Neutrino masses, cosmological inflation and dark matter in a $U(1)_{B-L}$ model with type II seesaw mechanism.

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Abstract

In this work we consider a $U(1)_{B-L}$ gauge model where the spontaneous breaking of the $U(1)_{B-L}$ symmetry engenders the type II seesaw mechanism. This allows that the right-handed neutrinos required by the model be completely dark in relation to the standard interactions. As immediate consequence, these neutrinos turn out to be natural candidates for the dark matter of the universe. We calculate their relic abundance and investigate their detections. Moreover, we show that by triggering the non-minimal coupling of gravity with the scalars, inflation is successfully driven by the neutral component of the scalar triplet.
I. INTRODUCTION

Experimental observation of solar and atmospheric neutrino oscillations surprisingly revealed that neutrinos are very light particles, with mass at sub-eV scale. From the theoretical point of view, seesaw mechanism is the most popular way of generating tiny neutrino masses.

The observation of galaxies rotation curves, cluster collisions and the precise measurements of the thermal anisotropy of the cosmic microwave background suggest the existence of dark matter (DM) permeating our universe. Recent results from PLANCK satellite indicates that 26.7% of the matter of the universe is in the form of non-luminous matter. The most popular DM candidate is a weakly interactive massive particle (WIMP). WIMPs can be any kind of particle since they are neutral, stable (or sufficiently long lived) and have mass in the range from few GeV’s up to some TeV’s.

Cosmological inflation is considered the best theory for explaining homogeneity, flatness and isotropy of the universe as required by hot big bang. Experiments in cosmology, as WMAP7 and PLANCK15, entered in an era of precision which allow us to probe proposal of physics scenario that try to explain the very early universe. Single-field slow-roll models of inflation coupled non-minimally to gravity appear to be an interesting scenario for inflation since it connects inflation to particle physics at low energy scale.

Although the standard model (SM) of particle physics is a very successful theory, its framework does not accommodate any one of the three issues discussed above. In other words, nonzero neutrino mass, dark matter and inflation require extensions of the SM.

In order to accomplish these three issues within a phenomenological viable particle physics model, we evoke the $U(1)_{B-L}$ gauge symmetry. Remember that the minimal $U(1)_{B-L}$ gauge model (B-L model), which is the natural framework of the type I seesaw mechanism, involves one neutral scalar singlet $S$ and three right-handed neutrinos (RHN) $N_i (i = 1, 2, 3)$ in addition to the SM content. Our proposal is modify the minimal B-L model by adding a triplet of scalar ($\Delta$) to its canonical scalar content. As nice results we have that small neutrino masses is achieved through the type II seesaw mechanism which is triggered by the spontaneous breaking of the B-L symmetry, the dark matter content of the universe is provided by the right-handed neutrinos and, by allowing non-minimal coupling of gravity with scalars, we show that the model perform inflation.
This work is organized as follows: in Section II, we describe the main properties of the B-L model. Section III is devoted to cosmological inflation. In Section IV we describe our calculation of the dark matter candidate, and Section V contains our conclusions.

II. THE B-L MODEL WITH SCALAR TRIPLET

A. The seesaw mechanism

Baryon number (B) and lepton number (L) are accidental anomalous symmetries of the SM. However, it is well known that only some specific linear combinations of these symmetries can be free from anomalies\cite{14, 19–21}. Among them, the most developed one is the B-L symmetry\cite{14–17} which is involved in several physics scenarios such as GUT\cite{22}, seesaw mechanism\cite{23–26} and baryogenesis\cite{27}. This symmetry engenders the simplest gauge extension of the SM, namely, the \textit{B-L model} which is based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$. In this work we consider an extension of the B-L model with the scalar sector being composed by a scalar triplet in addition to the doublet and scalar singlet of the minimal version. In this way, the particle content of the model is the SM one augmented by three RHNs, $N_i$, $i = 1, 2, 3$, one scalar singlet, $S$, and one scalar triplet,

$$\Delta \equiv \begin{pmatrix} \frac{\Delta^+}{\sqrt{2}} & \frac{\Delta^{++}}{\sqrt{2}} \\ \Delta^0 & -\frac{\Delta^+}{\sqrt{2}} \end{pmatrix}. \quad (1)$$

As far as we know, this is the first time the triplet $\Delta$ appears composing the scalar sector of the B-L model. Moreover, we impose the model to be invariant by a $Z_2$ discrete symmetry with the RHNs transforming as $N_i \rightarrow -N_i$ while the rest of the particle content of the model transforms trivially by $Z_2$.

