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Gravitational waves and Higgs boson couplings for exploring first order phase transition in the model with a singlet scalar field

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ABSTRACT

We calculate the spectrum of gravitational waves originated from strongly first order electroweak phase transition in the extended Higgs model with a real singlet scalar field. In order to calculate the bubble nucleation rate, we perform a two-field analysis and evaluate bounce solutions connecting the true and the false vacua using the one-loop effective potential at finite temperatures. Imposing the Sakharov condition of the departure from thermal equilibrium for baryogenesis, we survey allowed regions of parameters of the model. We then investigate the gravitational waves produced at electroweak bubble collisions in the early Universe, such as the sound wave, the bubble wall collision and the plasma turbulence. We find that the strength at the peak frequency can be large enough to be detected at future space-based gravitational interferometers such as eLISA, DECIGO and BBO. Predicted deviations in the various Higgs boson couplings are also evaluated at the zero temperature, and are shown to be large enough too. Therefore, in this model strongly first order electroweak phase transition can be tested by the combination of the precision study of various Higgs boson couplings at the LHC, the measurement of the triple Higgs boson coupling at future lepton colliders and the shape of the spectrum of gravitational wave detectable at future gravitational interferometers.

By the discovery of a Higgs boson [1,2] and the dedicated measurements of its property at the LHC, the mass generation mechanism for elementary particles in the standard model (SM) has been established. One of the next important targets of high energy physics is to explore the structure of the Higgs sector, dynamics of electroweak symmetry breaking (EWSB), and the nature of the Higgs boson (h).

In addition, the mechanism of electroweak phase transition (EWPT) is still a mystery, that is strongly related not only to the physics behind EWSB but also to various cosmological problems such as baryon asymmetry of the Universe and cosmic inflation. In particular, the strongly first order phase transition (1stOPT) is crucial for a successful scenario of electroweak baryogenesis (EWBG) [3]. With another requirement of additional CP violating phase, the EWBG scenario can be realized by introducing an extended Higgs sector. Therefore, in this scenario, the physics of EWBG can be in principle tested by exploring the Higgs sector.

It has been well known that the 1stOPT is realized by the non-decoupling thermal loop effects on the finite temperature effective potential and/or by the field mixing of the Higgs boson with additional scalar fields [4–26]. These effects also affect the effective potential at the zero temperature, so that they normally deviate the triple Higgs boson coupling (the hhh coupling) typically by larger than 10% [6,10–12,27,16,17,24,25]. It may be challenging for the (high luminosity) LHC to achieve this level of accuracy. However, the plan of the International Linear Collider (ILC) [28] includes the determination of the hhh coupling with 10% accuracy by upgrading the center-of-mass energy to $\sqrt{s} = 1$ TeV [29–31]. The Compact LInear Collider (CLIC) [32] also aims to reach the similar accuracy. The Future Circular Collider of electrons and positrons (FCC-ee) [33] will not address the precision measurement of the hhh coupling as its center-of-mass energy is insufficient. The possibility of testing the hhh coupling at future hadron colliders with $\sqrt{s} = 100$ TeV is also considered [34]. Therefore, the scenario of EWBG can be tested by precision measurements of the hhh coupling at future collider experiments. In a class of models where the

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1stOPT is caused by the field mixing, resulting predicted values for the Higgs boson couplings such as those with weak gauge bosons and with the SM fermions can also be deviated significantly from their SM values because of the field mixing. Therefore, this class of models for EWBG is expected to be tested by the data from the LHC not only those at future linear colliders.

