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Research Article

Strong Electroweak Phase Transition in a Model with Extended Scalar Sector

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We consider an extension of the Standard Model (SM) with additional gauge singlets which exhibits a strong first-order phase transition. Due to this first-order phase transition in the early universe gravitational waves are produced. We estimate the contributions such as the sound wave, the bubble wall collision, and the plasma turbulence to the stochastic gravitational wave background, and we find that the strength at the peak frequency is large enough to be detected at future gravitational interferometers such as eLISA. Deviations in the various Higgs boson self-couplings are also evaluated.

1. Introduction

The discovery of a narrow resonance, with a mass near 125 GeV, at the Large Hadron Collider (LHC) with properties similar to those of the Higgs boson predicted by the Standard Model (SM) [1, 2], sparked a lot of excitements among high energy physicists. But in spite of this important discovery the SM is considered to be incomplete. For instance, in the Standard Model of particle physics a strong first-order phase transition (SFPT) does not occur [3]. However SFPT is needed to justify the baryon asymmetry of our universe [4]; moreover, the SM does not have a candidate for dark matter (DM).

Therefore, some new models are required to address these issues. A popular model is to couple a singlet scalar to Higgs boson. In [5] it is shown that it is possible to modify the standard theory by adding a scalar which possesses a discrete Z_2 symmetry and to address the issue of the dark matter of the universe, within the frame work of singlet extended SM; the issue of dark matter has been studied in [6–13], while electroweak phase transition was studied in [14–23] and in [24, 25]; the authors attempt to explain electroweak phase transition EWPT and dark matter by singlet extended SM.

Another class of models are the multisinglet extensions of the SM [26–36]. These models have a larger parameter space in comparison to the singlet extended models; hence they can address several issues; in [26, 27] cosmological implications of such models with classical conformal invariance are presented. Electroweak phase transitions in two-Higgs doublet model are analyzed in [37, 38] and within supersymmetric models in [39–45]. A comprehensive review of EWPT within various models has been given in [46].

In order to investigate the dynamics of the electroweak phase transition EWPT one has to utilize techniques from the domain of thermal field theory [47–50]. The occurrence of a first-order phase transition requires that the electroweak breaking and preserving minima be degenerate, an event which happens at a critical temperature T_c. Moreover, to prevent the washout of any baryon asymmetry by electroweak sphalerons, the electroweak phase transition must be strongly first order. Namely, the ratio of vacuum expectation value of the Higgs field to the critical temperature needs to be greater than unity. As we describe in the next section if a first-order phase transition occurs in the early universe, the dynamics of bubble collision and subsequent turbulence of the plasma are expected to generate gravitational wave (GW). If we detect these GW, then we can obtain information about symmetry breaking in early universe. GW signals from phase transitions have been discussed in [51–59].

In [60] the authors study EWPT within several exotic models and in [61] an analysis of the EWPT of a large number of minimal extensions of the SM and their classically conformal limits is presented. Complex conformal singlet extension of SM with emphasis on the issue of dark matter and Higgs phenomenology is studied in [62]. Recently baryogenesis within a φ^6 model is addressed in [63]. An investigation of electroweak phase transitions in a singlet extended model in the 100 (TeV) range is given in [64]. The authors of [65] study strong first-order EWPT in a singlet scalar extension of the SM where the singlet scalar is coupled nonminimally to gravity. In this scheme the singlet field first derives inflation and at a later time causes a strong EWPT; in a new study a first-order EWPT in the SM is obtained by varying Yukawas during phase transition [66].

In this work we propose a new model and we investigate the strength of EWPT within this model. In this model which is a generalization of [19], N real gauge singlets are coupled to Higgs boson via trilinear interactions. Previous studies of multiscalar singlet extension of SM impose separate Z_2 symmetries on the singlets [26, 28, 29, 31, 33, 35, 36]; however in our model we do not require such symmetry. In spite of its simple form, this model has a very rich phenomenology. The main feature of the singlet extended SM (without Z_2 symmetry) is that the potential barrier between the true and the false vacua necessary for first-order EWPT can be formed mainly by tree-level interactions. But in the singlet extended SM (with Z_2 symmetry) nondecoupling loop effects are needed for the occurrence of a strong first-order EWPT; however, these models have DM candidate. In [67] a nonminimal composite model based on the coset $SO(7)/SO(6)$ has been considered. At low energy the scalar sector of their model is composed of two scalars, one with an unbroken Z_2 symmetry and another scalar with a broken Z_2 symmetry.