With these features, the Yukawa interactions involve the following terms

$$\mathcal{L}_{B-L} \supset Y_\nu \overline{f} C i \sigma^2 \Delta f + \frac{1}{2} Y_N \overline{N} S + h.c., \quad (2)$$

where $f = (\nu \ e)^T_L$. Perceive that both neutrinos gain masses when $\Delta^0$ and $S$ develop a nonzero vacuum expectation value ($v_\Delta$ and $v_S$). This yields the following expressions to the masses of these neutrinos

$$m_\nu = \frac{Y_\nu v_\Delta}{\sqrt{2}}, \quad m_{\nu_R} = \frac{Y_N v_S}{\sqrt{2}}. \quad (3)$$
Observe that there is no mixing mass terms involving $N$ and $\nu_L$ like in the type I seesaw mechanism. The energy scale of the neutrino masses is defined by $v_\Delta$ and $v_S$. Thus, small masses for the standard neutrinos requires a tiny $v_\Delta$. We are going to show that on fixing $v_h$ and $v_S$ we may obtain $v_\Delta$ around eV scale à la type II seesaw mechanism \cite{28–30}. For this we must develop the potential of the model which involves the following terms

$$
V(H, \Delta, S) = \mu_h^2 H^\dagger H + \lambda_h (H^\dagger H)^2 + \mu_s^2 S^\dagger S + \lambda_s (S^\dagger S)^2 \\
+ \mu_\Delta^2 Tr(\Delta^\dagger \Delta) + \lambda_\Delta Tr[(\Delta^\dagger \Delta)^2] + \lambda_\Delta^\prime Tr[(\Delta^\dagger \Delta)^2] \\
+ \lambda_1 S^\dagger S H^\dagger H + \lambda_2 H^\dagger \Delta^\dagger \Delta H + \lambda_3 Tr(\Delta^\dagger \Delta) H^\dagger H \\
+ \lambda_4 S^\dagger S Tr(\Delta^\dagger \Delta) + (k H \sigma^2 \Delta^\dagger HS + h.c.).
$$

where $S$ is the scalar singlet by the standard symmetry and $H = (h^+ h^0)^T$ is the standard Higgs doublet. We assume that all neutral scalars develop vev different from zero. To obtain the set of conditions that guarantee such potential have a global minimum, we must shift the neutral scalar fields in the conventional way

$$
S, h^0, \Delta^0 \rightarrow \frac{1}{\sqrt{2}} (v_{S,h,\Delta} + R_{S,h,\Delta} + i I_{S,h,\Delta}),
$$

and then substitute them in the potential above. Doing this we obtain the following set of minimum condition equations

$$
v_S \left( \mu_s^2 + \frac{\lambda_1}{2} v_h^2 + \frac{\lambda_4}{2} v_\Delta^2 + \lambda_s v_S^2 \right) - \frac{k}{2} v_h^2 v_\Delta = 0,
$$

$$
v_h \left( \mu_h^2 + \frac{\lambda_1}{2} v_S^2 + \frac{\lambda_2}{2} v_\Delta^2 + \frac{\lambda_3}{2} v_\Delta^2 + \lambda_h v_h^2 - k v_\Delta^2 v_S \right) = 0,
$$

$$
v_\Delta \left( \mu_\Delta^2 + \frac{\lambda_2}{2} v_h^2 + \frac{\lambda_3}{2} v_h^2 + \frac{\lambda_4}{2} v_\Delta^2 + (\lambda_\Delta + \lambda_\Delta^\prime) v_\Delta^2 \right) - \frac{k}{2} v_h^2 v_\Delta = 0.
$$

Remember that $v_\Delta$ modifies softly the $\rho$-parameter in the following way: $\rho = \frac{1 + 2 v_\Delta^2}{1 + \frac{4 v_\Delta^2}{\rho_s^2}}$. The electroweak precision data constraints require the value $\rho = 1.00037 \pm 0.00023$\cite{1}. This implies the following upper bound $v_\Delta < 4\text{GeV}$. Consequently, $v_S$ passes to contribute dominantly to the mass of the new neutral gauge boson $Z'$ associated to the B-L symmetry which also has a constraint due the LEP experiment\cite{31}

$$
\frac{m_{Z'}}{g_{B-L}} \gtrsim 6.9\text{TeV}.
$$
Perceive that the third relation in Eq. (6) provides
\[ v_\Delta \approx \frac{k v_h^2 v_S}{2 \mu_\Delta^2}. \] (8)

The role of the type II seesaw mechanism is to provide tiny vevs. In the canonical case, where \( v_\Delta = \frac{v_h^2}{\mu} \), tiny \( v_\Delta \) is a consequence of the explicit violation of the lepton number which must happens at GUT scale (\( \mu = 10^{14} \) GeV). Observe that in our case \( v_\Delta \) get suppressed by the quadratic term \( \mu_\Delta^2 \). This allows we have a seesaw mechanism occuring in an intermediate energy scale, as we see below.

