New Insights of Electroweak Phase Transition in NMSSM (Postprint)

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Full Text

Preamble

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New Insights of Electroweak Phase Transition in NMSSM

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Abstract: We perform a detailed semi-analytical analysis of the electroweak phase transition (EWPT) property in NMSSM, which serves as a good benchmark model in which the 126 GeV Higgs mixes with a singlet. In this case, a strongly first-order electroweak phase transition (SFOEWPT) is achieved by the tree-level effects and the phase transition strength c is determined by the vacua energy gap at T = 0. We make an anatomy of the energy gap at both tree-level and loop-level and extract out a dimensionless phase transition parameter R 4 vs/A , which can replace A in the parameterization and affect the light CP odd and even Higgs spectra. We find that SFOEWPT only occurs in R −1 and positive R (cid:46) O(10), which in the non-PQ limit case would prefer either a relatively light CP odd or CP even Higgs boson (60, 100) GeV, therefore serves as a smoking gun signal and requires new search strategies at the LHC.

1 Introduction and Motivations

In the last two years, the ATLAS and CMS collaborations have established the discovery of the long-expected standard model (SM)-like Higgs boson h, with a significance up to 6.1 and 6.9 , respectively [?]. This new resonance has a relatively light mass mh (cid:39) 126 GeV, and its observed production or decay rate is close to the SM one. With more data accumulating, we would enter into the territory of precise understanding of the electroweak symmetry breaking (EWSB) mechanism. In an orthogonal direction, one may wonder about its impacts on weak-scale cosmology, in particular the corresponding thermal property: the nature of the electroweak phase transition (EWPT).

This is not only a big question of early cosmology per se, but it also helps us understand the origin of baryon asymmetry in the sense that a strongly first-order EWPT (SFOEWPT) is required for successful electroweak baryogenesis (EWBG). Baryogenesis has a close relation with Higgs physics, and moreover successful baryogenesis implies a non-standard Higgs boson (for discussions on CP violation, see Ref. [?]). Broadly speaking, with the current LHC data on Higgs production and decay, we can specify three classes of SFOEWPT models based on their discovery potential through Higgs physics.

The first class is that there is a colored or electrically charged particle which couples to the 126 GeV Higgs boson. In this case, new particles which alter the Higgs production or decay through gluon fusion or di-photon decay channel will change the Higgs effective potential and potentially enhance the EWPT strength [?, ?, ?, ?, ?, ?, ?, ?, ?]. Comprehensive studies have been carried out after the LHC data, and it is found that for a single particle, a SFOEWPT requires an enhanced gluon fusion production rate and suppressed Higgs di-photon decay width [?, ?]. This problem can be cured if one introduces another particle with its loop contributions opposite to the first one while the EWPT strength is enhanced [?].

The second class is that we have a singlet scalar which couples to the Higgs but never develops a VEV [?]. In this case, future precision electroweak and Higgs measurements would constrain the overall kinematical renormalization of the 126 GeV Higgs induced by this model. The last class is that the extra scalar gets a VEV and mixes with the Higgs (or through a tadpole term which is essentially the same [?]) or there are multi-Higgs [?]. In this case, it is the mixing effect that changes the Higgs physics properties. Investigating its genuine features clearly is an important task.

Supersymmetry (SUSY) is a well-motivated example among beyond-SM models and it can also provide the SFOEWPT for successful EWBG. For instance, in the minimal supersymmetric SM (MSSM) the significant Higgs-stop coupling can lead to an acceptable EWPT strength in a tiny window, given a well-organized stop sector [?]. In light of the recent LHC Higgs discovery and stop exclusion, this window has been severely constrained and essentially ruled out [?, ?, ?] (for a remedy, see [?]). In this class of model, one challenge after the Higgs discovery is to lift mh with the least fine-tuning while still accommodating the Higgs constraints. One simple extension is the NMSSM, which provides a large tree-level Higgs mass and a natural solution to the µ problem. With an extra singlet in the Higgs sector, it is conceivable that SFOEWPT is still viable in NMSSM, and we are curious about the phase transition patterns constrained by the current data.

In this article, we have studied this problem in great detail and found a critical parameter R 4 vs/A where SFOEWPT only occurs in R −1 and positive R (cid:46) O(10), which in turn would prefer a lighter CP odd or even Higgs boson. This paper is organized as follows. In Section 2 we review the NMSSM in detail, including both its Higgs potential at zero and finite temperature. In Section 3, we first analyze the SFOEWPT in NMSSM semi-analytically through both the tree-level and loop effects, and then provide the numerical results of the parameter scan which includes all the current experimental and strong electroweak phase transition conditions. We also show the corresponding particle spectra patterns, LHC observations and dark matter in Section 4. Finally we conclude and give a discussion in Section 5, and some necessary details in the paper are given in the Appendices.

2 The NMSSM at Zero and Non-zero Temperature

As mentioned in the introduction, the NMSSM can accommodate natural SUSY with the current data constraints and provides viable dark matter candidates, thus receiving a lot of attention. It also provides a good benchmark model where the 125 GeV Higgs mixes with a singlet getting a VEV, therefore providing tree-level cubic terms in the Higgs effective potential to enhance the strength of EWPT. In this section we review the basic formulas for the Higgs effective potential setup without and with finite temperature corrections.

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2.1 Tree-level Higgs Potential

All of the above eminent features of the NMSSM are traced back to the Higgs sector, which in the Z3-invariant form is written as

WZ3 SHu · Hd + (cid:18) -Lsoft = A Hu · HdS + (cid:19) A S3 + h.c. (2.1)

After S gets a VEV vs (cid:104)S(cid:105) around the electroweak (EW) scale, an effective µ-term is generated. This is the original motivation of the singlet extension. But as a great bonus, the model provides a Higgs quartic term at tree level, i.e., 2|H0d|2, which can be significant for a large 1 and moreover a small tan 1. As a consequence, the tree-level mass of the SM-like Higgs boson becomes:

mh^2 = (cid:0)mZ^2 cos^2 2 + ^{2v}2 sin^2 2 (cid:1) + mmix^2, (2.2)

where mmix^2 stands for the mixing effects on the Higgs boson mass. It can be sizable (typically a few GeV), positive or negative depending on the mass ordering of the SM-like Higgs boson among the neutral Higgs bosons [?, ?]. If hSM is the lightest one (H1-scenario), the effect is a reduction. In contrast, if hSM is the next lightest one (H2-scenario), the effect is an enhancement. After the LHC Higgs 126 GeV discovery, the spectra patterns of those two scenarios have been studied intensively [?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?, ?]. Here we will investigate the status of EWPT in these two scenarios separately.

