We investigate the simplest models where baryon and lepton numbers are defined as local symmetries spontaneously broken at the low scale and discuss the implications for cosmology. We define the simplest anomaly-free theory for spontaneous baryon and lepton number violation which predicts the existence of lepto-baryons. In this context we study the new sphaleron condition on the chemical potentials and show the relation between the present baryon asymmetry and the $B - L$ asymmetry generated in the early universe. The properties of the cold dark matter candidate for which stability is a natural consequence from symmetry breaking are briefly discussed.

I. INTRODUCTION

The lack of new physics signals at the Large Hadron Collider questions most of the popular extensions of the Standard Model. There are appealing theoretical arguments in favor of physics at the TeV scale, but the current status of the field tells us that maybe the naturalness argument related to the Higgs mass is not a good guide for the search of new physics at current colliders.

One of the bonuses of the Standard Model is the conservation of the global baryonic and leptonic symmetries at the classical level, which forbids proton decay. Unfortunately, when high-dimension operators consistent with gauge symmetries of the Standard Model are considered, these global symmetries are lost. Then, one needs to postulate a suppression mechanism to be consistent with the experimental bounds on the proton lifetime. Typically, one postulates the existence of the great desert and the new scale is related to the grand unified scale, i.e. $M_{GUT} \sim 10^{14-16}$ GeV.

Recently, the authors of Ref. [4] proposed an alternative approach for physics beyond the Standard Model where the baryonic and leptonic symmetries are considered fundamental local symmetries. In this context these symmetries can be spontaneously broken near the electroweak scale and there is hope to test the predictions at colliders. The original idea related to the spontaneous breaking of local baryon number was proposed in Ref. [5], where the author speculated about the possibility to use the Higgs mechanism in these models. In the theories investigated in Refs. [4, 6, 7] the local baryon number is always broken in three units and the proton is stable. There are also interesting implications for cosmology since typically in
the new sector of the theory one has a candidate for the cold dark matter in the universe for which stability is automatic and the usual picture for baryogenesis can be different [8, 9].

In this article, we revisit the idea of constructing a simple theory where baryon and lepton numbers are local symmetries. We find that the simplest phenomenologically viable model that is free of anomalies consists of only four fermion multiplets, which we call “lepto-baryons”. We investigate the main features of this theory and discuss the main implications for cosmology. The details of the new sphaleron conditions on the chemical potentials when baryon number is broken at the low scale are discussed in details. The new fermions in the theory have the same quantum numbers, after symmetry breaking, as the gauginos and Higgsinos in the Minimal Supersymmetric Standard Model. In this context the dark matter candidate is the lightest Majorana fermion with baryon number.

II. LOCAL BARYON AND LEPTON NUMBERS

In this section, we argue that the simplest model with gauged baryon and lepton numbers requires four fermion multiplets to cancel all the anomalies. Subsequently, we will explore the implications for cosmology. Here we follow the discussion in Ref. [4] to understand the anomaly cancellation. Our main goal is to construct a theory based on the gauge group

\[ G_{BL} = SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_B \otimes U(1)_L. \]

To simplify the discussion we explain how to gauge the baryon number, and later we give the results in the case where lepton number is also a local symmetry. In order to define the baryon number as a local symmetry one has to introduce a new set of fermions to cancel the non-trivial baryonic anomalies

\[ A_1^{SM}(SU(2)_L^2 \otimes U(1)_B) = 3/2 \text{ and } A_2^{SM}(U(1)_Y^2 \otimes U(1)_B) = -3/2. \]

The others anomalies are automatically zero in the Standard Model, i.e.

\[ A_3^{SM}(SU(3)_C \otimes U(1)_B) = 0, A_4^{SM}(U(1)_Y \otimes U(1)_B) = 0, A_5(U(1)_B)^{SM} = 0, \text{ and } A_6(U(1)_B)^{SM} = 0. \]

One also has to make sure that the Standard Model anomalies still vanish once we modify the theory. In order to understand the possible solutions for anomaly cancellations we need to look for solutions which allow for vector-like masses for the new fields. Therefore, typically the simplest solutions correspond to the cases when we have an even number of colorless fields.
Let us start discussing the case where we add just two new fields

\[ \Psi_L \sim (1, N, Y_1, B_1) \text{ and } \Psi_R \sim (1, M, Y_2, B_2). \]