Remember that \( v_h \) is the standard vev whose value is 247 GeV while \( v_S \) define the mass of the neutral gauge bosons \( Z' \) and its value must lie around few TeVs. Here we take \( v_S \sim 10 \) TeV. This provides \( v_h^2 v_S \sim 10^9 \text{ GeV}^3 \). Consequently, \( v_\Delta \sim \text{eV} \) requires \( \mu_\Delta \sim 10^9 \) GeV. As conclusion, we have that type II seesaw mechanism engendered by the spontaneous breaking of the lepton number is associated to a new physics in the form of scalar triplet with mass around \( 10^9 \) GeV. Such energy scale is too high to be probed at the LHC. However such regime of energy may give sizable contributions in flavor physics and then be probed through rare lepton decays. This point will be discussed elsewhere.

B. Spectrum of scalars

Before go further, it makes necessary discuss briefly the scalar sector of the model. Let us first focus on the CP-even sector. In the basis \( (R_S, R_h, R_\Delta) \) we have the following mass matrix,
\[
M_{R_\Delta}^2 = \begin{pmatrix}
\frac{k v_\Delta v_h^2}{2 v_s} + 2 \lambda_S v_s^2 & -kv_h v_\Delta + \lambda_1 v_s v_h & -\frac{k v_h^2}{2} + \lambda_4 v_s v_\Delta \\
-kv_h v_\Delta + \lambda_1 v_s v_h & 2\lambda_h v_h^2 & -kv_s v_h + (\lambda_2 + \lambda_3) v_h v_\Delta \\
-\frac{k v_h^2}{2} + \lambda_4 v_s v_\Delta & -kv_s v_h + (\lambda_2 + \lambda_3) v_h v_\Delta & \frac{k v_h^2 v_\Delta^2}{v_S^2} + 2(\lambda_\Delta + \lambda_\Delta') v_\Delta^2
\end{pmatrix}.
\] (9)

Note that for values of the vevs indicated above, the scalar \( R_\Delta \) get very heavy, with \( m_\Delta^2 \sim \frac{k v_\Delta^2 v_h^2}{2 v_\Delta^2} \), which implies that it decouples from the other ones. The other two quadratic masses are
\[
m_h^2 \approx 2\lambda_h v_h^2 - \frac{1}{2} \lambda_1^2 v_h^2,
\]
\[
m_H^2 \approx 2\lambda_S v_S^2 + \frac{1}{2} \lambda_1^2 v_h^2,
\] (10)
where \( m_h \) stands for the standard Higgs boson with the allowed parameter space showed in FIG. 1

FIG. 1: Possible values of the quartic couplings that yield 125 GeV Higgs mass.

The respective eigenvectors are

\[
\begin{align*}
h &\simeq R_h - \frac{\lambda_1}{2\lambda_S} v_h R_S, \\
H &\simeq R_S + \frac{\lambda_1}{2\lambda_S} v_h R_h. \tag{11}
\end{align*}
\]

For the CP-odd scalars, we have the mass matrix in the basis \((I_S, I_h, I_\Delta)\),

\[
M_I^2 = \begin{pmatrix}
\frac{k v_H v_s^2}{2 v_s} & k v_h v_\Delta & -\frac{k}{2} v_h^2 \\
k v_h v_\Delta & 2 k v_s v_\Delta & -k v_s v_h \\
-\frac{k}{2} v_h^2 & -k v_s v_h & k \frac{v_s v_h^2}{v_\Delta^2}
\end{pmatrix}.
\tag{12}
\]

The mass matrix in Eq. (12) can be diagonalized providing one massive state \( A^0 \) with mass,

\[
m_A^2 = \frac{k}{2} \left( \frac{v_\Delta v_h^2}{v_s} + \frac{v_s v_h^2}{v_\Delta} + 4 v_s v_\Delta \right), \tag{13}
\]

and two goldstone bosons \( G^1 \) and \( G^2 \), absorbed as the longitudinal components of the \( Z \) and \( Z' \) gauge bosons. The eigenvectors for the cp-odd scalars are,

\[
\begin{align*}
G^1 &\simeq I_S + \frac{v_\Delta}{v_s} I_\Delta, \\
G^2 &\simeq I_h + \frac{v_h}{2 v_s} I_S, \\
A^0 &\simeq I_\Delta - \frac{2 v_\Delta}{v_h} I_h. \tag{14}
\end{align*}
\]
The charged scalars, given in the basis \((h^+, \Delta^+)\), have the mass matrix

\[
M_+^2 = \begin{pmatrix}
kv_s v_\Delta - \frac{\lambda_2}{2} v_\Delta^2 & \frac{\lambda_3}{\sqrt{2}} v_h v_\Delta - \frac{k}{\sqrt{2}} v_s v_h \\
\frac{\lambda_3}{\sqrt{2}} v_h v_\Delta - \frac{k}{\sqrt{2}} v_s v_h & \frac{k v_s^2 v_\Delta^2}{2} - \frac{\lambda_2}{2} v_h^2
\end{pmatrix}.
\] (15)

