $\mu_0 = 0$, since the Higgs boson condenses and the electroweak symmetry breaks spontaneously. Recall that we can introduce a nonzero chemical potential only for unbroken symmetries whose generators commute with all of the gauge ones. Here the Higgs boson is a composite particle, made of techni-up and techni-down quarks $(\bar{U}U + \bar{D}D)/\sqrt{2}$. Therefore when we refer to μ_0 as the chemical potential, we mean the chemical potential of the composite object.

From Eq. (21) we see that the number densities, in the leading approximation, are linear in the chemical potential for small μ/T. For convenience we express the baryon number density as

$$B \equiv \frac{n_B - n_{\bar{B}}}{gT^2/6}. \quad (24)$$

We shall use the same normalization (dividing the number density by $gT^2/6$) also for the lepton number, technibaryon number etc. Since in the end we care only for ratios of number densities, the normalization constant cancels out.

We conveniently define the function σ as follows

$$\sigma_i = \begin{cases} 6f(\frac{m_i}{T^*}) & \text{for fermions,} \\ 6g(\frac{m_i}{T^*}) & \text{for bosons,} \end{cases} \quad (25)$$

where f and g are those of Eq. (23) and (24), respectively, and the index i refers to the particle in question.

For all of the SM particles the statistical function is taken to be 1 and 2 for massless fermions and bosons, respectively, except for the top quark which we treat massive as m_t is of order T^*. The reason why we can take the other SM particles to be massless in the statistical function is that $m \ll T^*$. However, the technibaryons as well as the particles of the new lepton family have masses that cannot be ignored. We should emphasize that we calculate the baryon and lepton numbers at the temperature T^* where sphalerons die out.

The baryon density can be written as

$$\begin{aligned} B &= \frac{3}{3}[(2 + \sigma_t)(\mu_{uL} + \mu_{uR}) + 3(\mu_{dL} + \mu_{dR})] \\ &= (10 + 2\sigma_t)\mu_{uL} + 6\mu_W + (\sigma_t - 1)\mu_0, \end{aligned} \quad (26)$$

where Eqs. (8), (11), and (12) have been used and the factor in the first line includes number of colors and that the baryon number of each quark is $1/3$. The factor 3 of the down-type quarks is the number of families and equivalent the factor $2 + \sigma_t$ is the number of families taking into account the top mass effect.

Similarly the lepton number for the Standard Model leptons is

$$L = \sum_i (\mu_{\nu iL} + \mu_{\nu iR} + \mu_{iL} + \mu_{iR}) = 4\mu + 6\mu_W. \quad (27)$$

For the new lepton family we have

$$\begin{aligned} L' &= \sigma_\zeta(\mu_{\zeta L} + \mu_{\zeta R}) + \sigma_{\nu'}(\mu_{\nu' L} + \mu_{\nu' R}) \\ &= 2(\sigma_{\nu'} + \sigma_\zeta)\mu_{\nu' L} + 2\sigma_\zeta\mu_W + (\sigma_{\nu'} - \sigma_\zeta)\mu_0. \end{aligned} \quad (28)$$

Similarly for the technibaryons we get

$$TB = \frac{2}{3}(\sigma_{UU}\mu_{UU} + \sigma_{UD}\mu_{UD} + \sigma_{DD}\mu_{DD}) = \frac{2}{3}(\sigma_{UU} + \sigma_{UD} + \sigma_{DD})\mu_{DD} - \frac{2}{3}(\sigma_{UD} + 2\sigma_{UU})\mu_W. \tag{29}$$

The charge constraint for all the particles is

$$Q = \frac{2}{3}\cdot 3(2+\sigma_t)(\mu_{uL} + \mu_{uR}) - \frac{1}{3}\cdot 3\cdot 3(\mu_{dL} + \mu_{dR}) - \sum_i(\mu_{iL} + \mu_{iR}) - 2\cdot 2\mu_W - 2\mu_- + 2\sigma_{UU}\mu_{UU} + \sigma_{UD}\mu_{UD} - 2\sigma_\zeta(\mu_{\zeta L} + \mu_{\zeta R}) - \sigma_{\nu'}(\mu_{\nu' L} + \mu_{\nu' R}). \tag{30}$$