On the other hand, it has also been known that strongly 1stOPT at the early Universe is a discriminative origin of gravitational waves (GWs) [35–37,24,38,39,25,40–42]. Recently, the GWs has been directly detected at the Advanced LIGO experiment which has an astronomical origin [43]. By this discovery, measurements of GWs with various frequencies will be accelerated in the near future including KAGRA [44], Advanced LIGO [45] and Advanced VIRGO [46], by which new field of GW astronomy will be extensively developed. Furthermore, future space based GW interferometers such as eLISA [47], DECIGO [48] and BBO [49] provide us an opportunity of measuring GWs with a wider range of frequencies, which can cover GWs from the first order EWPT. Therefore, by precisely measuring the spectrum of GWs, we can test the physics of EWPT and further the scenario of EWBG.

In this Letter, we calculate the spectrum of GWs originated from strongly first order EWPT in a concrete renormalizable model, the extended Higgs model with a real singlet field. In order to calculate the bubble nucleation rate, we perform a two-field analysis to evaluate bounce solutions connecting the true and the false vacua using the one-loop effective potential at finite temperatures. We survey allowed regions of parameters of the model imposing the Sakharov condition of the departure from thermal equilibrium for baryogenesis [50]. We then investigate the GWs produced at electroweak bubble collisions in the early Universe, such as the sound wave, the bubble wall collision and the plasma turbulence. We find that in this model strongly first order EWPT can be well tested by the combination of the precision study of various Higgs boson couplings at the LHC, the measurement of the *hhh* coupling at future lepton colliders and the spectrum of GWs detectable at eLISA and DECIGO.

Let us begin with a brief review of the Higgs-singlet model (HSM), which is one of the simplest extensions of the SM [9,10,51,52,18,53,54,26,55]. The Higgs sector of the HSM is equipped with a real isospin scalar singlet S in addition to the Higgs doublet Φ. The general tree-level Higgs potential allowed by gauge invariance and renormalizability is given by

$$V_0 = -\mu_\Phi^2|\Phi|^2 + \lambda_\Phi|\Phi|^4 + \mu_{\Phi S}|\Phi|^2 S + \frac{\lambda_{\Phi S}}{2}|\Phi|^2 S^2 + \mu_S^3 S + \frac{m_S^2}{2}S^2 + \frac{\mu_S'}{3}S^3 + \frac{\lambda_S}{4}S^4, \tag{1}$$

with eight parameters $\mu_\Phi^2, m_S^2, \lambda_\Phi, \lambda_S, \lambda_{\Phi S}, \mu_{\Phi S}, \mu_S'$ and μ_S^3.[1] After the condensation of the two Higgs fields, they are expanded around the vacuum expectation values v_Φ and v_S as

$$\Phi = \begin{pmatrix} G^+ \\ \frac{1}{\sqrt{2}}(v_\Phi + \phi_1 + iG^0) \end{pmatrix}, \quad S = v_S + \phi_2. \tag{2}$$

There appear two physical degrees of freedom ϕ_1 and ϕ_2 that mix with each other in addition to Nambu–Goldstone (NG) modes $G^\pm$ and G^0 that are absorbed by the W- and Z-bosons. In the following, we analyze the phase structure of this HSM in the classical field space spanned by

$$\langle\Phi\rangle = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}\varphi_\Phi \end{pmatrix}, \quad \langle S\rangle = \varphi_S. \tag{3}$$

Radiative corrections modify the shape of the Higgs potential from the tree-level form. At zero temperature, the effective potential up to the one-loop level is [56]

$$V_{\mathrm{eff},T=0}(\varphi_\Phi,\varphi_S) = V_0(\varphi_\Phi,\varphi_S) + \sum_i n_i \frac{M_i^4(\varphi_\Phi,\varphi_S)}{64\pi^2}\left(\ln\frac{M_i^2(\varphi_\Phi,\varphi_S)}{Q^2} - c_i\right), \tag{4}$$