Again, diagonalizing this matrix gives us two goldstone bosons \(G^\pm\), responsible by the longitudinal parts of the \(W^\pm\) standard gauge bosons. The other two degrees of freedom give us a massive states \(H^\pm\) with mass

\[
m_{H^\pm}^2 = \left(\frac{v_\Delta}{2} + \frac{v_h^2}{4 v_\Delta}\right) \left(2 k v_s - \lambda_2 v_\Delta\right).
\] (16)

The respective eigenvectors are,

\[
G^\pm \simeq h^\pm + \frac{\sqrt{2} v_\Delta}{v_h} \Delta^\pm,
\]

\[
H^\pm \simeq \Delta^\pm - \frac{\sqrt{2} v_\Delta}{v_h} h^\pm.
\] (17)

Finally the mass of the doubly charged scalars \(\Delta^{\pm\pm}\) are expressed as,

\[
m_{\Delta^{\pm\pm}}^2 = \frac{kv_s v_h^2 v_\Delta - \lambda_2 v_h^2 v_\Delta^2 - 2 \lambda_\Delta v_\Delta^4}{2 v_\Delta^2}.
\] (18)

Once symmetries are broken and the gauge bosons absorb the goldstone bosons as longitudinal component, we have that the standard charged bosons get a contribution from the triplet vev, \(m_{w}^2 = g^2 \left(v_h^2 + 2 v_\Delta^2\right)\), while the neutral gauge bosons get mixed with \(Z'\) in the following way

\[
M_g^2 = \begin{pmatrix}
\frac{g^2 + g'^2}{4} (v_h^2 + 4 v_\Delta^2) & -g_{B-L} \sqrt{g^2 + g'^2} v_\Delta^2 \\
-g_{B-L} \sqrt{g^2 + g'^2} v_\Delta^2 & g_{B-L}^2 (2 v_S + v_\Delta^2)
\end{pmatrix}.
\] (19)

Recording the vev hierarchy discussed here \((v_S > v_h \gg v_\Delta)\), the mixing between the gauge bosons are very small, therefore, they decouple and we have the masses,

\[
M_Z^2 \approx \frac{(g^2 + g'^2)(v_h^2 + 4 v_\Delta^2)}{4}, \quad M_{Z'}^2 \approx 2 g_{B-L}^2 (v_S^2 + \frac{v_\Delta^2}{2}).
\] (20)

Observe that we have a B-L model which has as new ingredients scalars in the triplet and singlet forms and neutrinos with right-handed chiralities. Let us resume the role played by these new ingredients. The singlet \(S\) is responsible by the spontaneous breaking of the B-L symmetry and then define the mass of the \(Z'\). The triplet \(\Delta\) is responsible by the type II seesaw mechanism that generate small masses for the standard neutrinos. The right-handed
neutrinos is responsible by the cancellation of anomalies. It would be interesting to find new roles for these ingredients.

We argue here, and check below, that the right-handed neutrinos may be the dark matter component of the universe since the $Z_2$ symmetry protect them of decaying in lighter particles. In the last section we assume that the lightest right-handed neutrinos is the dark matter of the universe and then calculate its abundance and investigate the ways of detecting it.

We also argue here that once $\Delta^0$ has mass around $10^9$ GeV, it could be possible that it would come to be the inflaton and then drives inflation. We show in the next section that this is possible when we assume non-minimal coupling of $\Delta$ with gravity.