For the first order phase transition we will need also the neutrality with respect to the weak isospin charge which is

$$Q_3 = \frac{3(2+\sigma_t)}{2}\mu_{uL} - \frac{3\cdot 3}{2}\mu_{dL} + \frac{1}{2}\sum_{i=1}^{3}(\mu_{\nu iL} - \mu_{iL}) - 4\mu_W - (\mu_0 + \mu_-) + (\sigma_{UU}\mu_{UU} - \sigma_{DD}\mu_{DD}) + \frac{1}{2}(\sigma_{\nu'}\mu_{\nu' L} - \sigma_\zeta\mu_{\zeta L}). \tag{31}$$

The need for the isospin neutrality condition, in the first order case, comes from the fact that we are computing our final relic densities above the electroweak phase transition where the weak isospin is unbroken.

Since it is not clear whether the electroweak phase transition is first or second order, we should examine both cases. It is expected as in [16] that a strong first order phase transition occurs fast enough to "freeze" the baryon and technibaryon violating processes just at the transition. In this case one calculates the equilibrium conditions just before the transition. On the other hand, in a second order phase transition we expect the violating processes to persist below the phase transition and the equilibrium conditions are imposed after the phase transition. If the phase diagram as function of temperature and density of our technicolor theory would be known a specific order of the electroweak phase transition would be used.

When the ratio between the number densities of the technibaryons to the baryons is determined we have

$$\frac{\Omega_{TB}}{\Omega_B} = \frac{3}{2}\frac{TB}{B}\frac{m_{TB}}{m_p}, \tag{32}$$

here m_{TB} is the mass of the LTB (the m_{DD}) and m_p is the mass of the proton.

Note that a possible mixing between the new family and an ordinary standard model family would dilute the relative ν' abundance and eventually annihilate L'.

A. 2nd Order Phase Transition

Here the two conditions we have to impose are: Overall electrical neutrality and $\mu_0 = 0$ for the chemical potential of the Higgs boson. The ratio between the number density of the technibaryons to the baryons can be expressed as function of the L/B and L'/B ratios. In order to provide a simple and compact expression we consider the limiting case in which the UU and UD technibaryons are substantially heavier than the DD companion, the top is light with respect to the electroweak phase transition temperature and the new lepton family is degenerate, i.e. $\sigma_\zeta = \sigma_{\nu'}$. In this approximation the ratio simplifies to

$$-\frac{TB}{B} = \frac{\sigma_{DD}}{3(18+\sigma_{\nu'})}\left[(17+\sigma_{\nu'}) + \frac{(21+\sigma_{\nu'})}{3}\frac{L}{B} + \frac{2}{3}\frac{(9+5\sigma_{\nu'})}{\sigma_{\nu'}}\frac{L'}{B}\right]. \tag{33}$$

The results of the calculation are summarized in Fig. 1. This figure shows what are the allowed values of the parameter ξ defined below, as a function of the mass of the LTB, for a given T^*, if the LTB accounts for the whole dark matter density of the universe. The parameter ξ can be considered roughly speaking as the total ratio of lepton over baryon number, with the new lepton family number L' weighted "appropriately" due to the large mass that ν' and ζ carry.

$$\xi \equiv \frac{L}{B} + \frac{2}{\sigma_{\nu'}}\frac{9+5\sigma_{\nu'}}{21+\sigma_{\nu'}}\frac{L'}{B}. \tag{34}$$