The anomaly \( A_5(U(1)_B) = 0 \) requires \( B_1 = MB_2/N \). Now, using \( A_0(U(1)_B^3) = 0 \) the condition \( M = N \) holds and hence \( B_1 = B_2 \). Then \( A_1(SU(2)_L^2 \otimes U(1)_B) \) cannot be canceled. Thus, there is no solution with just two new fermionic fields. Next, we consider the set of three fields which cancel all the anomalies

\[ \Phi_L \sim (1, N, y, B_1), \quad \Phi_R \sim (1, N, y, B_2) \quad \text{and} \quad \eta_L \sim (1, M, 0, B_3), \]

where \( M = 2N, y^2 = N^2/4 \) and \( B_1 = -B_2 = -B_3 = 3/N^3 \). For example one can have the fields

\[ \Phi_L \sim (1, 2, 1/2, 3/2), \quad \Phi_R \sim (1, 2, 1, -3/2), \quad \text{and} \quad \eta_L \sim (1, 4, 0, -3/8). \]

Unfortunately, this model predicts the existence of fermionic fields with fractional electric charge which cannot decay into the Standard Model fields. Therefore, one always has a stable charged field which is ruled out by cosmology. Now, we are ready to look for a solution with only four fields.

### A. Minimal Theory of Lepto-Baryons

It is easy to show that there is an interesting solution with only four extra fermionic representations plus right-handed neutrinos

\[ \Psi_L \sim (1, 2, 1/2, 3/2, 3/2), \quad \Psi_R \sim (1, 2, 1/2, -3/2, -3/2), \]
\[ \Sigma_L \sim (1, 3, 0, -3/2, -3/2), \quad \text{and} \quad \chi_L \sim (1, 1, 0, -3/2, -3/2). \]

which cancel all anomalies including the leptonic anomalies. Here we show the baryon and lepton numbers for each field. Notice that in the Standard Model the leptonic anomalies are

\[ A_7^{SM}(SU(2)_L^2 \otimes U(1)_L) = 3/2 \text{ and } A_8^{SM}(U(1)_Y^2 \otimes U(1)_L) = -3/2. \]

The others anomalies are automatically zero in the Standard Model with right-handed neutrinos, i.e.

\[ A_9^{SM}(SU(3)_C \otimes U(1)_L) = 0, A_{10}^{SM}(U(1)_Y \otimes U(1)_L^2) = 0, A_{11}(U(1)_L)^{SM} = 0, \text{ and } A_{12}(U(1)_L^2)^{SM} = 0, \]
and in general one also must cancel the mixed anomalies

$$\mathcal{A}_{13}(U(1)^2_B \otimes U(1)_L) = 0, \mathcal{A}_{14}(U(1)_B \otimes U(1)^2_L) = 0, \text{ and } \mathcal{A}_{15}(U(1)_Y \otimes U(1)_B \otimes U(1)_L) = 0.$$ 

We call these fields “lepto-baryons” and in component form can be expressed as

$$\Psi_L = \begin{pmatrix} \psi_1^+ \\ \psi_0 \\ \psi_1 \end{pmatrix}, \quad (\Psi^c)_L = (\Psi_R)^c = \begin{pmatrix} \psi_2^0 \\ \psi_2 \end{pmatrix}, \quad \text{and } \Sigma_L = \frac{1}{2} \begin{pmatrix} \Sigma^0 & \sqrt{2} \Sigma^+ \\ \sqrt{2} \Sigma^- & -\Sigma^0 \end{pmatrix}.$$

The relevant interactions in this model are given by

$$-\mathcal{L} \supset h_1 \bar{\Psi}_R H \chi L + h_2 H^\dagger \Psi_L \chi L + h_3 H^\dagger \Sigma L \Psi_L + h_4 \bar{\Psi}_R \Sigma L H$$

$$+ \lambda_\Psi \bar{\Psi}_R \Sigma L S_B^* + \lambda_\chi \chi L \lambda L S_B + \lambda_\Sigma \text{Tr} \Sigma L^2 S_B$$

$$+ Y_\nu \ell_L H \nu^c + \lambda_R \nu^c \nu^c S_L + \text{h.c.},$$

where $\nu^c = (\nu_R)^c$ are the right-handed neutrinos. The needed scalar sector is composed of the fields

$$S_B \sim (1, 1, 0, 3, 3), \quad S_L \sim (1, 1, 0, 0, 2) \quad \text{and} \quad H \sim (1, 2, 1/2, 0, 0).$$

We define the vacuum expectation values as $\langle H \rangle = v/\sqrt{2}$, $\langle S_B \rangle = v_B/\sqrt{2}$ and $\langle S_L \rangle = v_L/\sqrt{2}$. In order to understand the implications for baryogenesis we will assume that $U(1)_L$ is broken at a scale much larger than the electroweak-scale. When $S_B$ acquires a vacuum expectation value, the local baryonic symmetry $U(1)_B$ is broken to the $Z_2$ symmetry which guarantees the stability of the lightest field with baryon number. This will be our dark matter candidate.