The plan of this paper is as follows.

In Section 2 we summarize the basic notions of EWPT. We describe the finite temperature effective potential at one loop. Then by emphasizing the underlying physical mechanisms, we describe the basic quantities of interest such as the strength of a phase transition, the rate of variation of bubble nucleation rate per volume, and the ratio of the latent heat released at the phase transition to the radiation energy density. In Section 3 we consider a simple extension of the SM by the addition of N real scalar gauge singlets with trilinear coupling to Higgs. We present the phenomenology of the model for $N = 2$ and due to lack of protective symmetry these gauge singlets are not candidates for DM and we discuss the issue of observability of gravitational wave of our model. And finally in Section 4 we present our conclusions. Technical details are explained in the Appendix.

2. Electroweak Phase Transition

In this section we summarize basic notions and definitions of the electroweak phase transitions.

The observable universe consists predominantly of matter. The asymmetry between the matter and antimatter content of the universe is expressed by the baryon to photon ratio

$$\rho = \frac{n_b - n_{\bar{b}}}{n_\gamma} \sim 10^{-9}, \tag{1}$$

where n_b, $n_{\bar{b}}$, and n_γ are the number densities of baryons, antibaryons, and photons. In a symmetric universe one expects $\rho = 0$ but experiments reveal that ρ has a tiny but nonzero value. This paradox can be resolved by requiring baryon number violation, C and CP violation, and departure from thermal equilibrium [4].

Baryogenesis is the physical process which is responsible for this observed baryon asymmetry of universe (BAU). In electroweak baryogenesis one assumes that the physical mechanism is the occurrence of a strong first-order electroweak phase transition (EWPT); namely, a smooth or a weak phase transition can not explain BAU [68, 69].

A convenient tool for investigation of EWPT is effective potential. For any quantum field theory if we replace the quantum field by its vacuum expectation value in the presence of a source the result for the potential energy to lowest order in perturbation theory is called the effective potential. The physical meaning of the effective potential is that it represents an energy density. In general effective potential contains other terms (loop corrections) [70, 71]. By using path integrals it is possible to find effective potential. In an Euclidean space-time at one-loop order the result is

$$V^{\text{eff}}(\varphi_c) = V_0(\varphi_c) + \frac{i}{2}\int \frac{d^4 p}{(2\pi)^4} \log\left[\frac{p^2 + m^2(\varphi_c)}{p^2}\right] + O\left(\hbar^2\right) + \cdots, \tag{2}$$

which is known as the Coleman-Weinberg potential. The first term in (2) is the classical tree-level potential.

While considering physical processes in a hot environment such as early universe a proper method is finite temperature field theory [47–50]. The expression for the one-loop effective potential at finite temperature is [72]

$$V_1^{\text{eff}}(\varphi_c, T) = \sum_i \frac{n_i T^4}{2\pi^2} J_\mp\left(\frac{m_i(\varphi_c)}{T}\right), \tag{3}$$

where n_i is the number of degrees of freedom of the particle, m_i is the field dependent particle mass, and

$$J_\mp = \pm\int_0^\infty dy y^2 \log\left[1 \mp \exp\left(-\sqrt{x^2 + y^2}\right)\right], \tag{4}$$

and $J_-(x)$ and $J_+(x)$ denote the contribution from bosons and fermions. In addition there is another contribution to thermal effective potential from the daisy subtraction [73, 74].

The best way for the evaluation of the finite temperature effective potential is to use special packages (codes) developed for a specific model.

But when the temperature is much greater than the masses of the various particles of the system under consideration it is possible to expand $J_\mp$ in a power series of m/T. In this work we use this high-temperature approximation. Recently, a new scheme for computation and resummation of thermal masses beyond the high-temperature approximation in general beyond SM scenarios has been proposed [75].

Let us consider the shape of the effective potential. For a generic model, as the universe cooled down at temperature above a critical temperature T_c, the effective potential had an absolute minimum which was located at the origin. At T_c there were two degenerate minima, which were separated by an energy barrier. At temperature below the critical temperature the second minimum became the global one and presently there is no energy barrier. Moreover the rate of expansion of the universe slowed, as the Hubble parameter H which characterizes the rate of expansion of the universe depends quadratically on the temperature.