For later convenience, we give the complete tree-level Higgs potential, which consists of the D-, F- and the soft SUSY breaking terms:

V0 = (cid:12)(cid:12) Hu · Hd - S$^{2(cid:12)(cid:12)}$2 + mHu^2H†uHu + mHd^2H†dHd + mS$^{2|S|}$2 + (cid:0) A Hu · HdS + A S^3 + h.c.(cid:1), (2.3)

where Hu (H+u, H0u) and g^2 = (cid:0)g2^2 + g1^2(cid:1)/2. Here we will not discuss the CP violation aspects of electroweak baryogenesis and assume , A , , A R for simplicity. An SU(2)×U(1) gauge is chosen such that at the physical vacuum

v+ (cid:104)H+u(cid:105) = 0, vu (cid:104)H0u(cid:105) R+, v- (cid:104)H-d(cid:105) = 0, vd (cid:104)H0d(cid:105) R+. (2.4)

This is a local minimum provided that the charged Higgs bosons have positive mass squared. Moreover, A > 0 and > 0 are assumed to realize vd, vS R+.

The angle is defined as in the MSSM:

vd = v cos , vu = v sin , (2.5)

with v (cid:39) 174 GeV.

For the Higgs mass squared structure, we decompose the Higgs fields as follows [?]:

Hu = vu + (S1 cos + S2 sin) + i(P1 cos + G0 sin), Hd = vd + (-S1 sin + S2 cos) + i(P1 sin - G0 cos), S = vs + S3 + iP2, (2.6)

where G0 is the Goldstone boson. In this basis, the doublet block has already been approximately diagonalized and S2 is the SM-like which carries electroweak VEV among the doublets. In the basis (S1, S2, S3), the elements of the CP-even Higgs mass squared matrix (M^2_S)ij are given by

(M^2_S)11 = M_A^2 + (mZ^2 - ^{2v}2) sin^2 2 , (M^2_S)12 = -(mZ^2 - ^{2v}2) sin 4 , (M^2_S)13 = -(M_A^2 sin 2 + 2 v^2s) cos 2 , (M^2_S)22 = mZ^2 cos^2 2 + ^{2v}2 sin^2 2 , (M^2_S)23 = (-M_A^2 sin^2 2 + 4 ^{2v}2) (cid:18)vs/v(cid:19) + 4 ^{2v}2s - v^2 sin 2 , (M^2_S)33 = A sin 2 (cid:18)vs/v(cid:19) + 4 ^{2v}2s + A vs - v^2 sin 2 . (2.7)

where M_A^2 = 2 vs(A + vs)/sin 2 defines the largest scale among these elements and is the heavy CP-odd Higgs mass. We can introduce an auxiliary parameter CA 1 - A sin 2 /2µ - sin 2 / to measure the mixing between singlet and doublet, i.e., (M^2_S)13 = 2CA µv (The mixing (M^2_S)12 can be safely neglected for moderately large tan). The other light Higgs diagonal mass square (M^2_S)33 can be written as

(M^2_S)33 = 4 ^{2v}2s (cid:18)1 + 1/R (cid:19) + ^{2v}2 - v^2 sin 2 , (2.8)

and we will use this formula again and again in later discussions. Here R 4 vs/A is a dimensionless critical variable defined for SFOEWPT.

For the CP-odd Higgs boson, A is theoretically upper bounded in order to keep the CP-odd singlet-like Higgs mass squared (M^2_P)22 positive [**?**]:

(M^2_P)11 = M_A^2, (M^2_P)22 = A sin 2 (cid:18)vs/v(cid:19) + 3 A vs, (M^2_P)12 = (A + 4 vs)v - v^2 sin 2 , (2.9)

where (M^2_P)11 corresponds to the mass squared M_A^2 of the only CP-odd Higgs in the MSSM.

2.2 Effective Potential at Finite Temperature

The starting point for the perturbative analysis of EWPT is the finite temperature effective potential. Up to one-loop order, it takes the form

V(l, T) = V0(l) + V1(l, T) + Vdaisy(l, T). (2.10)

where l, l = d, u, s are the classical field variables corresponding to H0d, H0u, and S. The tree-level part V0 follows directly from the Higgs potential in Eq. (2.3). We realize that our analysis at this precision may be subject to corrections from higher orders and the issue of gauge dependence [**?**, **?**], and a more complete analysis is left to a future study.

The one-loop part V1 consists of the Coleman-Weinberg potential at zero temperature and thermal corrections at finite temperature [**?**]:

V1 = (cid:88)i (-1)^(2si) (ni/64 ^2) m_i^4(l) [ln(m_i2(l)/Q2) - ci] + (cid:88)i (-1)^(2si) (niT4/2 2) J_i(m_i2(l)/T2), (2.11)

where i runs over all particles in the model, with each having degrees of freedom n_i, field-dependent mass m_i(l) and spin s_i. J_i is the thermal integral function J_B(F) for bosons (fermions)

J_B,F(y^2) = (cid:90)∞0 dx x^2 ln(1 e(-√(x2+y^2))). (2.12)

It tends to be zero in the non-relativistic limit, i.e., y^2 (cid:29) 1. By contrast, it has the high-temperature expansion and in particular gives rise to the well-known thermal cubic term in Eq. (2.11), given that i is a boson. Here we work in the Landau gauge and in the DR scheme.