After symmetry breaking the mass matrix for the neutral fermions in the basis $(\psi_1^0, \psi_2^0, \Sigma^0, \chi^0)$ is given by

$$\mathcal{M}_0 = \begin{pmatrix}
0 & \frac{1}{\sqrt{2}} \lambda_\Psi v_B & -\frac{1}{2 \sqrt{2}} h_3 v & \frac{1}{\sqrt{2}} h_2 v \\
\frac{1}{\sqrt{2}} \lambda_\Psi v_B & 0 & -\frac{1}{2 \sqrt{2}} h_4 v & \frac{1}{\sqrt{2}} h_1 v \\
-\frac{1}{2 \sqrt{2}} h_3 & -\frac{1}{2 \sqrt{2}} h_4 v & 0 & \frac{1}{\sqrt{2}} \lambda_\Sigma v_B \\
\frac{1}{\sqrt{2}} h_2 v & -\frac{1}{\sqrt{2}} h_1 v & 0 & \sqrt{2} \lambda_\chi v_B
\end{pmatrix},$$

(5)
and the mass matrix for the charged fermions reads as

\[
\mathcal{M}_\pm = \begin{pmatrix}
- \frac{1}{\sqrt{2}} \lambda \psi v_B & \frac{1}{2} h_3 v \\
- \frac{1}{2} h_4 v & \frac{1}{\sqrt{2}} \lambda \Sigma v_B
\end{pmatrix},
\]

in the basis \((\psi^+_1, \Sigma^+)\) and \((\psi^-_2, \Sigma^-)\). Notice that when all Yukawa couplings \(h_i\) are small the new fermions have pure vector-like masses proportional to the vacuum expectation value of \(S_B\). In this case the model can easily satisfy all possible constraints. Notice that these fields have the same quantum numbers as the gaugino-Higgsino fields in the minimal supersymmetric Standard Model.

As it has been pointed out in Ref. [4], this type of model is interesting because the symmetry breaking scale can be low and there is no need to postulate the existence of the great desert between the electroweak and Planck scales. The main reason is that the scalar field that breaks local baryon number has baryonic charge three, \(S_B \sim (1, 1, 0, 3, 3)\), and never generates any contribution to proton decay. The allowed higher-dimension operators are highly suppressed. For example, one can write down the operator

\[
\frac{c}{\Lambda^{15}} (QQQL)^3 S_B^*,
\]

which mediates the processes: \(ppp \rightarrow e^+e^+e^+, \ ppn \rightarrow e^+e^+\bar{\nu}, \ pnn \rightarrow e^+\bar{\nu}\bar{\nu},\) and \(3n \rightarrow 3\bar{\nu}\). Unfortunately, these processes are highly suppressed due to the large power \(\Lambda^{15}\), even when \(\Lambda\) is around the TeV scale. It is important to mention that the vacuum expectation value, \(v_B\), defines the mass of the leptophobic gauge boson in the theory and the bounds on the mass of this type of gauge bosons are weak. See Refs. [10–12] for details.

### III. BARYON ASYMMETRY, SPHALERONS AND DARK MATTER

In this section, we investigate the relation between the baryon asymmetry and the \(B - L\) asymmetry solving the chemical equilibrium equations. Here we will assume that the local leptonic symmetry is broken far above the electroweak scale. The baryon asymmetry is given by

\[
B_f = \frac{n_q - n_{\bar{q}}}{s} = \frac{15}{4\pi^2 g_s T} 3 \left( \mu_{u_L} + \mu_{u_R} + \mu_{d_L} + \mu_{d_R} \right),
\]

when the chemical potential is \(\mu \ll T\), \(g_s\) is the total number of relativistic degrees of freedom and \(s\) is the entropy density.

