Dark Matter and Neutrino Mass Models: Phenomenology of the Scalar Sector

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Abstract. Many standard model (SM) extensions that are motivated by the problems of dark matter (DM) and neutrino mass and oscillations include extra scalars. The existence of extra singlets may modify the strength of the Higgs decay \( h \to \gamma \gamma \), the Higgs triple coupling and the electroweak phase transition (EWPT). During the EWPT dynamics gravitational waves (GWs) could be generated, and could be detected by future experiments such as DECIGO and LISA. In this work, we investigate these aspects.

1. Introduction
Despite its success, the Standard Model (SM) of particle physics requires extensions and/or modifications to address many unanswered questions such as dark matter (DM), neutrino mass and oscillations data, baryon asymmetry at the Universe, gauge coupling unification .. etc. Then, the need to go beyond the SM is mandatory in order to address the above mentioned problems as well as others. The simplest mechanism to generate small neutrino masses is the so-called seesaw (type-I) mechanism \([7]\), where 3 massive right-handed neutrinos are coupled to the left handed neutrinos. In 2002, Krauss, Nasri and Trodden (KNT) \([8]\) had extend the SM by two singlet scalars and a Majorana fermion. In this model, a global \( Z_2 \) symmetry is imposed as \( \{ S^+_2, N \} \to \{-S^+_2, -N\} \), where DM gets stable and neutrino mass is generated at three-loop\(^1\). This model was promoted by replacing the singlet fields \( S^+_2 \) and \( N \) by \( SU(2)_L \) higher representation such as triplet fields \([9]\), quintuplet fields \([10]\) and septplet fields \([11]\). In addition this model is generalized within a scale invariant framework \([12]\).

In \([13]\), we proposed a simple SM extension where neutrino mass is generated via a radiative inverse seesaw scheme. The model involved three singlet chiral fermions and one singlet complex scalar obeying a global \( Z_4 \) symmetry, that is softly broken into invoked to a \( Z_2 \) symmetry.

\(^1\) For recent studies of the KNT model see Refs. \([7, 14, 15, 16]\).
that makes DM stable. The DM relic density is naturally obtained and the spin-independent scattering cross section of the DM off nucleus is consistent with the experimental limit reported by LUX [?], and yet within the reach of future DM direct detection searches. In a recent work [?], we investigated the case where the DM candidate is the lightest Majorana in the scotogenic model. We found that the phenomenology is significantly different than the scotogenic model with scalar DM. In this work, we investigate common phenomenological aspects such as the electroweak phase transition (EWPT) strength, which is required for a successful implementation of electroweak baryogenesis [?]. In addition, we discuss possible gravitational waves (GWs) that could be generated within the EWPT dynamics by considering the SM extension by a scalar with mass $m_S$, multiplicity $N$, and coupling $\omega$ to the Higgs doublet [?]. Many collider aspects of the charged scalars in these models are discussed at the LHC [?], and leptonic colliders [?].

In section II, we show briefly the three models and discuss some constraints. The EWPT dynamics is described in section III, and the corresponding GWs spectrum properties in section IV. We discuss our results in section V and give our conclusion in section VI.

2. $\nu$DM: Models & Constraints

Here, we present some examples where neutrino mass is generated at loop-level and DM candidate are defined as Majorana fermions.

$KNT \& KNT$-like models

This class of models is based on extending the SM by a singlet charged scalar $S_1^+$, a scalar multiplet $T$, and three generations of fermionic multiplets $E_i$, while retaining the global $Z_2$ symmetry $\{T, E_i\} \rightarrow \{-T, -E_i\}$, and the lightest neutral fermion $E_1^0$ plays the role of a stable DM candidate. The general Lagrangian reads as

$$\mathcal{L} = \mathcal{L}_{SM} + \left\{ f_{\alpha \beta} L_\alpha^T C \tau_2 L_\beta S_1^+ \right\} + g_{i \alpha E_i T} \ell_\alpha R - \frac{1}{2} E_i^c M_{ij} E_j + h.c \right\} - V,$$

where $L_\alpha$ is the left-handed lepton doublet, $f_{\alpha \beta}$ are Yukawa couplings which are antisymmetric in the generation indices $\alpha$ and $\beta$, $M_{ij}$ are the fermionic mass matrix elements, $C$ is the charge conjugation matrix, and $V(\Phi, S_1, T)$ is the tree-level scalar potential. Here $\Phi$ denotes the SM Higgs doublet.