As for Vdaisy, it is the daisy resummation contributions from the longitudinal components of gauge bosons and the scalar bosons [?, ?, ?]

Vdaisy = - (T/12) (cid:88)b [m_b^3(l, T) - m_b^3(l)], (2.13)

where m_b is the thermal mass.

Finally, it should be emphasized that in our analysis, we use the three VEVs v_l as the inputs and eliminate the Higgs soft masses via the minimization conditions for the three field variables l. Concretely, at one-loop order, they are given by

V1(T=0)/ u| l=vl = (A + vS)vS tan - 2(vS2 + v^2 sin^2) - mHu^2 = 0, V1(T=0)/ d| l=vl = (A + vS)vS cot - 2(vS2 + v^2 cos^2) + mZ^2 cos 2 = 0, V1(T=0)/ S| l=vl = A sin 2 - A vS - 2v2 - 2 2vS2 + v^2 sin 2 = 0. (2.14)

3 Electroweak Phase Transition in the NMSSM

With previous preparations, in this section we study EWPT in this model. It is well known that successful EWBG requires a SFOEWPT, namely c vc(Tc)/Tc (cid:38) 0.9. For a dedicated study of this condition based on gauge invariant quantities, see Ref. [?]. Here Tc is the critical temperature of SFOEWPT, with order parameter vc(Tc). In the SM, lattice simulation indicates that its EWPT is actually a crossover, failing to achieve any jumps in terms of order parameter. In the MSSM, in particular after the discovery of the 126 GeV Higgs boson, a single light stop alone would be ruled out by the current Higgs data because of too large enhancement on the Higgs production rate from gluon fusion. Nevertheless, a second colored light scalar can not only reduce the gluon-Higgs effective operator, but also enhance the EWPT strength [?]. While the NMSSM, by virtue of its tree-level effects, provides a simple way to enhance c. Such effects have been studied by many groups before [?, ?, ?, ?, ?, ?, ?, ?], but a detailed general analysis of SFOEWPT after the Higgs discovery is still absent and we fill the gap in this paper.

We will first introduce three types of EWPT and then propose a new way to investigate c from the zero-temperature Higgs effective potential. Following this approach, we make an anatomy of each type, giving semi-analytical treatment of tree-level effects and qualitative analysis of loop corrections. It is found that the latter plays a robust role in SFOEWPT, despite the dominated tree-level effects.

3.1 Vacua Energy Gap and SFOEWPT

The NMSSM contains three Higgs fields and thus possesses a rich vacua structure, which leads to a variety of EWPT patterns. There are mainly three patterns [**?**, **?**], classified by the course of the phase transition from the symmetric phase Ω0 to the EW symmetry breaking phase ΩEW (here we denote various phases with their VEVs):

Type-I: Ω0 ΩS ΩEW. At high temperature, the universe is in the symmetric phase. As the universe cools down, it may transit to the vacuum located in the singlet subspace, i.e., ΩS. As T further decreases to the critical temperature Tc, ΩS degenerates with ΩEW and then the universe transits into the phase ΩEW. Notice that if S = 0 in ΩEW this would induce a symmetry non-restoration effect in the S direction [**?**].

Type-II: Ω0 ΩU ΩEW. Type-II transition passes through the intermediate phase ΩU with Hu developing a VEV first. Here only the first step is relevant for EWPT, which recovers the SM case except that S contributes to the thermal cubic terms [**?**]. However, in this case, SFOEWPT requires large interactions between Higgs and singlet which induces sizable mixing between the two. This will change the transition type into Type-III. Generally, Type-II is hardly strong so we will not discuss this type.

Type-III: Ω0 ΩEW. The EW symmetry breaking vacuum develops first, and thus the universe in the symmetric phase transits directly into the phase ΩEW. It is worth pointing out that even though the transition does not undergo other phases, there are still extra local minima at T = 0. In particular, there usually exists a local minimum in the singlet subspace which makes the vacua structure Type-I-like. We will return to this point in later discussions.

Vacua structure at T = 0 should encode information on EWPT. For instance, the effective potential in Type-I is likely to have a metastable vacuum ΩS besides the EW vacuum ΩEW, with vacua energy gap ΔV VS - VEW. The T-dependent terms in the finite temperature potential need to smooth out this gap as T increases, until the critical temperature. Accordingly, a smaller ΔV (T = 0) may imply a lower Tc thus a larger c. This conjecture is confirmed by our final numerical results shown in Fig. 1 [Figure 1: see original paper]. In the three-dimensional field space, developing an analytical expression for c is mission impossible, except for some simplified cases like in the PQ-limit [**?**, **?**]. Therefore, our observation provides an important guideline for achieving a larger c.

This general correlation between the phase transition strength and the vacua energy gap ΔV can be understood in the $\Delta V \to 0$ limit. Let' s consider the vacua energy gap between the symmetry phase and the broken phase:

ΔV = (Vsym - VEW(v0))|T=0 (cid:39) Vsym(Tc) - VEW(T=0, v0) = VEW(Tc, vc) - VEW(T=0, v0) (cid:39) Tc VEW/ T|T=Tr (3.1)

where in the second line, we have used the approximation Vsym(Tc) (cid:39) Vsym(T=0) in the $\Delta V \to 0$ limit, which holds exactly if the symmetric phase is the origin of or is a good approximation if the symmetric vacuum is a weakly coupled singlet. The third line comes from the degenerate vacua condition at the critical temperature. In the last line, we set vc v0 and use the Lagrange mean value theorem with Tr [0, Tc]. Thus,

c vc/Tc (cid:39) (v0/ΔV) VEW/ T|T=Tc (3.2)

The dependence of the phase transition strength on ΔV can be revealed more explicitly in the simplified model of NMSSM in the PQ-limit [**?**]. In this model,

ΔV (2ã^{2mS}2 + ^{2v}2)^2 / (8 ^2) (3.3)

For an extremely small ΔV, vc v, and the phase transition strength can be rewritten as:

c (E/) (2ã^{2mS}2 + ^{2v}2) / v^2 (cid:39) (E/) (ΔV)^(1/2) / v^2 (3.4)

where E is the coefficient of the thermal cubic term ^3T.