The only conserved global symmetry after symmetry breaking in the theory is the usual \(B - L\) in the
Standard Model and we will use it to determine the final baryon asymmetry. The $B - L$ asymmetry is defined by

$$\Delta(B - L)_{\text{SM}} = \frac{15}{4\pi^2 g_s T} 3(\mu_{u_L} + \mu_{d_L} + \mu_{d_R} - \mu_{\nu_L} - \mu_{\nu_R}).$$

(9)

Assuming isospin conservation one has the conditions [13] on the chemical potentials

$$\mu_{u_L} = \mu_{d_L}, \quad \mu_{e_L} = \mu_{\nu_L}, \quad \text{and} \quad \mu_0 = \mu_+. \quad (10)$$

The Standard Model interactions for quarks and leptons with the Higgs give us the useful relations

$$\mu_{u_R} = \mu_0 + \mu_{u_L}, \quad \mu_{d_R} = -\mu_0 + \mu_{u_L}, \quad \text{and} \quad \mu_{e_R} = -\mu_0 + \mu_{e_L}. \quad (11)$$

The new interactions in our model proportional to $\lambda_i$ give the following new relations

$$2\mu_{\chi_L} + \mu_{SB} = 0, \quad 2\mu_{\Sigma_L} + \mu_{SB} = 0 \quad \text{and} \quad -\mu_{\Psi_L} + \mu_{\Psi_R} + \mu_{SB} = 0, \quad (12)$$

while the $h_i$ interactions impose the conditions

$$-\mu_{\Psi_R} + \mu_0 + \mu_{\chi_L} = 0, \quad -\mu_0 + \mu_{\Psi_R} + \mu_{\chi_L} = 0,$$

$$-\mu_0 + \mu_{\Sigma_L} + \mu_{\Psi_L} = 0, \quad \text{and} \quad -\mu_{\Psi_R} + \mu_{\Sigma_L} + \mu_0 = 0. \quad (13)$$

The sphaleron condition must conserve the total baryon number. Because the lepto-baryons carry $SU(2)_L$ quantum numbers, they are expected to contribute to the effective ’t Hooft operator. In Ref. [14], the ’t Hooft vertex for arbitrary representations was derived iteratively. Here we present a pedestrian approach to arrive at the same result. The basic idea is to integrate the anomaly equation to determine the change in the number of quanta corresponding to an instanton of unit winding number. Consider a left-handed fermion transforming as an isospin-$j$ representation under $SU(2)_L$. The anomaly equation reads

$$\partial_{\mu} j^\mu = \frac{g^2}{16\pi^2} A^{ab} W^a_{\mu} \tilde{W}^{\mu \nu b}, \quad (14)$$

where $A^{ab} = \text{Tr}[G_2 \{ T^a_j, T^b_j \}]$. The current carries a normalizing abelian charge $G$ which will drop out at the end of the computation. Upon inserting the $SU(2)$ identity $\{ T^a_j, T^b_j \} = \frac{2}{3} j(j + 1) \delta^{ab} \mathbf{1}_{2j+1}$, and
performing the trace, we have

$$\partial_\mu J^\mu = \frac{g^2}{16\pi^2} G \frac{1}{3} j(j+1)(2j+1) \delta^{ab} W^a_{\mu\nu} \tilde{W}^{\mu\nu b}. \quad (15)$$

We next express the field strength tensors in terms of matrix-valued fields by writing the Kronecker delta as a trace over a product of $SU(2)$ generators in the fundamental representation: $\delta^{ab} = 2 \text{Tr}(T^a_{1/2} T^b_{1/2})$. We then have

$$\frac{\partial \rho_G}{\partial t} + \nabla \cdot J = G \frac{2}{3} j(j+1)(2j+1) \left( \frac{g^2}{16\pi^2} \text{Tr}[W_{\mu\nu} \tilde{W}^{\mu\nu}] \right). \quad (16)$$

Inserting the explicit form of the instanton into the right hand side, we may integrate over the entire 4-volume. In the left hand side, we get the change of $G$-charge, and on the right hand side, the quantity inside the parenthesis yields the winding number, which in this case is unity. Upon performing the integration we arrive at

$$\frac{\Delta G}{G} = \Delta n_L = \frac{2}{3} j(j+1)(2j+1), \quad (17)$$

which is the change in the number of left-handed quanta corresponding to an instanton event, in agreement with Ref. [14]. The ’t Hooft vertex for a given multiplet is schematically of the form: $\mathcal{L}_{\text{eff}} = (\bar{\psi}_L) \Delta n_L$. Therefore, in this model the ’t Hooft operator is

$$(QQQL)^3 \bar{\Psi}_L \Psi_L \Sigma^4_L. \quad (18)$$

Notice that this condition is quite different from the one in the Standard Model. Hence, the relevant equilibrium condition imposed by the sphaleron processes is

$$3(3\mu_{uL} + \mu_{eL}) + \mu_{\Psi L} - \mu_{\Psi R} + 4\mu_{\Sigma L} = 0. \quad (19)$$