In 1976 't Hooft discovered that baryon number is violated in the Standard Model and it is due to transition between two topologically distinct $SU(2)_L$ ground states [76, 77]. While, at zero temperature, the probability for barrier penetration is vanishingly small, at nonzero temperature the transition between two ground states differing by a unit of topological charge can be achieved by a classical motion over the barrier. Unstable static solutions of the field equations with energy equal to the height of the barrier separating two topologically distinct $SU(2)_L$ ground states (sphalerons) have been reported in [78]. Hence in the electroweak baryogenesis, the baryon asymmetry of universe is generated through the sphaleron process in the symmetric phase during the EWPT. At the symmetric phase the rate of baryon violating process which we denote by $\widetilde{\Gamma}^{\text{sph}}$ is a quartic function of the temperature; hence in this phase $\widetilde{\Gamma}^{\text{sph}} \gg H$.

But the third conditions to generate BAU is the departure from thermal equilibrium. Hence, the baryon number changing sphaleron interaction must quickly decouple in the broken phase; that is, $\widetilde{\Gamma}^{\text{sph}} < H$. But the rate of sphaleron induced baryon violating process is suppressed by a Boltzmann factor

$$\widetilde{\Gamma}^{\text{sph}} \propto \exp\left[-\frac{E_{\text{sph}}(T)}{T}\right], \tag{5}$$

but at phase transition $E_{\text{sph}}(T) \propto \varphi_c$

$$\frac{\varphi_c}{T_c} \geq 1, \tag{6}$$

where φ_c is the broken phase minimum at the critical temperature T_c.

A first-order phase transition proceeds by nucleation of bubbles of the broken symmetry phase within the symmetric phase. The underlying mechanisms for bubble nucleation are quantum tunneling and thermal fluctuations. These bubbles then expand, merge, and collide. Gravitational waves are produced due to collision of bubble walls and turbulence in the plasma after the collisions. In addition while these bubbles pass through the plasma sound waves are created. These sound waves can provide additional sources of gravitational waves.

A crucial parameter for the calculation of the gravitational wave spectrum is the rate of variation of the bubble nucleation rate per volume, called β. It is common to use a normalized dimensionless parameter which is defined as

$$\widetilde{\beta} = \frac{\beta}{H_*}, \tag{7}$$

where H_* denotes the Hubble parameter at the time of phase transition.

The density of latent heat released into the plasma is

$$\epsilon_* = \left[-V_{\min}^{\text{eff}}(T) + T\frac{d}{dT}V_{\min}^{\text{eff}}(T)\right]_{T=T_*}, \tag{8}$$

where $V_{\min}^{\text{eff}}(T)$ is the temperature-dependent true minimum of the effective potential of the scalar fields which causes the phase transition; moreover, its value must be set to zero by adding a constant at each time. Another dimensionless parameter for characterizing the spectrum of gravitational wave is

$$\alpha = \frac{\epsilon_*}{\rho_{\text{rad}}}, \tag{9}$$

where the radiation energy density of the plasma $\rho_{\text{rad}} = (\pi^2/30)g_* T_c^4$. And the parameter g_* is the effective degrees of freedom in the thermal bath at the phase transition. In this work we assume $g_* = 106.75 + N_S$, where N_S denotes the number of singlet scalars that facilitates the electroweak phase transition.

3. The Model

In [19] the effects of a light scalar on the electroweak phase transition have been considered. In their model the scalar sector has been extended by an addition of a singlet, which has a trilinear interaction with the Higgs boson. Recently an extension of the SM with addition of N isospin-singlet, which has a quartic interaction with the Higgs boson, has been considered [57]. Here we consider a generalization of model of [19]. At zero temperature the effective potential of the scalar sector of our model is

$$V_0 = -DT_0^2\varphi^2 + \frac{\lambda}{4}\varphi^4 + \frac{1}{2}\sum_{i=1}^{N} m_i^2 s_i^2 + \varphi^2\sum_{i=1}^{N}\kappa_i s_i. \tag{10}$$

Since we have not included quartic self-coupling for the extra scalars, in order to have a stable potential, we assume that the squares of the mass parameters of the extra scalars (m_i^2) are positive definite (see the Appendix for details).