Using interactions in (??) together with the scalar interaction $V \supset \lambda_\alpha S_1^+ S_1^- T^\dagger T$, the neutrino mass matrix elements can arise from the three-loop diagram in Fig. ??, that are given by [?]

$$M_\nu_{\alpha \beta} = \frac{(2n + 1) \lambda_\alpha m_\ell_i m_\ell_k}{(4\pi^2)^2 M_T} f_{\alpha i} f_{\beta k} g_{ij} g_{kj} F \left( \frac{M_{E_i}^2}{M_T^2}, \frac{M_{S_1}^2}{M_T^2} \right),$$

where $\rho, \kappa(= e, \mu, \tau)$ are the charged leptons flavor indices, $i = 1, 2, 3$ denotes the three $E_i$ multiplets, and the function $F$ is a loop integral which is $O(1)$ [?]. Here $n = 0$ corresponds to the KNT model, while $n = 1, 2, 3$ gives generalizations where $E_i$ and $T$ are $SU(2)_L$ triplets, quintuplets and septuplets, respectively (i.e. $T$ and $E_i$ are both assigned to the $(2n + 1)$ representation under $SU(2)_L$ and carry two units of hypercharge).

The Lagrangian (??) induces flavor violating processes, such as $\ell_\alpha \rightarrow \gamma \ell_\beta$ and $\ell_\alpha \rightarrow \ell_\beta \ell_\beta \ell_\beta$ for $m_{\ell_\alpha} > m_{\ell_\beta}$, and an extra contribution to the muon anomalous magnetic moment. Both are generated at one loop via the exchange of the charged scalar $S_1^+$, and the members of the multiplets $T$ and $E_i$. The considered parameter space of the model must respect the existing experimental bounds [?], and match the neutrino oscillation parameters [?].

2 Except for the septuplet case where the global symmetry $Z_2$ is accidental [?].
Figure 1. The three-loop diagram that generates the neutrino mass.

Figure 2. Radiative inverse seesaw with DM.

Dark Radiative Inverse Seesaw
Here, the SM is extended by three generations of chiral fermion pairs $N_R$ and $N_L$, one other chiral fermion $\chi_R \equiv \chi$ and a complex scalar $S$. The fields have the charges under a global $Z_4$ symmetry as

$$
\begin{array}{|c|c|c|c|c|}
\hline
 & L, \ell_R, N_R, N_L & \chi & H & S \\
\hline
Z_4 & i & -1 & +1 & i \\
\hline
\end{array}
$$

This allows to write the Lagrangian

$$
-\mathcal{L} \supset y_\nu \bar{\nu}_R \tilde{T} H N_R + M_N N_R N_L + y_N S \bar{\chi} N_L + \frac{m_\chi}{2} \chi^T C^{-1} \chi + \text{h.c.},
$$

where $\tilde{T} \equiv i \sigma_2 H^*$ and $C$ is the charge conjugation operator. The scalar potential is written as:

$$
V = -\mu_H^2 H^\dagger H + \frac{1}{2} \lambda_H (H^\dagger H)^2 + \mu_S^2 S^* S + \frac{\mu_\chi^2}{2} (S^2 + \text{h.c.}) + \lambda_S (S^* S)^2 + \lambda_H S^\dagger H S^* S.
$$

The neutrino mass can be generated a la inverse seesaw mechanism as shown in Fig. 2. The neutrino mass matrix is given by

$$
m_\nu \simeq m_D^\top \frac{1}{M} \epsilon_L \frac{1}{M} m_D,
$$

where $m_D = y_\nu \langle H \rangle$ is the usual Dirac neutrino mass, and the matrix $\epsilon_L$ is generated radiatively (Fig. 1-right), which is proportional to the soft breaking mass $\mu_\nu^2$ [1]. Here, the chiral fermion $\chi$ plays the role of a DM candidate.