Finally, the conservation of electric charge implies

$$6(\mu_{uL} + \mu_{uR}) - 3(\mu_{dL} + \mu_{dR}) - 3(\mu_{eL} + \mu_{eR}) + 2\mu_0 + \mu_{\Psi L} + \mu_{\Psi R} = 0, \quad (20)$$

while the total baryon number conservation above the symmetry breaking scale gives us

$$3(\mu_{uL} + \mu_{uR} + \mu_{dL} + \mu_{dR}) + \frac{3}{2}(2\mu_{\Psi L} - 2\mu_{\Psi R} - 4\mu_{\Sigma L} - \mu_{\chi L}) + 6\mu_{S_B} = 0. \quad (21)$$
Therefore, we have the following four equations

\[ \Delta (B - L)_{\text{SM}} = \frac{15}{4\pi^2 g_* T} (12\mu_{uL} - 9\mu_{eL} + 3\mu_0), \]  

\[ Q_{em} = 0 \implies 3\mu_{uL} - 3\mu_{eL} + 8\mu_0 = 0, \]  

\[ B_T = 0 \implies \mu_{uL} - 2\mu_{\chi L} = 0, \]  

\[ \text{Sphalerons} \implies 9\mu_{uL} + 3\mu_{eL} + 2\mu_{\chi L} = 0. \]

We can express the final baryon asymmetry as a function of the initial \( B - L \) asymmetry

\[ B_f = \frac{15}{4\pi^2 g_* T} (12\mu_{uL}) = \frac{32}{99} \Delta (B - L)_{\text{SM}} \approx 0.32 \Delta (B - L)_{\text{SM}}. \]

Notice that the conversion factor is different from the case where one has the usual Standard Model sphaleron condition, i.e. smaller than \( \frac{28}{79} \approx 0.35 \) [13]. As it is well-known this conversion factor is important to understand how the initial \( B - L \) asymmetry is transferred into the baryon asymmetry without assuming any particular model. Notice that in general there is no reason to expect a simple solution for the conversion factor in this model because the sphalerons conserve the total baryon number.

In this model the lightest new field with baryon number is automatically stable. Here we discuss briefly how to explain the observed relic density in the Universe in the case when the dark matter candidate is the \( \chi_L \) field, with mass approximately equal to \( \sqrt{2\lambda_{\chi} v_B} \). The simplest dark matter candidate in this model is a Majorana fermion \( \chi \). After symmetry breaking the three CP-even physical Higgses can have interactions with all Standard Model fields and the dark matter candidate. Therefore, we can have the following annihilation channels,

\[ \chi \chi \rightarrow H_i \rightarrow \bar{q}q, hh, WW, ZZ, \ldots \]

From the interaction with the leptophobic gauge boson, the following additional channel

\[ \chi \chi \rightarrow Z_B \rightarrow \bar{q}q \]

is open. Unfortunately, this channel is velocity suppressed since \( \chi \) is a Majorana fermion. The dark matter candidate can annihilate into two gauge bosons or into Higgs bosons through the \( t \)-channel

\[ \chi \chi \rightarrow Z_B Z_B, H_i H_j \]
if it heavy enough.

The direct detection of this dark matter candidate can be through the Higgs portal or the leptophobic gauge boson. The elastic nucleon-dark matter cross section mediated by the gauge boson is velocity suppressed and it is difficult to test the predictions at current experiments. The most optimistic scenario for placing constraints arise from the Higgs portal interaction. In this case, predictions for direct detection of a Majorana fermion has been studied by many authors in the field. See for example Ref. [15]. The detailed analysis of this dark matter candidate will appear in a future publication.

IV. SUMMARY

In this article, we have proposed the simplest extension of the Standard Model where baryon and lepton numbers can be defined as local symmetries spontaneously broken at the low scale. The new fermionic sector of the theory is composed of four new representations—the lepto-baryons—and right-handed neutrinos. The baryon number is broken in three units allowing for a low scale symmetry breaking consistent with experimental bounds on proton lifetime. After the breaking of baryon number the lepto-baryons have the same quantum numbers as the gaugino-Higgsino fields in the minimal supersymmetric Standard Model. Therefore, most of the phenomenological studies on gauginos and Higgsinos can be relevant for this model.

We have investigated the implications for cosmology when the local baryon number is broken near the electroweak scale. In this case, the sphaleron condition is different, changing the relationship between the initial \( B - L \) asymmetry generated by a mechanism such as leptogenesis and the final baryon asymmetry. This model also predicts the existence of a cold dark matter candidate which is a Majorana fermion.

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