III. COSMOLOGICAL INFLATION

As argued above, $\Delta^0$ may be the inflaton and then drives inflation. For previous works of inflation conducted by $\Delta^0$, see [32–34]. However, for successful inflation we have to evoke non-minimal coupling of the inflaton with gravity. For sake of simplicity, we assume that $\Delta^0$ provides the dominant coupling. Thus, the relevant terms in the potential pass to be the quartic one $V(\Delta^0) = \frac{\lambda + \lambda'}{4} \Delta^0 \Delta^{04}$. In this case we have that the lagrangian of the model must involves the following terms in the Jordan frame

$$L \supset \frac{1}{2} (\partial_{\mu} \Delta^0) (\partial^{\mu} \Delta^0) - \frac{M_P^2 R}{2} - \frac{1}{2} \xi \Delta^0 R - V(\Delta^0).$$

(21)

In order to calculate the parameters related to inflation we must recover the canonical Einstein-Hilbert gravity. This process is called conformal transformation and can be understood by two steps. First we re-scale the metric $\tilde{g}_{\alpha \beta} = \Omega^2 g_{\alpha \beta}$. In doing so the coupling to gravity vanish but the inflaton gain a non-canonical kinetic term. The process is finished by transforming the inflaton field to a form with canonical kinetic energy. Such transformation involves the relations [35, 36]

$$\tilde{g}_{\mu \nu} = \Omega^2 g_{\mu \nu} \quad \text{where} \quad \Omega^2 = 1 + \frac{\xi \Delta^0}{M_P^2}$$

$$\frac{d\chi}{d\Delta^0} = \sqrt{\frac{\Omega^2 + 6\xi^2 \Delta^0^2 / M_P^2}{\Omega^4}}.$$

(22)
The lagrangian in Einstein frame is given by,
\[
\mathcal{L} \supset -\frac{M_P^2 \tilde{R}}{2} + \frac{1}{2} (\partial_\mu \chi)^\dagger (\partial^\mu \chi) - U(\chi),
\] (23)
where \(U(\chi) = \frac{1}{4\xi^4} V(\Delta[\chi])\). There is some discussion about which frame is the physical one \[37\], however both frames agrees in the regime of low energy.

As we will see ahead, similar to the canonical Higgs inflation, our model demands a sizable non-minimal coupling (\(\xi > 1\)), even though \(\xi\) is much smaller than the canonical case. For that reason, the potential in Einstein frame can be approximated to the canonical form,
\[
U(\chi) \approx \frac{\lambda M_P^4}{4\xi^2} \left( 1 - \exp \left( -\frac{2\chi}{\sqrt{6}M_P} \right) \right)^2,
\] (24)
where we have made use to the relation among the fields,
\[
\Delta_0 \approx \left( \frac{M_P^2}{\xi} \exp \left( -\frac{2\chi}{\sqrt{6}M_P} \right) - \frac{M_P^2}{\xi} \right)^{\frac{1}{2}}.
\] (25)

Inflation occurs whenever the field \(\chi\), or equivalently \(\Delta_0\), rolls slowly in direction to the minimum of the potential. The slow roll parameters can be written as,
\[
\epsilon = \frac{M_P^2}{2} \left( \frac{V'}{V} \right)^2, \quad \eta = M_P^2 \left( \frac{V''}{V} - \frac{V'V''}{V^2} \right),
\] (26)
where \(^'\) indicate derivative with respect to \(\Delta_0\). Those forms of slow roll parameters permit us to calculate the spectral indices in an exact way using Eq. \[22\]. Inflation starts when \(\epsilon, \eta \ll 1\) and stop when \(\epsilon, \eta = 1\). In the slow roll regime we can write the spectral index and the scalar-to-tensor ratio as \[38\],
\[
n_S = 1 - 6\epsilon + 2\eta, \quad r = 16\epsilon.
\] (27)

The spectral index gives the dependence of the perturbation with the scale, while the scalar-to-tensor ratio gives the ratio between the scalar and tensor perturbation. Planck2015 have measured \(n_S = 0.9644 \pm 0.0049\) and gave the bound \(r < 0.149\) for a pivot scale equivalent to \(k = 0.05\) Mpc\(^{-1}\) \[11\]. Any inflationary model that intend to be realistic must recover these values.

The amount of “visible” inflation is quantified by the number of e-folds,
\[
N = -\frac{1}{M_P^2} \int_{\Delta_0}^{(\Delta_0)'_i} \frac{V(\Delta_0)}{V'(\Delta_0)} \left( \frac{d\chi}{d\Delta_0} \right)^2 d\Delta_0,
\] (28)
where again we use Eq. (22) to evaluate $N$. The number of e-folds is not a free parameter and it depends on the form of the potential. Its major source of uncertainty comes from reheating phase, considering that the equation of state of this period is model dependent. For a $\phi^4$ theory the universe is radiative dominated during the reheat, so it is plausible to affirm that in our model the number of e-folds should be around 60 \cite{39,40}. Using $N = 60$ we can solve Eq. (28) for $(\Delta^0)_i$. Substituting it in Eqs. (27) we obtain the spectral index and scalar-to-tensor ratio. Before going to this point, we consider the effect of radiative corrections for the scalar potential. This makes necessary because the amplitude of curvature perturbation $A_s = V^{24\pi^2 M^{-4}_{\text{P}} \Omega^4_{\text{4}} \epsilon}$, measured by Planck to be $A_s = 2.21 \times 10^{-9}$ for $k = 0.05$ Mpc$^{-1}$, translates in strong constrain on the couplings of inflationary potential ($\lambda_\Delta$, $\lambda'_\Delta$ and $\xi$).