where Q is the renormalization scale, which is set at v_Φ in our analysis. Here, n_i and $M_i(\varphi_\Phi,\varphi_S)$ denote the degrees of the freedom and the field-dependent masses for particles i, respectively. We take the $\overline{\mathrm{MS}}$ scheme, where the numerical constants c_i are set at 3/2 (5/6) for scalars and fermions (gauge bosons). We impose the tadpole conditions using the one-loop level effective potential as $\langle\partial V_{\mathrm{eff},T=0}/\partial\varphi_i\rangle = 0$, with $i = \Phi$ or S. Here, the angle bracket $\langle\cdots\rangle$ represents the field-dependent quantity evaluated at our true vacuum $(\varphi_\Phi,\varphi_S) = (v_\Phi, v_S)$. The mass squared matrix of the real scalar bosons in the (ϕ_1,ϕ_2) basis is diagonalized as

$$m_{ij}^2 = \left\langle\frac{\partial^2 V_{\mathrm{eff},T=0}}{\partial\varphi_i\partial\varphi_j}\right\rangle = \begin{pmatrix}\cos\theta & -\sin\theta \\ \sin\theta & \cos\theta\end{pmatrix}\begin{pmatrix} m_h^2 & 0 \\ 0 & m_H^2\end{pmatrix}\begin{pmatrix}\cos\theta & \sin\theta \\ -\sin\theta & \cos\theta\end{pmatrix}, \tag{5}$$

leading to one-loop improved mass eigenvalues of the Higgs bosons m_h and m_H, and mixing angle θ ($m_h < m_H$, $-\pi/4 \le \theta \le \pi/4$). The lighter boson h is identified with the discovered Higgs boson with the mass 125GeV in this Letter, and the alternative case where H is the discovered one will be examined elsewhere. From the above equations, we use v_Φ, v_S, m_h, m_H and θ as the input parameters instead of μ_Φ^2, m_S^2, λ_Φ, $\mu_{\Phi S}$ and μ_S'.

Due to finite temperature effects, the effective potential is modified to [57]

$$V_{\mathrm{eff},T}[M_i^2(\varphi_\Phi,\varphi_S)] = V_{\mathrm{eff},T=0}(\varphi_\Phi,\varphi_S) + \sum_i n_i \frac{T^4}{2\pi^2} I_{B,F}\left(\frac{M_i^2(\varphi_\Phi,\varphi_S)}{T^2}\right), \tag{6}$$

where

$$I_{B,F}(a^2) = \int_0^\infty dx\, x^2 \ln\left(1 \mp \exp^{-\sqrt{x^2+a^2}}\right), \tag{7}$$

for boson and fermions, respectively. In order to take ring-diagram contributions into account, we replace the field-dependent masses in the effective potential as [58]

$$M_i^2(\varphi_\Phi,\varphi_S) \to M_i^2(\varphi_\Phi,\varphi_S,T) = M_i^2(\varphi_\Phi,\varphi_S) + \Pi_i(T), \tag{8}$$

where $\Pi_i(T)$ stand for the finite temperature contributions to the self energies. We consider loop contributions from the fields $i = h, G^\pm, G^0, H, W^\pm_{T,L}, Z_{T,L}, \gamma_{T,L}, t$ and b. As for the scalar sector particles, the thermally corrected field-dependent masses are given by [51]

$$M_{h,H}^2(\varphi_\Phi,\varphi_S,T) = \frac{1}{2}\left(M_{11}^2 + M_{22}^2 \mp \sqrt{(M_{11}^2 - M_{22}^2)^2 + 4M_{12}^2 M_{21}^2}\right), \tag{9}$$

$$M_{G^0,G^\pm}^2(\varphi_\Phi,\varphi_S,T) = -\mu_\Phi^2 + \lambda_\Phi\varphi_\Phi^2 + \mu_{\Phi S}\varphi_S + \frac{\lambda_{\Phi S}}{2}\varphi_S^2 + \frac{T^2}{48}(9g^2 + 3g'^2 + 12(y_t^2 + y_b^2) + 24\lambda_\Phi + 2\lambda_{\Phi S}), \tag{10}$$

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[1] One of the mass parameters can be removed by the field redefinition of the singlet field without loss of generality [52,53].

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