3.2 Anatomy of Type-I

In this subsection we will lead the way to generate a smaller energy gap ΔV (T = 0) in Type-I. A full understanding requires analysis at both tree-level and loop-level. We also give numerical results, which are consistent with those semi-analytical understandings.

3.2.1 Tree-level Analysis First, let' s investigate the vacuum energy VEW of ΩEW. Substituting Eq. (2.14) into Eq. (2.3), one can eliminate the Higgs soft masses in VEW, and then VEW can be divided into three parts: V^H_EW, V^S_EW, and V^HS_EW. The first part denotes the contribution completely from the Higgs doublets:

V^H_EW = - (1/8)(g^{2v}2 cos^2 2 + ^{2v}2 sin^2 2) (cid:39) -v$^{2m—h}$2/4 (3.5)

To derive the second approximation we have used Eq. (2.2) where we neglect the mixing effects for the Higgs boson mass. As one can see, V^H_EW is definitely negative. Moreover, its value is determined by the Higgs quartic coupling and thus related to the SM-like Higgs boson mass, mh (cid:39) 126 GeV. Therefore, this part is almost fixed to be around -1.18 $\times$ 10^8 GeV^4.

The second part V^S_EW is the contribution from the singlet sector, taking the form:

V^S_EW = - (1/3) A vS^3 - ^{2vS}4 (3.6)

The third part V^HS_EW is a result of the doublet-singlet mixing, and it can be cast into a simple form:

V^HS_EW = -CAμ^{2v}2. (3.7)

As we have mentioned before, CA measures the mixing between singlet and doublet: (M^2_S)13 = 2CA μv. From the current Higgs data, we expect this auxiliary parameter CA is usually much smaller than 1 [**?**], so the singlet or the SM-like Higgs mass is not largely pushed down in the H2- or H1-scenario respectively. This fact will help us to simplify discussions and furthermore find out a crucial variable R which has a close relation with the vacua energy gap ΔV and the EWPT strength c.

Next we discuss VS, the tree-level potential energy of the absolute minimum uS in the singlet subspace. To compare with VEW, it is convenient to eliminate mS^2 through the third equation of Eq. (2.14), rewritten as:

mS^2 = -CA ^{2v}2 - A vS - 2 ^{2vS}2 (3.8)

Then from the potential with only singlet S:

V(S) = mS^{2S}2 + A S^3 + ^{2S}4, (3.9)

we can get:

VS = (cid:2)- A (vS - 2uS/3) - $^{2(2vS}$2 - uS^2)(cid:3) uS^2 - CA ^{2uS}2 (3.10)

It is also illustrative to express VS in terms of the inputs only:

VS = - (A $^{4/384}$ 2)(1 + √(1 - 8x))^2 (1 + √(1 - 8x) - 12x), (3.11)

which holds for x   mS$^{2\ /A}$ 2 < 1/8, see more details in Appendix A. VS is definitely negative for x < 1/9. Moreover, it is an even function of both A and . For a given x , Eq. (3.11) indicates that VS becomes more negative as A ^2/ (or A / to some extent) increases.

With all the above expressions of potential energy, we proceed to discuss the vacua energy gap at tree level, which is found to be related to the deviation of uS from vS. To see it, consider the small deviation case and write uS = (1 +)vS (| | (cid:28) 1), then the gap is approximated as:

ΔVtree = VS - V^S_EW - V^HS_EW - V^H_EW (cid:39) -CA ^{2v}2(uS^2 - vS^2) + $^{2(2uS}$2(uS - vS) + vS(vS^2 - uS^2)) (3.12)   4 ^{2} 2(1 + 1/R)vS^4 - 2 CAv^{2μ}2 + v$^{2m—h}$2/4 (3.13)

Obviously, ΔVtree goes to the doublet limit v$^{2m—h}$2/4 as   → 0. In other words, a substantial deviation is necessary to decrease the energy gap away from the doublet limit. In fact, ΔVtree can even be negative (we will see this soon). A negative ΔVtree is somewhat welcome since loop corrections will be found to favor uplifting VS relative to VEW.

One could have a closer inspection into the deviation. Consider the minima structure of the singlet subspace at tree level, whose details are listed in Appendix A. Its absolute minimum locates at the origin or

uS = -(A /2)(1 + √(1 - 8x)), (3.14)

In Type-I the latter is just the case, which requires x < 1/9. Using Eq. (3.8) one can rewrite x as:

x = (1 + R)^2/4 - CA ^{2v}2/A ^2 (3.15)

In the limit CA → 0, one gets the following simple relation between uS and vS:

uS { -vS(1 + 2/R) + O(CA) if R (cid:38) -1; vS + O(CA) if R (cid:46) -1. } (3.16)

which shows that uS usually deviates from vS significantly in the first case while in the second case they should be close to each other, given suppressed corrections from nonzero CA. Note that at leading order O(CA) is given by -CA ^{2v}2/|1 + R | A , which indicates that the approximation breaks down for R near -1. In this case, a positive CA in x can also generate a deviation.

Arguably, a substantial VEV deviation, i.e., for the first case in Eq. (3.16), tends to drive ΔVtree < 0. Notice that in the decoupling limit CA → 0, S = vS is always either a minimum or maximum (A < 0 and R (cid:38) -1) in the singlet subspace since the first derivative of V(S) from Eq. (3.9) over S is zero. In the latter case, it is not surprising that V^S_EW(vS) > VS(uS); in the former case, a large negative VS(uS) is also possible for uS (cid:46) -vS (R > 0). Recall that the singlet-doublet mixing term is suppressed by small CA, thus the above difference tends to dominate in ΔVtree, rendering it negative. This is particularly true in the A < 0 region when VS(vS) is a maximum, where R > -1 requires -A / > 4vS (cid:38) O(TeV) or even order of magnitude larger for a larger µ. That large -A / , in terms of the naive argument below Eq. (3.11), renders ΩS well below ΩEW. Therefore substantial loop corrections are indispensable to flip the order.