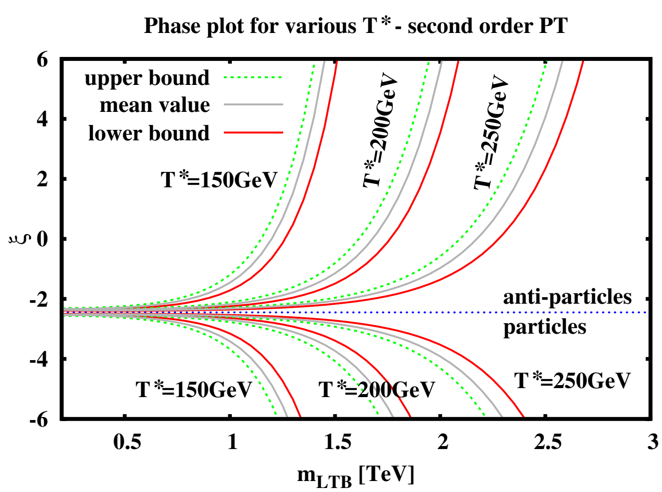

FIG. 1 (color online). Plot representing the region of the parameters according to which the fraction of technibaryon matter density over the baryonic one takes on the values [3.23, 5.55]. We consider a second order phase transition. The parameters in the plot are the mass of the LTB DM particle and ξ of Eq. (34). The plot includes various values of T^*. The dotted line separates areas of abundant particles and antiparticles.

From Fig. 1 we see, for example, that if we set $L' = 0$ (no new leptons present) while also setting $L/B = -4$, we need a mass for the LTB somewhere between 1.1 to 2.2 TeV, according to what is the freeze out temperature T^*. We should emphasize that there are two branches of allowed values for ξ, separated by the dotted horizontal line. The lower branch, as, for example, the one we just described with $\xi = -4$, corresponds to a relic density made by technibaryons DD. The upper set of allowed values, (as for $\xi = 2$), corresponds to the DD antiparticle. In Fig. 3 we show the dependence of the neutral technibaryon matter density as a function of its mass for a fixed value of the parameter ξ. We see that if the LTB mass is lighter than roughly a TeV the density of the particle is very large, giving a too large ratio Ω_{TB}/Ω_B. So, for a given value of ξ and T^*, WMAP data put constraints on the allowed mass of the technibaryon. On the other hand if we increase enough the mass of the technibaryon, we can get a ratio less than 4–5, which means that the technibaryon can be a component of the dark matter density.

B. 1st Order Phase Transition

If the electroweak phase transition is predicted to be first order, then the baryon, lepton and technibaryon violating processes freeze slightly above the phase transition. For this reason, we have to impose two conditions, the overall charge neutrality $Q = 0$ and $Q_3 = 0$, where Q_3 is the charge associated with the T_3 isospin generator of the weak interactions. This charge has to be zero because above the phase transition the electroweak symmetry is not broken and therefore $Q_3 = 0$.

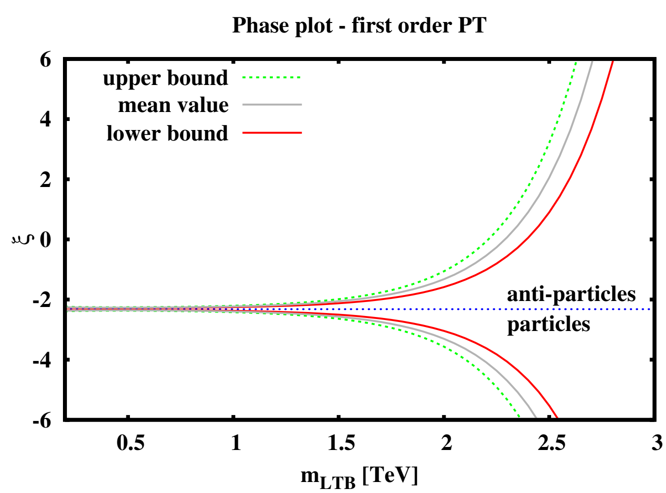

FIG. 2 (color online). Plot representing the region of the parameters according to which the fraction of technibaryon matter density over the baryonic one takes on the values [3.23, 5.55]. Here we consider the case of a first order phase transition. The parameters in the plot are the mass of the LTB DM particle and ξ of Eq. (36). The dotted line separates areas of abundant particles and antiparticles.