3 For simplicity we add the iso-singlet pairs sequentially, though two pairs would suffice to account for the neutrino oscillations data.
Scotogenic Model with Majorana DM

The SM is extended by three singlet Majorana fermions \( N_i \sim (1, 1, 0) \), and an inert doublet \( \Phi^T = \begin{pmatrix} H^+ & \frac{1}{\sqrt{2}}(H^0 + iA^0) \end{pmatrix} \), both odd under the discrete \( Z_2 \) symmetry, that ensures the DM stability. The neutrino mass can be generated at one loop due to the Lagrangian

\[
\mathcal{L} \supset h_{ij} \bar{L}_i \epsilon \Phi N_j + \frac{1}{2} M_i N_i^C N_i + h.c.,
\]

where \( \bar{L}_i \) is the left-handed lepton doublet and \( \epsilon = i\sigma_2 \) is an antisymmetric tensor. The scalar potential can be written as

\[
V = -\mu_1^2 |H|^2 + \mu_2^2 |\Phi|^2 + \frac{\lambda_1}{6} |H|^4 + \frac{\lambda_2}{6} |\Phi|^4 + \lambda_3 |H|^2 |\Phi|^2 + \frac{\lambda_4}{2} |H^\dagger \Phi|^2 + \frac{\lambda_5}{4} [(H^\dagger \Phi)^2 + h.c.].
\]

In this setup, we have two CP-even scalars \((h, H^0)\), one CP-odd scalar \(A^0\) and a pair of charged scalars \(H^\pm\). Their tree-level masses are given by:

\[
m_{H^\pm}^2 = \mu_2^2 + \frac{1}{2} \lambda_3 v^2, \quad m_{H^0, A^0}^2 = m_{H^\pm}^2 + \frac{1}{4} (\lambda_4 \pm \lambda_5) v^2.
\]

The neutrino mass can be generated a la scotogenic as shown in Fig. 3.

![Feynmann diagram](image)

**Figure 3.** Feynmann diagram responsible for the neutrino mass.

The one loop neutrino mass matrix elements are given by [2]

\[
m_{\alpha \beta}^{(\nu)} = \sum_k \frac{h_{\alpha k} h_{\beta k} M_k}{16 \pi^2} \left[ \frac{m_{H^0}^2}{m_{H^0}^2 - M_k^2} \ln \frac{m_{H^0}^2}{M_k^2} - \frac{m_{A^0}^2}{m_{A^0}^2 - M_k^2} \ln \frac{m_{A^0}^2}{M_k^2} \right].
\]

The fact that the light Majorana fermion is our DM candidate, the relic density value does not allow the new Yukawa couplings to be suppressed, and hence the smallness of neutrino mass is achieved by the mass degeneracy between \(m_{H^0}^2/m_{A^0}^2\), i.e., by considering very small \(\lambda_5 \sim O(10^{-10})\). The phenomenology of this setup is fully different from the case of IHDM or scotogenic model with scalar DM. In our case, the relic density is mainly dictated by the new Yukawa couplings, and the scalar parameter space is constrained only by the EW precision tests, vacuum stability and the perturbativity, which makes the collider signatures difference from the IHDM [2].