Before heading towards the loop corrections, let' s make some observations of the tree-level results on the R - ΔVtree plane (see the upper panel of Fig. 2 [Figure 2: see original paper]). They are in accord with the analysis above: (I) In R (cid:28) -1 region, uS (cid:39) vS, so ΔVtree clearly takes the doublet limit v$^{2m–h}$2/4; (II) In R -1 region, the magnitude of ΔVtree can blow up, in particular within the window -1 < R < 0 and for a relatively large µ; (III) In R > 0 region, as argued before, ΔVtree can also be negative and of order of a few 10^8 GeV^4, significantly smaller than case (II). More complementary analysis is left to the part of numerical study.

3.2.2 Loop-level Analysis Previously it was shown at tree level that ΩS usually lies above ΩEW. Here we will demonstrate that the tree-level order is going to be flipped by loop effects, which tend to lower ΩEW but lift up ΩS. In the DR scheme, the former is mainly ascribed to the remnant of Coleman-Weinberg potential after correcting the Higgs soft mass terms, while the latter is mainly due to the shift of uS. In the following we describe their details respectively.

On one hand, loop effects can lower ΩEW. To offset the shift of VEVs in ΩEW due to V1(T=0)/ l, with Eq. (2.11), we may need to add the corresponding Higgs quadratic term, -½Δm_l^2 l^2, with Δm_l^2 determined to be

Δm_l^2 = 2 (cid:88)i A_i(vl) [m_i^2(l)]' (L_i(vl) - 1). (3.17)

We have introduced A_i   (-1)$^{(2si)/64}$ 2 and L_i   ln(m_i$^{2(vl)/Q}$2) for short. As a consequence, the remnant of Coleman-Weinberg potential, V^R_CW = V_CW - ½Δm_l$^{2\,l}$2, results in a shift to the tree-level VEW vacuum energy:

ΔV^R_CW = (cid:88)i A_i(vl) m_i^4(vl) [L_i(vl) - ½(L_i(vl) - 1)]. (3.18)

The above expression can be simplified greatly for two limits of m_i^2(vl):

Strong v_l-dependence: In this limit the mass of particle i dominantly originates from coupling to Higgs fields, such as the SM particles and Higgsinos. Then we have

ΔV^R_CW = - (cid:88)i A_i(vl) m_i^4(vl) (L_i(vl) - 3/2). (3.19)

Bear in mind that we have fixed Q = 2 TeV, thus the relatively light fermions, e.g., top quark and Higgsinos, contribute a positive V^R_CW and make for flipping. The resulting decrease in VEW can be up to order 10^9 GeV^4 for a heavy µ   500 GeV. By contrast, the light weak gauge bosons hamper flipping but numerically it is unimportant due to their lightness.

Weak v_l-dependence: Some particles like stops have large (soft) mass terms, so they typically have quite weak dependence on v_l. In this limit one may write m_i^2(vl) = m_i^2(1 + f(x)) with x = v_l/m_i and f(x) (cid:28) 1. With that, we get an approximation

ΔV^R_CW   A_i m_i^4 [L_i - 3/2] - A_i m_i^4 [(1 - L_i)(2f(x) - xf' (x)) - L_i f(x)(f(x) - xf' (x))], (3.20)

where L_i   ln(m_i$^{2/Q}$2). The term in the first line is a constant thus contributing null to the energy gap. While for the second line, heavier CP even/odd Higgs bosons with their mass dependences on the VEVs (f(x)   x) will benefit the reduction of V^R_CW at the VEVs. If f(x) = _l^2 x^2 (stop without trilinear soft mixing), the leading x-dependence in the second row will vanish, with energy shift proportional only to powers of v_l, i.e., A_i L_i _l^4 v_l^4.

In summary, viewing from our particle spectrum, loop effects tend to decrease the energy of ΩEW. In the following we discuss the uS-shift effect on VS.

On the other hand, loop effects can lift VS up. Here the discussion is different from the previous case, because in ΩS the singlet VEV changes after loop corrections and the effective radiative potential plays an important role. It induces a shift to mS^2, inherited from the previous discussions in ΩEW. On top of that, it affects other tree-level couplings, as can be seen by expanding V_CW(S) into polynomials of S. We give corresponding typical examples:

The heavy Higgs bosons (still lighter than µ) and Higgsinos with mass s increase mS^2 and by an amount, respectively,

ΔmS^2 TeV$^{2/16}$ 2, $\rightarrow$ [1 + (1 - L) $^{2/8}$ 2]^(1/2). (3.21)

From Eq. (3.14) and Eq. (3.11) we know that both |uS| and |VS| are monotonically decreasing functions of x (and as well from Eq. (3.15) in CA $\rightarrow$ 0 limit). Thus, in the region with relatively small ^2 ((cid:46) 0.01), |uS| may be decreased and negative VS is increasing so ΩS is lifted up.

The loop-level numerical results on the R - ΔVtree plane are shown in the lower panel of Fig. 2 [Figure 2: see original paper]. The second role of lifting VS which effectively increases is crucial to flip vacua order with an especially large tree-level gap, which is characterized by -1 < R < 0. As mentioned before, it is usually accompanied with a relatively smaller and large µ, which yields a sizable increase of from Eq. (3.21). Recall that R ; increasing may drag R out of the window -1 < R < 0 and make R < -1. Therefore from Eq. (3.12), it is reasonable for us to draw a conclusion that the loop-level gap goes to the doublet limit hv^2/4 > 0 for those points. In this way, the tree-level order is flipped.