As we are dealing with an inflaton that is triplet under $SU(2)_L$, such corrections will involve the standard gauge couplings $g$, $g'$ and $g_{B-L}$. Here we consider one-loop radiative correction in Jordan frame, as done in \cite{41–43}. The complete potential involving the relevant terms for inflation is given by

$$V = \frac{1}{4} \lambda_\Delta + \frac{1}{4} \lambda'_\Delta + \sum_i \frac{3(g^4 + g'^4 + g_{B-L}^4)}{32\pi^2} - 4 Y_\nu \epsilon^4 + \sum_i \frac{\lambda_i^2}{M_{\text{P}}} \ln \frac{\Delta^0}{M_{\text{P}}} \Delta^0,$$

(29)

where $M_{\text{P}}$ is chosen for renormalization scale and $i$ runs for the scalar contributions ($\lambda_2$, $\lambda_3$, $\lambda_4$, $c_\Delta \lambda_\Delta$ and $c_\Delta \lambda'_\Delta$). Note the suppression factor in the loop term coming from the inflaton ($c_\Delta$). It arises from the suppression in the propagator of inflaton field \cite{43}. The main contributions to the radiative corrections involve the standard couplings $g$ and $g'$. At grand unification scale the standard gauge couplings are evaluated at $g^2 \approx g'^2 \approx 0.3$ \cite{42}. Here we assume $g_{B-L} < g, g'$ at any energy scale. Furthermore, for $v_\Delta$ at eV scale, the Yukawa coupling $Y_\nu$ must develop value in the range $10^{-2} - 10^{-3}$ in order to generate neutrino mass capable of explaining solar and atmospheric neutrino oscillation. Finally, it is natural to assume $\lambda_i << 1$. In view of this, the dominant terms in the potential are

$$V = \frac{1}{4} \lambda_\Delta + \frac{1}{4} \lambda'_\Delta + \sum_i \frac{3(g^4 + g'^4)}{32\pi^2} \ln \frac{\Delta^0}{M_{\text{P}}} \Delta^0,$$

(30)

Remember that, in the case of the standard Higgs inflation\cite{13}, as the standard quartic coupling , $\lambda_i$ is already fixed at 0.6, then the tree level contribution in the effective potential of the inflaton get dominant over the radiative corrections, and then cosmological constraints imply $\xi \sim 10^4$ which yield loss of unitarity at $\Lambda_U = \frac{M_{\text{P}}}{\xi} \sim 10^{14}$ GeV \cite{44,45}.

Differently from the standard case, $\lambda_\Delta$ and $\lambda'_\Delta$ are free parameters. In this case it is very likely that the magnitude of the gauge coupling at high scales overcomes the tree level
contributions in the potential. For a radiative dominance case, the inflaton potential at inflation time reads,

$$V = \frac{3(g^4 + g'^4)}{32\pi^2} \ln \frac{\Delta^0}{M_P} \Delta^0.$$  \hspace{1cm} (31)

We can use the value of $A_s$ observed by Planck satellite to constrain the non-minimal coupling $\xi$. For $A_s = 2.21 \times 10^{-9}$ we must have $\xi \sim 1880$.

For these values of the couplings, inflation happens for $\Delta^0$ in the range: $3.20 \times 10^{18} < \Delta^0 \lesssim 2.12 \times 10^{20}$ GeV. Although this energy scale is much higher than the Planck one, the potential energy is fixed at $V^{\frac{1}{4}} \sim 1.64 \times 10^{16}$ GeV. Note that in our case the model face problem with the loss of unitarity in $\Lambda_U \sim 1.28 \times 10^{15}$, which is one order of difference in relation to $V^{\frac{1}{4}}$. This is a gain in comparison to the standard Higgs inflation case[13].

Finally we calculate the spectral index and tensor to scalar ratio. For $\xi = 1880$, we obtain

$$r = 0.066 \quad \text{and} \quad n_s = 0.975.$$  \hspace{1cm} (32)

This set of values is in agreement with the 68% CL contour of the most stringent data set of Planck2015 ($Planck \; TT + TE + EE + LowP$) [11]. In particular, our model gives a sizable tensor to scalar ratio. This suggest observable primordial gravity waves.