These models are subject of many theoretical and experimental constraints such as perturbativity, perturbative unitarity, vacuum stability, EW precision tests, LFV, the di-photon Higgs decay and the Higgs invisible decay. For instance, the existence of the charged scalar \(H^\pm\) in the scotogenic \((S_1^2\) in the KNT-like models) modifies the value of the branching ratio.
\[ B(h \rightarrow \gamma \gamma), \text{ where the ratio } R_{\gamma \gamma} := \frac{B(h \rightarrow \gamma \gamma)}{B(h \rightarrow \gamma \gamma)}^{SM} = 1.09 \pm 0.12 \text{ is reported by ATLAS and CMS [?]}. \text{ This ratio is given by in the scotogenic model by} \]

\[ R_{\gamma \gamma} = \left| 1 + \frac{\lambda_3 \mu^2}{2 m_H^2} A_{\gamma \gamma}^0 \left( \frac{m_H^2}{4 m_H^2} \right) \right|^2, \tag{10} \]

where the functions \( A_{\gamma \gamma}^0 \) are given in [?]. In the KNT-like models, \( m_{H^\pm}^2 \) is replaced by \( m_{S^\pm}^2 \), and \( \lambda_3 \) by the relevant quartic coupling \( |H|^2 |T|^2 \).

In the scotogenic and the KNT-like models, the LFV decay processes arise at the one-loop order with the exchange of \( H^\pm \) (or \( S^\pm \)) and \( N_k \) particles. The branching ratio of the decay \( \ell_\alpha \rightarrow \ell_\beta + \gamma \) due to the contribution of the interactions \( ?? \) is [?].

\[ B(\ell_\alpha \rightarrow \ell_\beta + \gamma) = \frac{3 \alpha \nu^4}{32 \pi m_{H^\pm}^2} \left| \sum_{i=1}^{3} h_{\beta i} h_{\alpha i} F \left( \frac{M_i^2}{m_{H^\pm}^2} \right) \right|^2, \tag{11} \]

where \( \alpha = \epsilon^2/4\pi \) is the electromagnetic fine structure constant and \( F(x) = (1 - 6x + 3x^2 + 2x^3 - 6x^2 \log x)/6(1-x)^4 \). We will consider also the LFV decays \( \ell_\alpha \rightarrow \ell_\beta \ell_\beta \ell_\beta \), where their branching ratio formulas are given in [?]. In our numerical scan, we will impose all the experimental limits on both \( B(\ell_\alpha \rightarrow \ell_\beta + \gamma) \) and \( B(\ell_\alpha \rightarrow \ell_\beta \ell_\beta \ell_\beta) \) [?]. The EW precision tests are easy to fulfill for the scotogenic and the KNT-like models, whereas they do not apply for the Dark Radiative Inverse Seesaw.

### 3. EW Phase Transition

It is well known that the SM cannot explain baryogenesis [?] for two reasons: (1) too small CP violating source in the CKM matrix and (2) weak electroweak phase transition. The latter is required to make the B+L violating processes suppressed in the broken phase during the bubble wall expansion. The EWPT is strongly first order only if [?],

\[ v(T_c)/T_c > 1, \tag{12} \]

which is not fulfilled in the SM, since this ratio is given by \( v_c/T_c \sim \lambda \). This last condition requires a Higgs mass below 42 GeV. Here \( T_c \) is the critical temperature at which the effective potential exhibits two degenerate minima, one at zero and the other at \( v(T_c) \). Both \( T_c \) and \( v(T_c) \) are determined using the full effective potential at finite temperature, which is given by [?]

\[ V_{\text{eff}}(h, T) = V_{\text{eff}}^{T=0}(h) + \frac{T^4}{2\pi^2} \sum_i n_i J_{B,F} \left( \frac{m_i^2}{T^2} \right) + V_{\text{ring}}(h, T); \tag{13} \]

\[ J_{B,F}(\alpha) = \int_0^\infty x^2 \log(1 + \exp(-\sqrt{x^2 + \alpha})). \tag{14} \]

The last contribution in (??) represents a leading contribution from the higher-order loop corrections, which can play an important role during the EWPT dynamics, which is the so-called daisy contributions [?]