At high temperature we have

$$\begin{aligned} V_T = {} & D\left(T^2 - T_0^2\right)\varphi^2 - ET\varphi^3 + \frac{\lambda_T}{4}\varphi^4 + \frac{1}{2}\sum_{i=1}^{N} m_i^2 s_i^2 \\ & + \varphi^2\sum_{i=1}^{N}\kappa_i s_i + \frac{T^2}{12}\sum_{i=1}^{N}\kappa_i s_i. \end{aligned} \tag{11}$$

The parameters of (11) are given by

$$D = \frac{1}{8v^2}\left(2m_W^2 + m_Z^2 + 2m_t^2 + 2\lambda v^2\right),$$

$$E = \frac{1}{8\pi v^3}\left(4m_W^3 + 2m_Z^3\right),$$

$$\lambda_T = \lambda - \frac{1}{16\pi^2 v^4}\left(6m_W^4 \ln\frac{m_w^2}{a_B T^2} + 3m_Z^4 \ln\frac{m_Z^2}{a_B T^2} - 12m_t^4 \ln\frac{m_t^2}{a_F T^2}\right),$$
$$\ln a_B = 3.91,$$
$$\ln a_F = 1.14. \tag{12}$$

The critical temperature of model is given by

$$T_c = \frac{T_0}{\sqrt{1 - E^2/D\left(\lambda_T - 2\zeta\right) - \zeta/12D}}, \tag{13}$$
$$\text{where } \zeta = \sum_{i=1}^{N}\frac{\kappa_i^2}{m_i^2}.$$

The strength of the phase transition is denoted by ξ and it is given by

$$\xi = \frac{\varphi_c}{T_c} = \frac{2E}{\lambda_T - 2\zeta}. \tag{14}$$

Moreover, s_{ic} is the *vev* of the scalar field s_i at the second minimum of the effective potential at T_c:

$$s_{ic} = -\frac{\kappa_i}{m_i^2}\left(\varphi_c^2 + \frac{T_c^2}{12}\right). \tag{15}$$

3.1. Phenomenology of the Models with N Trilinear Interactions. The structure of the scalar mass matrix is

$$\begin{pmatrix} 2\lambda v^2 & 2\kappa_1 v & 2\kappa_2 v & 2\kappa_3 v & \cdots & \cdots \\ 2\kappa_1 v & m_1^2 & 0 & 0 & 0 & 0 \\ 2\kappa_2 v & 0 & m_2^2 & 0 & 0 & 0 \\ 2\kappa_3 v & & 0 & m_3^2 & \cdots & \cdots \\ \cdots & \cdots & \cdots & \cdots & \cdots & \cdots \\ \cdots & \cdots & \cdots & \cdots & \cdots & \cdots \end{pmatrix}. \tag{16}$$

The physical mass squares of the scalar sector of our model are the eigenvalues of this matrix. The characteristic equation is

$$\det\begin{pmatrix} \omega - 2\lambda v^2 & -2\kappa_1 v & -2\kappa_2 v & -2\kappa_3 v & \cdots & \cdots \\ -2\kappa_1 v & \omega - m_1^2 & 0 & 0 & 0 & 0 \\ -2\kappa_2 v & 0 & \omega - m_2^2 & 0 & 0 & 0 \\ -2\kappa_3 v & & 0 & \omega - m_3^2 & \cdots & \cdots \\ \cdots & \cdots & \cdots & \cdots & \cdots & \cdots \\ \cdots & \cdots & \cdots & \cdots & \cdots & \cdots \end{pmatrix} = 0. \tag{17}$$

Hence in general one should solve a polynomial of degree $(N + 1)$ of ω. From the above equation we see that the coefficient of ω^{N+1} is unity. If we designate the coefficient of ω^N by α, then $-\alpha$ is equal to the sum of physical masses of the model; namely,

$$2\lambda v^2 + \sum_{i=1}^{N} m_i^2 = m_H^2 + \sum_{i=1}^{N}\chi_i^2, \tag{18}$$

where the physical mass of the ith scalar is denoted by χ_i. From this relation we find an important relation about the Higgs self-coupling; namely,

$$\lambda - \lambda_{\mathrm{SM}} = \frac{1}{2v^2}\sum_{i=1}^{N}\left(\chi_i^2 - m_i^2\right), \tag{19}$$

where λ_{SM} is the Higgs quartic self-coupling of the Standard Model.

We see that if the mass parameters of the scalars as well as the physical masses of the scalars are much smaller than v, then the deviation of the parameter λ from the Higgs self-coupling of the Standard Model will be very small.