We end this section advocating the necessity of new physics to reconcile the loss of unitarity[45]. As possible solutions, we refer the proposals in refs. [46] [47]

IV. DARK MATTER

In our model the three RHNs transform non trivially by the $Z_2$ symmetry. Consequently the model does not perform the type I seesaw mechanism. As we saw above, neutrino mass is achieved through an adapted type II seesaw mechanism. In view of this the role played by the RHNs is to cancel gauge anomalies and, as consequence of the $Z_2$ symmetry, provide the dark matter content of the universe in the form of WIMP. However, differently from the minimal B-L model[18], here the three RHNs may be stable particles since we chose $Y_N$ diagonal in Eq. (2). In other words, the model may accommodate multiple DM candidates.

Although the three RHNs are potentially DM candidate, for simplicity reasons we just consider that the lightest one, which we call $N$, is sufficient to provide the correct relic abundance of DM of the universe in the form of WIMP. This means that $N$ was in thermal equilibrium with the SM particles in the early universe. Then, as far as the universe expands
and cools the thermal equilibrium is lost causing the freeze out of the abundance of \(N\). This happens when \(N\) annihilation rate, whose main contributions are displayed in FIG. 2 becomes roughly smaller than the expansion rate of the universe. In this case the relic abundance of \(N\) is obtained by evaluating the Boltzmann equation for the number density \(n_N\),

\[
\frac{dn_N}{dt} + 3H n_N = -\langle \sigma v \rangle (n_N^2 - n_{\text{EQ}}^2),
\]

where

\[
H^2 \equiv \left( \frac{\dot{a}}{a} \right)^2 = \frac{8\pi}{3M_P^2} \rho,
\]

with \(n_{\text{EQ}}\) and \(a(t)\) being the equilibrium number density and the scale factor in a situation where the radiation dominates the universe with the energy density \(\rho = \rho_{\text{rad}}\), i.e., the thermal equilibrium epoch. \(\langle \sigma v \rangle\) is the thermal averaged product of the annihilation cross section by the relative velocity. As usually adopted, we present our results in the form of \(\Omega_N\), which is the ratio between the energy density of \(N\) and the critical density of the universe.

![FIG. 2](image)

**FIG. 2:** The main contributions to the DM relic Abundance. The SM contributions stands for fermions, Higgs and vector bosons.

We proceed as follow. We analyze numerically the Boltzmann equation by using the micrOMEGAs software package (v.2.4.5)\[48, 49\]. For this we implemented the model in Sarah (v.4.10.2)\[50–53\] in combination with SPheno (v.3.3.8)\[54, 55\] package, which solves all mass matrix numerically. Our results for the relic abundance of \(N\) are displayed in FIG. 3. The thick line in those plots correspond to the correct abundance. Perceive that, although the Yukawa coupling \(Y_N\) is an arbitrary free parameter, which translates in \(N\) developing any mass value, however what determine if the model provides the correct abundance of \(N\) is the resonance of \(Z'\) and \(H\). This is showed in the top left plot of FIG. 3. In it the first resonance correspond to a \(Z'\) with mass around 3 TeVs. The second resonance correspond to \(H\) with mass in the range from 6 TeV up to 7 TeV. In the top right plot we show the density
of DM varying with its mass but now including the region excluded by LEP constraint given in Eq. (7). On the bottom left of FIG. 3 we made a zoom in the resonance of $H$. Perceive that we included the LEP constraint. For completeness reasons, on the bottom right we show the dependence of $M'_Z$ with $g_{B-L}$ including the LEP exclusion region, too. In the last three plots we show two benchmark point localized exactly in the line that gives the correct abundance. They are represented in red square and orange star points and their values are displayed in the table. Observe that LEP constraint imposes $N$ reasonably massive with mass around few TeVs. To complement these plots, we show the ones in FIG. 5 which relates the resonance of $H$ (all points in color) and the mass of the DM. The mixing parameter $\sin \theta$ is given in Eq. (11). Observe that LEP exclusion is a very stringent constraint imposing $H$ with mass above 5800 GeV and requiring DM with mass above 2800 GeV. All those points in colors give the correct abundance but only those in black recover the standard Higgs with mass of 125 GeV. The benchmark points in the red square and orange star are given in the table I. In summary, for the set of values for the parameters choose here $N$ with mass around 3 TeVs is a viable DM candidate once provides the correct relic abundance required by the experiments [4]. That is not all. A viable DM candidate must obey the current direct detection constraints.

| $M_{DM}$ (GeV) | $Y_{N1}$ | $M_{Z'}$ (GeV) | $g_{B-L}$ | $M_H$ (GeV) | $\Omega h^2$ | $v_S$ | $\sigma_{DMq}$ | $q$ |
|----------------|----------|----------------|------------|-------------|-------------|------|----------------|-----|
| 3050           | 0.291    | 3840           | 0.518      | 6279        | 0.116       | 7400 | 5.4 10^{-11}  | 🂽  |
| 3190           | 0.294    | 3904           | 0.509      | 6470        | 0.122       | 7658 | 5.4304e-11    | ⚫  |

TABLE I: Benchmark points for parameters values added on plots.