\[ V_{\text{ring}}(h, T) = \frac{T}{12\pi} \sum_i n_i \left( \tilde{m}_i^3(h, T) - m_i^3(h) \right). \tag{15} \]

Here, the summation is performed over scalar and longitudinal gauge degrees of freedom, with \( \tilde{m}_i^2(h, T) = m_i^2(h) + \Pi(T) \) their thermal masses, and \( \Pi(T) \) are the thermal parts of the self-energies. In our work, we include this effect similar to [?], where the thermal masses replace
the field dependent masses of the scalar and longitudinal gauge fields in the full effective potential (16). In addition, we evaluate the integrals (17) numerically in order to account for all the (heavy and light) degrees of freedom.

By considering the tadpole condition at one-loop level, the Higgs mass is given by

\[ m_h^2 = 2\lambda \nu^2 + \frac{\nu^2}{32\pi^2} \sum_i n_i \alpha_i^2 \log \frac{m_i^2}{\Lambda^2}, \tag{16} \]

and the triple Higgs coupling can be simplified as

\[ \lambda_{hhh} = \frac{3m_h^2}{\nu} \left( 1 + \frac{\nu^4}{96\pi^2 m_S^2} \sum_i n_i \alpha_i^3 \right), \tag{17} \]

where all field dependent masses are written in the from

\[ m_i^2(h) = \mu_i^2 + \alpha_i^2 h^2/2. \]

For the case of SM extended by a scalar field \( S \) with multiplicity, \( N \), and a coupling to the Higgs doublet \( \omega \), the triple Higgs coupling gets enhanced by the ratio [?]

\[ \Delta_{hhh} = \frac{\lambda_{hhh} - \lambda_{hhh}^{SM}}{\lambda_{hhh}} = \frac{N \omega^3 \nu^2}{32\pi^2 m_S^2} \lambda_{hhh}^{SM} = 3 \times 10^{-6} N \left( \frac{\omega}{0.1} \right)^3 \left( \frac{m_S}{300 GeV} \right)^{-2}, \tag{18} \]

with the one-loop triple Higgs coupling in the SM is given by [?]

\[ \lambda_{hhh}^{SM} \approx \frac{3m_h^2}{\nu} \left[ 1 - \frac{m_t^4}{\pi^2 \nu^2 m_0^4} \right]. \tag{19} \]

According to (19), with more scalar degrees of freedom coupled to the Higgs doublet, the Higgs quartic coupling \( \lambda \) gets smaller and the EWPT gets stronger, in addition the Higgs triple coupling gets enhanced with respect to SM value.

### 4. Gravitational Waves from a Strong EW Phase Transition

In order to analyze the spectra of the gravitational waves (GWs) that are generated from first order EWPT in our model, we use the parameters, \( \alpha \) and \( \beta \), which characterize the GWs from the dynamics of vacuum bubble [?]. The parameter \( \beta \) describes the inverse of the time duration of the PT

\[ \beta = - \left. \frac{dS_E}{dt} \right|_{t=t_t} \simeq \left. \frac{1}{\Gamma} \frac{dT}{dt} \right|_{t=t_t}, \tag{20} \]

where \( S_E \) and \( \Gamma \) are the Euclidean action of a critical bubble and vacuum bubble nucleation rate per unit volume and unit time at the PT time \( t_t \), respectively. This parameter is usually normalized by the Hubble parameter \( H_T \) as

\[ \bar{\beta} = \frac{\beta}{H_T} = T_t \left. \frac{d}{dT} \left( \frac{S_3(T)}{T} \right) \right|_{T=T_t}, \tag{21} \]

where \( T_t \) is the transition temperature which is defined from \( \Gamma/H_T^3|_{T=T_t} = 1 \). The other parameter \( \alpha \) is proportional to the released energy density

\[ \epsilon(T) = -V_{eff}(\phi_t(T), T) + T \left. \frac{\partial V_{eff}(\phi_t(T), T)}{\partial T} \right|_{T=T_t}, \tag{22} \]
where $\phi_t(T)$ is the true minimum at the temperature $T$, which is normalized by the radiation energy density $\rho_{rad} = (\pi^2/30)g_\ast T^4$, where $g_\ast$ is relativistic degrees of freedom in the thermal plasma, at $T_t$ as

$$\alpha = \frac{\epsilon(T_t)}{\rho_{rad}(T_t)}.$$  \hspace{1cm} (23)