3.2.3 Numerical Results The SFOEWPT is the result of complicated interplay among quite a few parameters, including , , A , etc. We have turned to numerical methods for a global understanding, and for a cross-check with previous qualitative analysis. Using the NMSSMTools package [**?**, **?**, **?**], we scan the parameter space of the model with constraints from various relevant experiments, including the constraints on Higgs signatures. The parameter setting is listed as follows:

: (0.01, 0.5), : (0.3, 0.8), tan : (1.5, 10), A : (200, 2000) GeV, A : (-1000, 1000) GeV, µ: (100, 600) GeV. (3.22)

To minimize the uncertainty from the soft spectrum and explore the EWPT properties from the genuine Higgs-singlet sector, we assume that they are irrelevantly heavy. In particular, the parameters in the stop sector are taken to be m_Q^2 = 2 TeV and A_t = 0. The numerical points used in the previous sections are actually from this parameter space. We calculate the phase transition strength following the textbook approach: search for the minima of the complete one-loop Higgs potential (2.10) at each temperature, and then find out Tc and c by the degenerate vacua condition.

In the following we show the parameter distributions favored by SFOEWPT, and try to give interpretations of them.

Figure 3 [Figure 3: see original paper]. R versus c in Type-I transition, with color code denoting µ. Left panel: H1-scenario; Right panel: H2-scenario.

Firstly, we present plots on the R - c plane in Fig. 3. From them one can see that R is a helpful variable to judge c. In other words, the previous tree-level analysis, despite very complicated loop corrections, still provides valid insights. A possible large c is accommodated in two distinct regions:

R O(1) - O(10): c > 1 in Type-I favors the R > 0 (thus A / > 0) region, which could have a sizable (uS - vS) deviation to avoid the doublet limit. One may wonder why large R ' s fail. Consider R (cid:29) 1 and still take CA → 0; from Eq. (3.16) one gets uS -vS. Then it is straightforward to derive

ΔVtree m_h^{2v}2/4 - 16 ^{2vS}4/3R + O(CA). (3.23)

So ΔVtree, with its singlet contribution suppressed by the large R , goes to the doublet limit again. Numerical results show that in H2-scenario, c > 1 occurs for R (cid:46) 10, while in H1-scenario c > 1 can still occur for R as large as 30.

Note that R cannot be too small either, owing to phenomenological reasons. A large positive A threatens the positivity of the light CP-odd Higgs boson mass (we will return to this later). Moreover, a very small leads to a light (singlino-like) neutralino, which opens too large 126 GeV Higgs exotic decay branching ratio in the non-PQ limit. Besides, a singlino-like LSP may be over-abundant because of too small annihilation cross section. In conclusion, A / cannot be too large and accordingly R gets a lower bound, about 4.0 and 2.0 in H1-scenario and H2-scenario respectively, as shown in Fig. 3.

R -1: For R < 0, it is not surprising that points with c > 1 crowd around R -1. On one hand, R (cid:28) -1 fails to achieve SFOEWPT owing to the doublet limit, as argued at the end of Section 3.2.1. On the other hand, R cuts off before approaching zero because -A is not allowed to be very large here owing to the singlet-like CP-even Higgs boson. Therefore, R is preferred to be around -1.

H1-scenario and H2-scenario demonstrate a remarkable distinction in this region, namely c (cid:38) 1 is accommodated in the latter but not in the former. If the point is in the H1-scenario with c (cid:38) 1, a large M33 term is required and from Eq. (2.8), we can see that this further requires a large MA (very large µ and A , see Fig. 3 and Fig. 6 [Figure 6: see original paper]) since the term 4 ^{2vS}2(1 + 1/R) is close to its minimal.

Notice that those points in H1-scenario have obviously large µ' s, which yield considerable loop corrections according to the discussion at the end of Section 3.2.2. Those loop corrections would bring back to zero (see Fig. 4 [Figure 4: see original paper]) and therefore forbid the SFOEWPT. As a comparison, in H2-scenario could have a very large deviation from zero and even approach > 1. It is precisely those points which have a small ΔV and trigger a SFOEWPT.

Figure 4. Loop-level gap ΔVnum versus (uS - vS)/vS for the Type-I transition for the H1- and H2-scenario respectively.

Next, we outline the distributions of relevant parameters favored by c > 1 in Type-I:

- Light µ and small tan are preferred (see Fig. 5 [Figure 5: see original paper]), in particular in H1-scenario where µ (cid:46) 250 GeV and tan (cid:46) 3.5. In H2-scenario they can extend to a bit larger regions, i.e., µ (cid:46) 250 GeV and tan (cid:46) 5. In this sense, Type-I agrees with the most natural NMSSM scenario [?], which has a slice of parameter space with 1, tan 1 and µ mZ.
- A (200, 600) GeV (see Fig. 6 [Figure 6: see original paper]). Due to the suppressed mixing effect |CA| (cid:28) 1, A is strongly correlated with µ, i.e., A 2µ/ sin 2 .
- As for and A , most of the preferences can be traced back to the discussion on R = 4 vS/A . We emphasize again that both H1- and H2-scenario accommodate SFOEWPT for A > 0, while for A < 0, c can hardly achieve O(1) in H1-scenario (see Fig. 8 [Figure 8: see original paper]).

3.3 Type III: Results and Analysis

In this subsection, we turn our attention to the Type-III transition. This type of EWPT is studied in detail by early works [?, ?, ?, ?] due to its compatibility with the near-PQ symmetry limit and its one-step nature. Here we revisit this type of transition in the spirit of energy gap. Most of the analysis is similar to that of Type-I.

It is worth pointing out that Type-III arises not only in the case that the origin is indeed the absolute minimum in the singlet subspace but also in the case that the origin is metastable. The latter has a Type-I-like vacua structure but belongs to Type-III due to thermal evolution: At T = 0, the absolute minimum in the subspace locates at uS 0, but as temperature increases it will exceed the origin and become energetically disfavored. So the degeneracy eventually happens between the origin and ΩEW, which determine the gap. Such type-crossing phenomena are not difficult to understand. The critical case is when mS^2 is positive and small (typically mS (cid:46) 100 GeV), x = $mS^{2\ /A}$ 2 (cid:28) 1 means that mS^2 is quite sensitive to temperature. Although it is unable to distinguish quantitatively this case from Type-I at T = 0, our numerical results tell that the models of this case have a smaller gap between the origin and ΩS than those in Type-I.