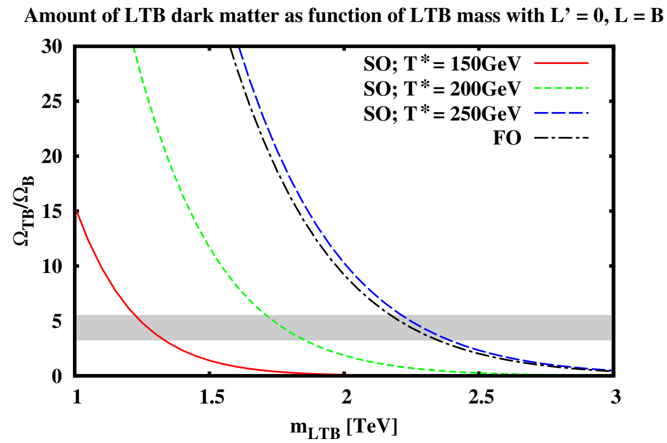

FIG. 3 (color online). Amount of LTB DM as function of the mass of the LTB particle. The plot is shown for $L' = 0$ and $L = B$ for second order (SO) phase transitions with various temperatures T^* and a for first order (FO) phase transition as well.

The technibaryon over baryon number density ratio is, in the same approximation used for the second order phase transition:

$$-\frac{TB}{B} = \sigma_{DD}\frac{22 + \sigma_{\nu'}}{9(22 + 2\sigma_{DD} + \sigma_{\nu'})}\left[3 + \frac{L}{B} + \frac{1}{\sigma_{\nu'}}\frac{L'}{B}\right]. \tag{35}$$

T^* is expected to be larger than that of the second order case, i.e. it should be identified with the critical temperature T^c of the electroweak phase transition. This fact forces the mass of the LTB to be larger than that of the second order case to describe the whole DM. Our results are summarized in Fig. 2. As in the case of the 2nd order phase transition, we have plotted the allowed values of the ξ parameter

$$\xi \equiv \frac{L}{B} + \frac{1}{\sigma_{\nu'}}\frac{L'}{B}, \tag{36}$$

which is slightly different from the previous case, as a function of the LTB mass, under the WMAP constraints regarding the overall density of dark matter in the universe. Using the previous example of $L' = 0$ and $L/B = -4$, we get an LTB mass of around 2.2 TeV.

IV. DETECTION OF THE NEUTRAL TECHNIBARYON

Apart from the possibility of detecting a technibaryon in the LHC experiment it would certainly be interesting to detect the neutral technibaryon in earth based experiments such as the CDMS [17–20]. Two are the basic ingredients affecting the detection of a cold DM object in these kinds of experiments. The first one is how large is the cross section of the object to be observed with the matter in the detector. The second has to do with the local density of

DM in general and of the specific component of DM, in particular. Current estimates suggest that the local density for a single component should be somewhere between 0.2–0.4 $\mathrm{GeV/cm^3}$. It is evident that the higher the cross section and the local density of dark matter are, the larger are the chances for the detection of the particle. The CDMS collaboration, for example, can identify WIMPs by observing the recoil energy produced in elastic scattering between the WIMP and a nucleus in the detector. The expected rates of events per unit time and mass of the detector has been calculated in several places and we refer to the review paper [21] for a complete list of relevant references. The number of counts reported by the detector per unit time, mass of the detector and recoil energy is

$$\frac{dR}{dT} = \frac{R_0}{E_0 r} e^{-T/E_0 r}, \tag{37}$$

where T is the recoil energy of the nucleus, E_0 is the kinetic energy of the WIMP and $r = 4mM_n/(m + M_n)^2$, m and M_n being the masses of the WIMP and the nucleus, respectively. The parameter R_0 is the total rate containing the information about the cross section and is given by

$$R_0 = \frac{2}{\pi^{1/2}} \frac{N_0}{A} \frac{\rho_{dm}}{m} \sigma_0 v_0, \tag{38}$$