There are mainly three mechanisms to produce GWs during a strong first order PT occurs: (1) form the collisions of bubbles walls and shocks in the plasma, where the so-called “envelope approximation” [?] is a good approach to describe this phenomenon and estimate the contribution of the scalar field to the GWs spectrum. (2) Before the bubbles expansion has dissipated the kinetic energy in the plasma, and just after their their collision, the sound waves could result a significant contribution to the GWs spectrum [?]. (3) The magnetohydrodynamic (MHD) turbulence in the plasma that is formed after the bubbles collision may also have give rise to the GWs spectrum [?]. These three mechanisms could co-exist during the PT. Then the full stochastic GWs background is given by

$$\Omega_{GW} h^2 = \Omega_\phi h^2 + \Omega_{sw} h^2 + \Omega_{tur} h^2.$$  \hspace{1cm} (24)

In our analysis, we focus on the contribution to GWs from the compression waves in the plasma (sound waves), since it is the strongest GWs spectrum among the other GW sources. It can be fitted as [?]

$$\Omega_{sw}(f) h^2 = \tilde{\Omega}_{sw} h^2 \times (f/\tilde{f}_{sw})^3 \left(\frac{7}{4 + 3(f/\tilde{f}_{sw})^2}\right)^{7/2},$$  \hspace{1cm} (25)

where the peak energy density is

$$\tilde{\Omega}_{sw} h^2 \simeq 2.65 \times 10^{-6} v_b \tilde{\beta}^{-1} \left(\frac{\kappa \alpha}{1 + \alpha}\right)^2 \left(\frac{100}{g_\ast}\right)^{1/3},$$  \hspace{1cm} (26)

at the peak frequency, we have

$$\tilde{f}_{sw} \simeq 1.9 \times 10^{-5} \text{Hz} \frac{1}{v_b} \tilde{\beta} \left(\frac{T_t}{100 \text{GeV}}\right) \left(\frac{g_\ast}{100}\right)^{1/6},$$  \hspace{1cm} (27)

where $v_b$ is the velocity of the bubble wall and $\kappa$ is the fraction of convertable vacuum energy. It has been shown that a succesful electroweak baryogenesis and a sizable gravitational wave signal can be achived in the same setup [?], where the GWs signal is in the reach of BBO and LISA for very fine-tuned scenario [?]. Here, since we are interested in probing any possible GWs signal in a general SM extension that can map many models, we adopt the value $v_b = 0.95$. In order to investigate the effect of extra scalar degrees of freedom on the GWs spectrum in these models, as well in other SM extension by scalars, we consider the SM exyended by a scalar with mass $m_S$, multiplicity $N$, and a coupling to the Higgs doublet $\omega$.

5. Numerical Results

For the KNT-like and Dark Radiative inverse seesaw, we considered all the previous constraints, we find that the EWPT is easily first order as shown in Fig. ??.