3.2. The Special Case $N = 2$. Here we consider the case $N = 2$. The scalar mass matrix of the model at zero temperature for this case is

$$M^2 = \begin{pmatrix} 2\lambda v^2 & 2\kappa_1 v & 2\kappa_2 v \\ 2\kappa_1 v & m_1^2 & 0 \\ 2\kappa_2 v & 0 & m_2^2 \end{pmatrix}. \tag{20}$$

The physical mass squares of the model can be obtained from

$$\omega^3 + A\omega^2 + B\omega + C = 0, \tag{21}$$

where

$$A = -\left(2\lambda v^2 + m_1^2 + m_2^2\right),$$
$$B = 2v^2\left[\lambda\left(m_1^2 + m_2^2\right) - 2\left(\kappa_1^2 + \kappa_2^2\right)\right] + m_1^2 m_2^2, \tag{22}$$
$$C = 4v^2\left(\kappa_1^2 m_2^2 + \kappa_2^2 m_1^2\right) - 2\lambda^2 v^2 m_1^2 m_2^2.$$

Now one of the eigenvalues is equal to m_H^2; hence we obtain

$$m_H^6 + Am_H^4 + Bm_H^2 + C = 0. \tag{23}$$

Moreover, by minimizing the potential with respect to variables s_1, s_2, we obtain

$$s_i = -\frac{\kappa_i v^2}{m_i^2}, \quad i = 1, 2. \tag{24}$$

By minimizing the potential with respect to variable φ we obtain

$$-2DT_0^2 + \lambda v^2 - 2\sum_{i=1}^{2}\frac{\kappa_i^2 v^2}{m_i^2} = 0. \tag{25}$$

The parameters of the model must satisfy (23) and (25). Hence the parameter space of the model in this case contains four independent parameters $(\kappa_1, \kappa_2, m_1, m_2)$.

Table 1: Different configurations associated with the onset of a strong first-order phase transition ($\xi = 1$).

Set	κ_1	κ_2	m_1 (GeV)	m_2 (GeV)	χ_1 (GeV)	χ_2 (GeV)	$\cos(\varphi)$	$\Delta\lambda_{\text{hhh}}$	$\Delta\lambda_{\text{hhhh}}$
I	0.2	3.7	69.2	16.0	6.2	69.2	0.993	13.9%	4.1%
II	0.5	3.1	18.7	13.5	5.2	18.6	0.995	14.5%	3.0%
III	0.8	2.3	11.2	10.4	4.0	11.1	0.997	15.1%	1.8%
IV	1.2	0.7	18.1	3.2	17.5	1.3	0.999	5.7%	0.6%

Table 2: Values of Higgs invisible branching ratios for various decay modes, total value of Higgs invisible branching ratio, Higgs invisible width, Higgs total width, and deviation of the Higgs width from the SM value are shown for various configurations. The unit for width in this table is MeV. All four configurations are consistent with current experimental data.

Set	$B(h \to s_1 s_1)$	$B(h \to s_1 s_2)$	$B(h \to s_2 s_2)$	$B_{\text{tot}}(h \to \text{inv})$	$\Gamma_{\text{inv}}(h)$	$\Gamma_{\text{tot}}(h)$	$\Delta\Gamma_{\text{tot}}$
I	1.45%	0.31%	0.0%	1.49%	0.06	4.006	0.006
II	0.75%	0.04%	3.6×10^{-5}%	0.8%	0.03	3.992	0.008
III	0.24%	0.004%	1.3×10^{-6}%	.24%	0.01	3.99	0.01
IV	0.02%	4.6×10^{-7}%	6.3×10^{-12}%	0.02%	5.8×10^{-4}	3.993	0.007

If we subtract (23) from (21) we get

$$\chi^4 + \left(A + m_H^2\right)\chi^2 + m_H^4 + m_H^2 A + B = 0. \quad (26)$$

Therefore, the physical mass squares of the singlets are determined.

But models with extended Higgs sectors predicting strongly first-order phase transition simultaneously predict a significant deviation in the triple Higgs boson coupling as well [79]. This deviation at the tree level in [17] and at loop level in [33, 79] has been studied, with

$$\Delta_{\text{hhh}} = \frac{\lambda_{\text{hhh}}^{\text{MSM}} - \lambda_{\text{hhh}}^{\text{SM}}}{\lambda_{\text{hhh}}^{\text{SM}}}, \quad (27)$$

where $\lambda_{\text{hhh}}^{\text{SM}}$ is the Higgs triple coupling of the SM and $\lambda_{\text{hhh}}^{\text{MSM}}$ is the Higgs triple coupling of the multisinglet extension of the SM. Collider experiments could measure the Higgs triple coupling. Here we want to explore the region of intermediate mass of the singlets. But there are bounds on Higgs-Portal models from the LHC Higgs data [80–82]. For instance, the Higgs doublet is mixed with the extra singlet scalars and the mixing element $\cos(\varphi)$ between the CP-even component of the doublet (φ) and the physical Higgs (whose $m_H = 125.09$ GeV) is not arbitrary and it is subject of a constraint coming from the Higgs coupling to the W gauge bosons. Current data [81] suggests $\cos(\varphi) > 0.86$.