where N_0 is the Avogadro number, A is the atomic number of the nucleus of the detector, ρ_{dm} is the local dark matter density, σ_0 is the cross section for an elastic collision between the WIMP and the nucleus and v_0 is the thermal velocity of the WIMPs. One should note here that Eq. (37) is an approximate expression. In reality the calculation is more elaborate. For example, in principle one has to assume a Maxwell distribution for the velocities of the WIMPs up to the escape velocity for our galaxy. In addition, the effect of the motion of the earth relatively to the halo should be considered. These factors can change the expected rate. The total rate of counts can be more usefully rewritten in convenient units as

$$R_0 = \frac{503}{M_n m} \left(\frac{\sigma_0}{1\ \mathrm{pb}}\right) \left(\frac{\rho_{dm}}{0.4\ \mathrm{GeV\,cm^{-3}}}\right) \left(\frac{v_0}{230\ \mathrm{kms^{-1}}}\right) \frac{\mathrm{GeV}^2}{\mathrm{kg.days}}. \tag{39}$$

Since our prospective dark matter component is a Goldstone boson, we are interested only in the spin-independent elastic cross section. This is given in natural units by [22]

$$\sigma_0 = \frac{G_F^2}{2\pi} \mu^2 \bar{Y}^2 \bar{N}^2 F^2, \tag{40}$$

where G_F is the Fermi constant and $\bar{Y} = 2Y$. For a Dirac fermion $\bar{Y} = Y_L + Y_R$ and μ is the reduced mass of the latter and the nucleus target. $\bar{N} = N - (1 - 4\sin^2\theta_w)Z$, where N and Z are the number of neutrons and protons in the target nucleus and θ_w is the Weinberg angle. The parameter F^2 is a form factor for the target nucleus. The cross section can be written as

$$\sigma_0 = 8.431 \times 10^{-3} \frac{\mu^2}{\mathrm{GeV}^2} \bar{Y}^2 \bar{N}^2 F^2\ \mathrm{pb}. \tag{41}$$

The Ge atom has 41 neutrons and 32 protons, giving an $\bar{N} = 38.59$. Our LTB has $\bar{Y} = 1$ [23]. Since Standard Model neutrinos have $\bar{Y} = 1/2$, this means that the cross section for the technibaryon will be 4 times larger than the one corresponding to a heavy neutrino. As we already mentioned, for typical values of the L/B ratio, in order to get the whole density for the dark matter, the mass of the technibaryon should be of the order of a TeV. The form factor F^2 for the nucleus of Ge depends on the recoil energy T. It models the loss of coherence of the scattering for large recoil energies. For typical values of the recoil energy around 20–50 keV one expects F^2 to be around 0.58. We estimated the nuclear form factor F using the solid sphere approximation -proper for the spin-independent WIMP interaction- summarized in [21]. To be more precise the nuclear form factor ranges from 0.72 to 0.43 when the recoil energy ranges from 20 to 50 keV [24].

The number of counts that are detectable is given by

$$\text{counts} = \frac{dR}{dT} \Delta T \times \tau, \tag{42}$$

where τ is the exposure of the detector measured in kg.days and ΔT is the energy resolution of the detector. In the CDMS experiment a 19.4 kg.days exposure was achieved for the Ge detectors with an energy resolution of $\Delta T = 1.5$ keV. So far no counts have been found. The 90% level of confidence would lead to 2.3 counts.

If we assume that our LTB constitutes the whole DM in the Universe we have seen from our previous computations that a typical value of the mass is about 2 TeV, for the second order phase transition case. Taking a recoil energy around 50 keV, $\rho_{dm} = 0.3\ \mathrm{GeV/cm^3}$ and $F^2 \sim 0.43$ the number of counts predicted is around 13 which is a value few times larger than the 90% confidence value presented before. By stretching the parameters we can reduce, or even annihilate the gap, between our prediction and experimental bounds. Using still a mass around 2 TeV but choosing a different set of input, i.e. $\rho_{dm} = 0.1\ \mathrm{GeV/cm^3}$, $F^2 \sim 0.3$ and $T = 70$ keV one finds around two predicted counts. Hence we would be within the 90% confidence level. Under these rather extreme conditions one cannot yet completely exclude the possibility that our WIMP can constitute the whole DM. Another simple way to reduce the gap between experiment and our LTB particle, if we imagine it to be a component of DM, is to increase its mass. In doing so, however, we neglect the relevant information gained in the previous sections in which we related the mass of the LTB to the fraction of DM in the Universe it can account for.