It is clear that the strength of the EWPT gets improved when new bosonic degrees of freedom are present. For large values for the couplings of the doublet with extra scalars and/or small mass-values for the extra (singlet and multiplet) scalars, the one-loop correction to the Higgs mass in (??) is significant, and the Higgs self-coupling gets smaller. This helps to avoid Higgs
mass bound form the strong first order EWPT while fulfilling the criterion (??). In similar analyses [?, ?], it has been shown that extra scalars can help to generate a strongly first order EWPT due to: (a) the relaxation of the Higgs self-coupling $\lambda$ to be as small as $O(10^{-4})$; and/or (b) the enhancement of the value of the effective potential at the wrong vacuum at the critical temperature, without suppressing the ratio $v(T_c)/T_c$. This will easily relaxe the severe bound on the mass of the SM Higgs. For both models, the critical temperature is smaller or at the same order as in the SM. For the scotogenic model where DM candidate is a heavy Majorana fermion, the inert members mass are less restricted with respect the IHDM case with scalar DM. Then, one expect the same conclusion for the scotogenic model with DM Majorana candidate.

For the GWs spectrum, we consider different values for the multiplicity $N = 2, 6, 12, 24$, the scalar-Higgs doublet coupling range $|\omega| \leq 5$, and the scalar mass to lie in the range $m_S \in [100 \text{GeV}, 550 \text{GeV}]$. In Fig. ??-right, we present the predicted values of $\alpha$ and $\beta$. In the left panel, we show the region where the GW could be detected for $N = 12$.

In Fig. ??-left, the labels “DECIGO” and “Correlation” are DECIGO designs [?]. One notices that for stronger PT with $\phi_c/T_c \geq 1.38$ the generated GWs can be seen by DECIGO, and for $\phi_c/T_c \geq 2.95$ the corresponding GWs can be seen by LISA. Here, one has to mention two patterns of the EWPT, the usual one (type-I), and another case, where the EW symmetry is borken, then at lower temperature another transition to a deeper minimum occurs, which may leave detectable GWs (type-II). If the new singlet is electrically charged $Q_S \neq 0$, the EWPT could not be strongly first order while fulfilling the constraint from $h \rightarrow \gamma\gamma$ for large multiplicity values. Depending on the multiplicity $N$, the scalar mass should lie between 120 GeV and 380 GeV in order to have a strong first order EWPT, i.e., type-I PT, and therefore detectable GWs signal at DECIGO and may be LISA. Detectable GWs signal implies positive enhancement on the triple Higgs coupling (??) with ratio between 10% to more than 150%, depending on the multiplicity $N$. One can learn from these results that if the EWPT is strongly first order (type-I PT), the GWs signal would be able to be detected by the future space based GW interferometers, in addition to a non-negligible enhancement in triple Higgs coupling (??).

6. Conclusion
In many SM extensions motivated by DM and neutrino mass and oscillations, extra scalars are involved, where many collider and comlogical aspects are modified. In this work, we considered
Figure 5. Right: the predicted values of $\alpha$ and $\tilde{\beta}$ for different values of $\{N, m_S, \omega\}$. Right: the occurrence area of detectable GW spectrum together with different constraints: the vacuum stability condition of $\lambda \geq 0$ and the constraint on $R_{\gamma\gamma}^6$. The dotted lines, here, are the deviation of the triple coupling ($\gamma \gamma$), and the red region describes the vacuum stability condition. The green and yellow regions represent the current experimental data of $h \rightarrow \gamma \gamma$ at the accuracy of $1\sigma$ and $2\sigma$. The brown region represents the area of a strong PT without possible GWs detectability, and the black line shows the border of Type-I and Type-II PT.

3 SM extensions that address both DM and neutrino mass problems, and consider all theoretical and experimental constraints. We found that the EWPT could be easily strongly first order, as required for baryon asymmetry generation. In addition, the Higgs triple coupling could be significantly modified with respect to the SM, which may be tested at future colliders. Strong EWPT may generate detectable GWs.

We noticed also that GWs are large enough to be detected by future space-based interferometers, such as LISA and DECIGO, for the condition of strongly phase transition, in either type-I for a successful scenario of electroweak baryogenesis or type-II where baryogenesis can be fulfilled via another mechanism. This implies a significant deviation in the Higgs triple coupling with respect to the SM. Moreover, the detection of the GWs and the measurement of $h\gamma\gamma$ can test the model.

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