The results are presented in Table 1 for the onset of a strong EWPT; namely, $\xi = 1$. For each configuration in the table we present our results for the deviation of Higgs triple coupling as well. Hence by adding one scalar to the model we can have singlets in the intermediate mass region.

For completeness the predictions of the model for the deviation of Higgs boson quartic coupling $\Delta\lambda_{\text{hhhh}}$ is given in Table 1.

Moreover, the existence of the extra singlet scalars could affect the total Higgs decay if they are light enough, which becomes

$$\Gamma_{\text{total}}(h) = \cos^2(\varphi)\,\Gamma_{\text{total}}(h_{\text{SM}}) + \sum \Gamma\left(h \longrightarrow s_i + s_k\right), \quad (28)$$

where s denote all the scalars. The deviation from the SM value $\Delta\Gamma_{\text{total}} < 1.4$ (MeV).

In the SM, a total Higgs decay width around 4 MeV is predicted. In this work we assume the decay $h \to s_i + s_k$ is an invisible decay. However, current analysis [82] suggests that the Higgs invisible decay branching ratio should be less than 17%. In Table 2 we present Higgs invisible decay branching ratio in various decay modes, as well as the total Higgs invisible branching ratio for all of Higgs invisible decay modes. The invisible Higgs width, total Higgs width, and the deviation of total Higgs width of our model from that of the SM are also shown. The unit for the various width in this table is MeV; moreover in our calculation we have assumed $\Gamma_{\text{total}}(h_{\text{SM}}) = 4$ (MeV).

It turns out that the critical temperature for this model $T_c < 100$ (GeV). Hence it is a good approximation to use λ instead of λ_T and in [19] this approximation is used to study SFPT for the case $N = 1$ but the mass of the light scalar is up to 12 GeV; in [20] a one-loop study of the same model has been presented but the mass of the light scalar to catalyze a SFPT is up to 20 GeV. But for the model presented in this work and using this approximation for the case $N = 2$ the mass of the scalar to catalyze a SFPT is 69 GeV. By adding more scalars we expect to have heavy singlets.

3.3. Detection of Gravitational Waves. The direct detection of gravitational waves by LIGO [83] generated a lot of interest among researchers in cosmology, astrophysics, and particle physics. Major sources of gravitational waves are inflation, compact binary systems, or cosmological phase transitions.

Table 3: The spectra of gravitational wave as predicted by our model. The emitted GW are within the reach of eLISA C1.

GW spectra	f_{sw}	f_{col}	f_{turb}	Ωh^2_{sw}	Ωh^2_{col}	Ωh^2_{turb}	$\Omega h^2_{\text{total}}$
$\widetilde{\beta} = 50$	0.65	0.12	1.85	2.34×10^{-11}	1.98×10^{-13}	1.67×10^{-15}	2.36×10^{-11}
$\widetilde{\beta} = 100$	1.30	0.24	3.69	1.17×10^{-11}	9.9×10^{-14}	8.35×10^{-16}	1.18×10^{-11}
$\widetilde{\beta} = 250$	3.25	0.60	9.23	4.67×10^{-12}	3.96×10^{-14}	3.34×10^{-16}	4.71×10^{-12}

In this section under phenomenological description we aim to make an estimate for the observability of the gravitational waves produced by the model presented in this work.

In all of previous multisinglet models a Z_2 symmetry (either broken or unbroken) is imposed on the fields. However, neither of the two fields of our model has this symmetry. This is the main difference between our model and previous studies.

The main parameters that are of value for obtaining the spectrum of gravitational waves are the parameters α and β, which we discussed in Section 2. And they are obtained from the thermal effective potential.