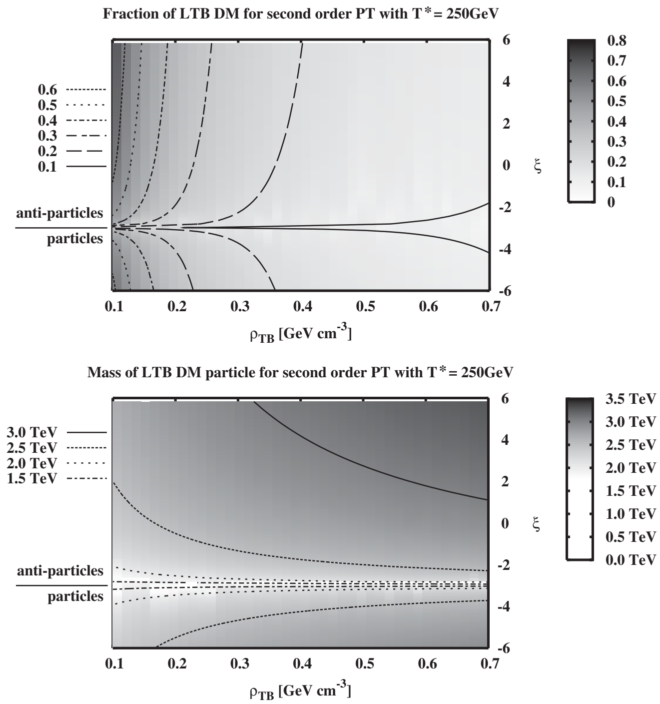

FIG. 4. *Top Panel*: The maximal fraction of local DM density allowed by the 90% experimental constraint as function of the local DM density and the parameter ξ of Eq. (34). *Bottom Panel*: For the corresponding maximal fraction of local DM density currently allowed by the 90% experimental constraint as function of the local DM density and ξ we plot the associated LTB mass. Both plots are presented with second order phase transition with $T^* = 250$ GeV and a recoil energy $T = 50$ keV.

We now take into account, in a more careful way, such a dependence on the mass of the LTB. From the previous section we learned that the general trend is that the amount of DM saturated by our LTB object decreases when increasing the mass of the LTB. In the absence of a complete computation of how DM distributes itself in the Universe we make the oversimplifying assumption that the fraction of local DM density of our LTB follows the same fraction of DM in the Universe. At this point we impose the 90% experimental constraint. Our results are reported in Fig. 4. In the figure we have chosen $F^2 \simeq 0.43$ and the thermal velocity is 230 km/s. We present both the maximal fraction of local DM density determined imposing the 90% experimental constraint (Fig. 4) and the associated value of the LTB mass as functions of ξ. We have allowed for variations of the parameters to make our analysis more complete. Note that we have allowed the local DM density to reach, in the plots, very large values although a more modest range (i.e. up to 0.4 GeV/cm^3) is probably sufficient.

Summarizing we can say that for reasonable values of the input parameters the 90% experimental constraint allows for a 10% to 65% of LTB-DM component in the Universe. Our DM component allowed mass ranges between 1.4 and 3.3 TeV depending on the order of the associated electroweak phase transition as well as the exact value of the local DM density and experimental parameters range. Our conclusions are slightly affected if we use 20 keV as recoil energy. In this case at most we can account for 30% of the DM in the Universe and the masses of the LTB would be slightly higher.