The energy gap is quite simple, given by:

ΔVtree = -V^H_EW - V^HS_EW - V^S_EW (cid:39) v^{2m_h}2/4 + $CA\mu^{2v}$2 + ^{2vS}4(4/3 + R) (3.24)

Obviously, the singlet part dominates the energy gap for a large µ, and thus a small gap appeals to a large negative A such that -4/3 (cid:46) R < 0. For a moderate µ, the mixing part becomes important and a negative CA can help to decrease the gap, which requires a large A . The features outlined above are well reflected in Fig. 9 [Figure 9: see original paper].

The above analysis can be well adjusted in the near PQ-limit, where the mixing part from the second term in Eq. (3.24) could play a dominant role, so a large A and a moderate µ play key roles in decreasing the energy gap. Notice that in Ref. [**?**] where a relatively larger tan > 10 region is considered, the corresponding A (2, 5) TeV is even larger in order to achieve a small ΔV for SFOEWPT. We observe that Type-III can only be accommodated in a very restricted region. This is because that as soon as R becomes a bit larger, ΩS will become so deep that the transition changes into Type-I.

The loop-level analysis in Section 3.2.2 is also applicable to Type-III, i.e., there are mainly two kinds of loop-level effects. On the one hand, loop corrections lower ΩEW and make it to be the global minimum in the same way as in Type-I. On the other hand, the loop corrections in the singlet subspace lift up ΩS by increasing mS^2 and . However, the latter has no influence on the energy gap concerned here, and just makes ΩS shallow and ready to exceed the origin in thermal evolution.

Figure 9 [Figure 9: see original paper]. R versus c in Type-III transition, with color code denoting µ. Left panel: H1-scenario; Right panel: H2-scenario.

Finally, we summarize the parameter preference of SFOEWPT in Type-III:

- Unlike Type-I, SFOEWPT in Type-III prefers a larger µ (cid:38) 250 GeV, especially in the region -4/3 (cid:46) R < 0 (see Fig. 9). As argued above, ΔV here is effectively decreasing by the singlet part. Similar to Type-I, tan (cid:46) 3.5 is favored (see Fig. 10 [Figure 10: see original paper]).
- The favored values of A here are larger than in Type-I since A (cid:39) 2µ/ sin 2 or even larger for a negative CA (see Fig. 11 [Figure 11: see original paper]).
- Another obvious distinction between Type-I and -III can be observed in Fig. 13 [Figure 13: see original paper]: A > 0 barely accommodates Type-III, not to mention realizing c > 1. Eq. (3.8) may provide a simple interpretation. It indicates that given a positive A we have mS^2 < 0 and consequently x < 0, which yields a large deep minimum in the singlet direction, therefore strongly favoring Type-I.
- The vast majority of surviving points are in H2-scenario. As a matter of fact, H1-scenario seemingly has a strong tension with Type-III. This may be blamed on its parameter configurations that make it difficult to achieve a large M33 unless µ and A are large, which forbids the SFOEWPT just like the case of Type-I.

4.1 Higgs Spectra

The most remarkable connection may lie in the existence of a relatively light Higgs boson in the spectrum, though their masses are heavier than mhSM/2; otherwise the Higgs exotic decay channel would dominate and get severely constrained from the current Higgs signal rates.

For the CP-odd Higgses, it is well known that the absence of tachyon states in the Higgs CP-odd sector strongly favors negative A , and now it is more constrained from the observed Higgs signal rate as mentioned above. Such a sharp contradiction substantially compresses the allowed parameter space and leads to a remarkable prediction on the Higgs spectrum. The CP-odd Higgs sector contains two physical states a1,2. In the case of M_A^2 (cid:29) M_Z^2 we have

m_a1^2 (A + 4 vS)vuvd/vS - 3 A vS (4.1)

where the second term can be expressed as -12 ^{2vS}2/R . Thus, a positive small R drives the lighter CP-odd Higgs boson mass downwards.

For the CP-even Higgses, R can also affect the singlet-dominated CP-even Higgs mass, as we mentioned in Eq. (2.8):

(M^2_S)33 = 4 ^{2vS}2 (1 + 1/R) + ⋯ , (4.2)

where ⋯ stands for other terms that are either small or not relevant here. Therefore, the preference of R from SFOEWPT may result in some specific distribution in the Higgs spectrum. Especially we can see that -1 < R < 0 is the region where the above term gets its minimum and will drive down the h1 mass significantly. Therefore, it is clear that in the NMSSM, the SFOEWPT will impose a specific Higgs spectrum through the critical parameter R :

- R > 0 prefers a light CP-odd Higgs mass (small m_a1) with no strong preference on m_h1.
- 0 > R > -1 prefers a light CP-even Higgs mass (small m_h1 in the H2-scenario) with no preference on m_a1.

The above speculation is confirmed by the histograms Fig. 14 [Figure 14: see original paper] and 15, in which we have imposed the CMS Higgs signal data on surviving samples in various channels including hSM → ZZ, WW, , bb, at 2 level [?] (ATLAS constraints have a relatively smaller number of points with the same distributions). In Fig. 14, we compare the normal histograms for m_h1, m_a1 and those with SFOEWPT. We can see that the latter has more points concentrated in the light m_h1, m_a1 region. For those histograms where we have both R > 0 and R < 0, we further distinguish them in the histograms Fig. 15 [Figure 15: see original paper]. We can see that there is a very distinct pattern: R > 0 prefers small m_a1 while R < 0 prefers small m_h1, which is coincident with our analysis above.