For the set I of Table 1 and from (8) and (9) we obtain $\alpha = 0.16$. But the velocity of the bubble wall and the efficiency factor, the fraction of the latent heat which is converted to the kinetic energy of the plasma, are determined from

$$v_b = \frac{1/\sqrt{3} + \sqrt{\alpha^2 + 2\alpha/3}}{1+\alpha},$$
$$\kappa = \frac{1}{1+0.715\alpha}\left(0.715\alpha + \frac{4}{27}\sqrt{\frac{3\alpha}{2}}\right). \tag{29}$$

Thus for this configuration $v_b = 0.81$ and $\kappa = 0.17$. The peak frequency for bubble collision contribution is given by [84]

$$f_{\text{col}} = 16.5 \times 10^{-6} \frac{0.62}{v_b^2 - 0.1v_b + 1.8} \frac{\beta}{H_*} \frac{T_*}{100} \left(\frac{g_*}{100}\right)^{1/6}. \quad \text{Hz}. \tag{30}$$

And the energy density at this frequency is given by

$$\Omega h^2_{\text{col}} = 1.67 \times 10^{-5} \left(\frac{\beta}{H_*}\right)^2 \frac{0.11 v_b^3}{0.42 + v_b^2} \left(\frac{\kappa\alpha}{1+\alpha}\right)^2 \left(\frac{g_*}{100}\right)^{-1/3}, \tag{31}$$

where h is the reduced Hubble constant at present.

Another source of gravitational waves is the compression waves in the plasma (sound waves) and the peak frequency is

$$f_{\text{sw}} = 1.9 \times 10^{-5} \frac{\beta}{H_*} v_b^{-1} \frac{T_*}{100} \left(\frac{g_*}{100}\right)^{1/6}, \quad \text{Hz}; \tag{32}$$

the energy density at this peak frequency is obtained from

$$\Omega h^2_{\text{sw}} = 2.65 \times 10^{-6} \left(\frac{\beta}{H_*}\right)^{-1} \left(\frac{\kappa\alpha}{1+\alpha}\right)^2 \left(\frac{g_*}{100}\right)^{-1/3} v_b. \tag{33}$$

And finally the peak frequency for gravitational waves which are caused by the turbulence of the plasma is

$$f_{\text{turbo}} = 2.7 \times 10^{-5} \frac{\beta}{H_*} v_b^{-1} \frac{T_*}{100} \left(\frac{g_*}{100}\right)^{1/6} \quad \text{Hz}, \tag{34}$$

and the peak energy density of this part is given by

$$\Omega h^2_{\text{turbo}} = 3.35 \times 10^{-4} \left(\frac{\beta}{H_*}\right)^{-1} \left(\frac{\epsilon\kappa\alpha}{1+\alpha}\right)^{3/2} \left(\frac{g_*}{100}\right)^{-1/3} \cdot v_b \frac{1}{2^{11/3}\left(1 + 8\pi\left(f_{\text{turbo}}/H_*\right)\right)}, \tag{35}$$

where ε denotes the fraction of latent heat that is transformed into turbulent motion of the plasma. We choose $\varepsilon = 0.05$.

Hence by knowing the values of the parameters α, $\widetilde{\beta}$, and T_*, we can obtain the spectra of the GW. Even though the nucleation temperature T_* is lower than the critical temperature in this work we assume $T_* \approx T_c$.

The standard method of calculation of the parameter $\widetilde{\beta}$ from the effective potential is to compute the Euclidean action of the model and it is explained in [36, 85], but in an approximate scheme it is found that [59]

$$\widetilde{\beta} \approx 170 - 4\ln\left(\frac{T_*}{1\,\text{GeV}}\right) - 2\ln g_*, \tag{36}$$

and in our case when the phase transition happens at weak scale, $\widetilde{\beta} \approx 144$.

Hence, in order to assess the implications of the model on the spectra of GW we present results by varying this parameter in the interval $50 \leq \widetilde{\beta} \leq 250$. In Table 3 we present our results. The unit for various frequencies is mHz.

The results of Table 3 show that in our model the contribution from the turbulent motion of always a few orders of magnitude is smaller than the previous two. Moreover, the contribution from the sound waves is the dominant source of the total GW spectrum. By studying the frequency dependent spectra [84], we find that the peak of energy density of the sound wave contribution and the peak of energy density due to collision are well separated, a desirable feature while detecting these waves. These waves are within the reach of future gravitational wave interferometers (eLISA C1).

4. Conclusions

In this work we have presented a new extension of the SM. In this model we amend the SM by N gauge singlets. In order to avoid proliferation of the parameters we considered the most economical model. For each scalar we allowed a mass term with a positive squared mass parameter to insure vacuum stability in the direction of that scalar and a triple coupling with the Higgs field to facilitate a strong electroweak phase transition, and for the special case $N = 2$ we find SFPT with gauge singlets in the intermediate mass range. We also investigated the deviations of the Higgs coupling constants from the SM values. And we find that the deviation of the triple Higgs boson coupling can be as large as 15%.