The question to be answered at this point is: What makes the rest of the DM in the Universe? We speculate that a techni-axion, needed for the solution of the strong *CP* problem, could be a natural candidate (see for example [25]). In this way the two components for DM are associated to two natural and complementary extensions of the SM. An explicit model containing axions from technicolor-like dynamics has been constructed in [26].

V. CONCLUSIONS

We have investigated in much detail dark matter candidates emerging in our recently proposed technicolor theories. We have determined the contribution of the lightest, neutral, stable technibaryon to the dark matter density having imposed weak thermal equilibrium conditions and overall electric neutrality of the Universe. Sphaleron processes have been taken into account. We performed the analysis in the case of a first order electroweak phase transition as well as a second order one. In both cases, the new technibaryon contributes to the dark matter in the Universe. We have also examined the problem of the constraints from earth based experiments. We find that quite a substantial amount of DM can be explained within our model. The new generation of DM-detection experiments can put very strict limits or even rule out the present type of DM component. We should stress that our framework can be applied to any model featuring a new baryonic type particle at the electroweak scale whose new baryon-type charge is violated only by weak interactions.

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[1] C. Alcock *et al.* (MACHO Collaboration), Astrophys. J. **542**, 281 (2000).
[2] S. Nussinov, Phys. Lett. B **165**, 55 (1985).
[3] S. Nussinov, Phys. Lett. B **279**, 111 (1992).
[4] S. M. Barr, R. S. Chivukula, and E. Farhi, Phys. Lett. B **241**, 387 (1990).
[5] F. Sannino and K. Tuominen, Phys. Rev. D **71**, 051901 (2005).
[6] D. K. Hong, S. D. H. Hsu, and F. Sannino, Phys. Lett. B **597**, 89 (2004).
[7] D. D. Dietrich, F. Sannino, and K. Tuominen, Phys. Rev. D **72**, 055001 (2005).
[8] S. B. Gudnason, C. Kouvaris, and F. Sannino, Phys. Rev. D **73**, 115003 (2006).
[9] D. K. Hong and H. U. Yee, Phys. Rev. D **74**, 015011 (2006).
[10] J. Hirn and V. Sanz, Phys. Rev. Lett. **97**, 121803 (2006).
[11] M. Piai, hep-ph/0608241.
[12] T. Appelquist, M. Piai, and R. Shrock, Phys. Rev. D **69**, 015002 (2004).
[13] G. Rybka and P. Fisher, hep-ex/0507086.
[14] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. B **155**, 36 (1985).
[15] F. R. Klinkhamer and N. S. Manton, Phys. Rev. D **30**, 2212 (1984).
[16] J. A. Harvey and M. S. Turner, Phys. Rev. D **42**, 3344 (1990).
[17] D. S. Akerib *et al.* (CDMS Collaboration), Phys. Rev. Lett. **93**, 211301 (2004).
[18] V. Mandic, FERMILAB Report No. FERMILAB-THESIS-2004-53 (unpublished).
[19] G. s. Wang, FERMILAB Report No. FERMILAB-THESIS-2005-09 (unpublished).
[20] D. S. Akerib *et al.* (CDMS Collaboration), Phys. Rev. Lett. **96**, 011302 (2006).
[21] J. D. Lewin and P. F. Smith, Astropart. Phys. **6**, 87 (1996).
[22] M. W. Goodman and E. Witten, Phys. Rev. D **31**, 3059 (1985).
[23] We have directly computed this value for $\bar{Y}$ using the effective Lagrangian in [8]. For a generic y we have $\bar{Y} = 2y - 1$ and it coincides with the value one deduces by constructing the LTB wave function as follows: $\sqrt{2}|LTB\rangle = |D_L D_L\rangle - |D_R D_R\rangle$.
[24] We may have overestimated the nuclear form factor [18]. If the physical value of F^2 is lower than the one we used then the allowed fraction of LTB-DM increases.
[25] H. Davoudiasl, hep-ph/0607111.
[26] S. D. H. Hsu and F. Sannino, Phys. Lett. B **605**, 369 (2005).