In addition to the relic density of the DM candidate, which involves gravitational effects, only, we need to detect it directly in order to reveal its nature. Here we restrict our studies to direct detection, only. In attempting to detect directly signal of DM in the form of WIMPs, many underground experiment, using different sort of targets, were built. Unfortunately no signal has been detected, yet. The results of such experiments translate in upper bounds into the WIMP-Nucleon scattering cross section. In view of this any DM candidate in the form of WIMPs must be subjected to the current direct detection constraints. Direct detection theory and experiment is a very well developed topic in particle physics. For a review of the theoretical predictions for the direct detection of WIMPs in particle physics models,
FIG. 3: Plots relating DM relic Abundance, Yukawa coupling of the RHN and the dark matter candidate mass. The thick horizontal line correspond to the correct relic abundance[1].

FIG. 4: The WIMP-quark scattering diagram for direct detection.

see [56, 58]. For a review of the experiments, see [59]. In our case direct detection requires interactions among $N$ and quarks. This is achieved by exchange of $h$ and $H$ via t-channel as displayed in FIG. 4. Perceive that $Z'$ t-channel gives null contribution because $N$ is a Majorana particle. In practical terms we need to obtain the WIMP-quark scattering cross section which is given in [60]. Observe that the scattering cross section is parametrized by four free parameters, namely $M_N, M_H, v_S$ and the mixing angle $\theta$ given in eq. (11). However this cross section depends indirectly on other parameters. For example, for $g_{B-L}$
FIG. 5: Physical parameter space of $M_{H_2} \times M_{DM}$. In colors we show the points that provide the correct relic abundance in the resonant production of the singlet scalar according with the diagrams in Fig. 4. The black points recover a Higgs with mass of 125 GeV.

in the range $0.1 - 0.55$ (as explained in Section III) the LEP constraint implies $v_S > 7$ TeV and $M_{Z'}$ around 3 TeV. Considering this, and using the micrOMEGAs software package [48, 49] in our calculations, we present our results for the WIMP-Nucleon cross section in FIG. 6. All color points conduct to the right abundance and are in accordance with the Lux (2017) exclusion bound. Those points in pink are excluded by LEP constraint. However only those points in black recover a Higgs with mass of 125 GeV. Observe that the black points may be probed by future XenonNnT and DarkSide direct detection experiments. This turn our model a phenomenological viable DM model.

V. CONCLUSIONS

In this work we implemented the type II seesaw mechanism for generation of small neutrino masses in a B-L gauge model. For achieve this we added one triplet of scalars ($\Delta$) to the canonical scalar content of the B-L model. The mechanism is triggered when the B-L symmetry is broken spontaneously. We showed that neutrino masses at eV scale require that $\Delta$ belongs to an energy scale around $10^9$ GeV. This characterizes a seesaw mechanism
FIG. 6: The spin-independent (SI) WIMP-nucleon cross sections constraints. The lines corresponds to experimental upper limits bounds on direct detection. For LUX [61] (black line with blue fill area), Xenon1T [62] (blue line, prospect), XenoNnT [63] (prospect, green line), Dark Side Prospect [64] (red and dark red lines for different exposure time) and the neutrino coherent scattering, atmospheric neutrinos and diffuse supernova neutrinos [65] (orange dashed line with filled area).

at intermediate energy scale and can be probed through rare lepton decays.

One interesting advantage of this model is that we can evoke a \( Z_2 \) discrete symmetry and leave the right-handed neutrinos completely dark in relation to the standard model interactions. In this case these neutrinos turn out to be the natural candidate for the dark matter of universe in the form of WIMP. In this case we showed that the correct abundance of dark matter is obtained thanks to the resonant production of \( Z' \) and of the heavy Higgs \( H \). Although our scenario is in accordance with LUX exclusion bound, however prospect direct detection experiments will be able to probe it.

In what concern inflation, we showed that by allowing non-minimal coupling of the neutral component of the scalar triplet with gravity we succeeded in accommodating the inflation cosmological parameters in a scenario where the loss of unitarity occurs one order of magnitude below the energy density during inflation[47][45]. That is a great gain in what concern Higgs inflation.
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