Finally, we have obtained the gravitational wave spectrum from the electroweak phase transition. We have shown that the gravitational wave signal can be detected by eLISA. The present model has a large parameter space and as a result a richer phenomenology in comparison to [19]. It would be of interest to amend the model by inclusion of quartic self-coupling of the scalars as well as the mixing term between scalars (see the Appendix for details).

It would be of interest to consider higher values of N for the model proposed in this work. These and other related issues are presently under consideration.

Appendix

Vacuum Stability Conditions

For the special case $N = 2$ of the model the weak condition for vacuum stability is that if the value of the fields tends to infinity, then the tree-level potential is not unbounded-from-below directions. This leads to $\lambda > 0$. But if only the field s_1 tends to infinity, then the condition vacuum stability will be $m_1^2 > 0$ and by similar argument for the field s_2 the restriction will be $m_2^2 > 0$.

However, the strong condition for vacuum stability as stated in [35, 86, 87] is that the masses of the particles of a particular model are positive.

The necessary conditions for a symmetric matrix A of order 3 to have real eigenvalues are

$$\begin{aligned} & a_{11} > 0, \\ & a_{22} > 0, \\ & a_{33} > 0, \\ & \bar{a}_{12} = a_{12} + \sqrt{a_{11}a_{22}} > 0, \\ & \bar{a}_{13} = a_{13} + \sqrt{a_{11}a_{33}} > 0, \\ & \bar{a}_{23} = a_{23} + \sqrt{a_{22}a_{33}} > 0, \end{aligned} \tag{A.1}$$

$$\begin{aligned} & \sqrt{a_{11}a_{22}a_{33}} + a_{12}\sqrt{a_{33}} + a_{13}\sqrt{a_{22}} + a_{23}\sqrt{a_{11}} \\ & \quad + \sqrt{2\bar{a}_{12}\bar{a}_{13}\bar{a}_{23}} > 0. \end{aligned} \tag{A.2}$$

The constraints stated in (A.1) apply to quantities with dimension of squared mass. From our mass matrix of Section 3.2, the first constraint is $2\lambda v^2 > 0$, which leads to $\lambda > 0$. Hence, for the model presented in this work, the criteria which provide the necessary and sufficient vacuum stability conditions are given by

$$\begin{aligned} & \lambda > 0, \\ & m_1^2 > 0, \\ & m_2^2 > 0, \\ & \kappa_1 > -\sqrt{\frac{\lambda m_1^2}{2}}, \\ & \kappa_2 > -\sqrt{\frac{\lambda m_2^2}{2}}, \end{aligned} \tag{A.3}$$

$$\begin{aligned} & \sqrt{\lambda m_1^2 m_2^2} + \kappa_1\sqrt{2m_2^2} + \kappa_2\sqrt{2m_1^2} \\ & \quad + \sqrt{\left(2\kappa_1 + \sqrt{2\lambda m_1^2}\right)\left(2\kappa_2 + \sqrt{2\lambda m_2^2}\right)\sqrt{m_1^2 m_2^2}} \\ & \quad > 0, \end{aligned} \tag{A.4}$$

where the conditions stated in the first line of (A.3) have been obtained by using weak conditions for vacuum stability.

In our model we have not included quartic terms for the gauge singles such as

$$V_{\text{quartic}} = \sum_{i=1}^{2} \frac{\lambda_i s_i^4}{4} + \delta s_1^2 s_2^2 \tag{A.5}$$

in the effective potential, where λ_1, λ_2, and δ are dimensionless coupling parameters; in this case the vacuum stability in the direction of the extra scalars is maintained if

$$\begin{aligned} \lambda_1 &> 0, \\ \lambda_2 &> 0, \\ \sqrt{\lambda_1\lambda_2} &> -2\delta, \end{aligned} \tag{A.6}$$

and the squares of the mass parameters m_1^2 and m_2^2 in principle can assume any value (positive, zero, or negative); for instance, in the multiscalar model described in [26] all of the extra scalars do not have a mass term, while in the absence of quartic terms the squared mass parameters of the model presented in Section 3 must be positive as required by vacuum stability.

Conflicts of Interest

The authors declare that they have no conflicts of interest.

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