4.2 Dark Matter Consideration

If the lightest NMSSM neutralino is assumed to be the WIMP Dark Matter (DM) candidate, an important consideration is the DM relic density in the present epoch. Combining the PLANCK [?] and WMAP 9-year data [?] and also including a 10% theoretical uncertainty, the 2 range of the WIMP DM relic density can be considered in the following range:

0.091 Ωhˆ2 0.138 (4.3)

Since the bino and wino mass parameters M1, M2 have been fixed at 2 TeV in our analysis, the lightest neutralino basically consists of higgsino and/or singlino. From the neutralino mass matrix one can tell that can significantly affect the singlino component in the DM:

M_neutralino = [M1 0 -g1vd/√2 g1vu/√2 0 M2 g2vd/√2 -g2vu/√2 -g1vd/√2 g2vd/√2 0 -µ g1vu/√2 -g2vu/√2 -µ 0] (4.4)

Since all squarks and sleptons have decoupled in our analysis, the DM can annihilate in the early universe only through the light Higgs bosons in the s-channel process or as the final states when the process is kinematically opened. For a highly higgsino-like DM, the coupling of the Higgs with the DM can be sizable and it is very easy to obtain a small relic density. We have checked that for each SFOEWPT scenario in our discussion, it is always possible to pick out several samples which can produce a relic density that does not overclose the universe, or even lies in the band shown in Eq. (4.3).

5 Conclusion and Discussion

After the discovery of the 126 GeV Higgs boson, the NMSSM is an attractive supersymmetric theory by virtue of its specific tree-level effect to enhance the Higgs boson mass and allowing a more natural µ parameter. On top of that, the tree-level effect can readily enhance the strength of EWPT c (cid:38) 1.0, which is required for a successful EWBG mechanism to generate the baryon asymmetry. In this article we have concentrated on studying SFOEWPT in the NMSSM, paying special attention to its relation with Higgs phenomenology. We have calculated the EWPT strength c with the one-loop finite temperature effective potential and found that a larger c requires a smaller gap ΔV. Then, in terms of the vacua structure and its evolution with temperature, we divide EWPT into three categories—Type-I, II, and III—along with two Higgs spectra patterns: H1-scenario and H2-scenario. We use our semi-analytical analysis for intuitive understanding and then use our numerical results to confirm those understandings.

We have observed a dimensionless critical parameter R 4 vS/A which demonstrates a clear correlation between the different types of SFOEWPT in NMSSM and the Higgs spectra as follows:

- In H1-scenario, the Type-I phase transition prefers R > 0 ((5, 30)) and very few points exist in the Type-III phase transition with small negative R .

- In H2-scenario, the Type-I phase transition has two distinct regions, in which either 0 > R > -1 or R > 0 ((2, 7)). For the Type-III phase transition, most of them lie in the region 0 > R > -4/3.

A SFOEWPT in general prefers a relatively light CP-odd or CP-even (H2-scenario) Higgs mass. In particular, R > 0 prefers small m_a1 while R < 0 prefers small m_h1.

The current classification of the EWPT patterns and the Higgs spectra has great importance to guide us to the next step of understanding the nature of EWPT. In particular, the Higgs spectra in SFOEWPT in the NMSSM, or even more broadly, in some SM/2HDM + singlet models, prefers either a light CP-odd or CP-even Higgs with mass slightly more than mhSM/2. We hope that our work can help bring more attention to the observation of those (60, 100) GeV light Higgses, and works along this direction will be presented in the future.

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A Minimum in the Singlet Subspace

For convenience, we present the minimum structure in the singlet subspace here. The potential has the form

V(S) = b2 S^2 + b3 S^3 + b4 S^4. (A.1)

The minimum condition V/ S = 0 has the following solutions:

S = 0, (A.2) S = -b3/(2b4) [1 ± √(1 - 8x)], (A.3)

where x b2b4/(2b3^2). The latter is physical as long as Δ b3^2 - 4b2b4 > 0, i.e., x 1/8.

The minimum structures are listed in Table 1 . Sometimes it is useful to express the potential at the nontrivial minimum as

Vmin(S±) = - (b3$^{4/64b4}$2) f(x), (A.4)

where f(x) = 4 + 4√(1 - 8x) + 96x^2/(1 + √(1 - 8x) - 48x - 32x^2), which decreases monotonically from one to zero as x increases from zero to 2/9. Obviously, Vmin = 0 for x > 2/9.

Table 1 . Minimum structures in singlet direction

x range	b2	b3	b4	Minimum
x > 1	>0	<0	>0	S=0

x range	b2	b3	b4	Minimum
x < 1	<0	>0	>0	S=S-
x < 1	<0	<0	>0	S=S+
x < 1	>0	>0	>0	S=0
x < 1	>0	<0	>0	S=0

References

[1] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012); S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012).

[2] J. Shu and Y. Zhang, Phys. Rev. Lett. 111, no. 9, 091801 (2013) [arXiv:1304.0773 [hep-ph]].

[3] D. J. H. Chung, A. J. Long and L. -T. Wang, Phys. Rev. D 87, 023509 (2013) [arXiv:1209.1819 [hep-ph]].

[4] W. Huang, J. Shu and Y. Zhang, JHEP 1303, 164 (2013) [arXiv:1210.0906 [hep-ph]].

[5] H. Davoudiasl, I. Lewis and E. Ponton, Phys. Rev. D 87, 093001 (2013) [arXiv:1211.3449 [hep-ph]].

[6] M. Laine, G. Nardini and K. Rummukainen, JCAP 1301, 011 (2013) [arXiv:1211.7344 [hep-ph]].

[7] R. Huo, arXiv:1305.1973 [hep-ph]; X. Chang and R. Huo, Phys. Rev. D 89, 036005 (2014) [arXiv:1402.4204 [hep-ph]].

[8] S. S. AbdusSalam and T. A. Chowdhury, arXiv:1310.8152 [hep-ph].

[9] M. Laine, G. Nardini and K. Rummukainen, arXiv:1311.4424 [hep-lat].

[10] A. Katz and M. Perelstein, arXiv:1401.1827 [hep-ph].

[11] T. Li and Y. -F. Zhou, arXiv:1402.3087 [hep-ph].

[12] B. Henning, X. Lu and H. Murayama, arXiv:1404.1058 [hep-ph].

[13] P. H. Damgaard, D. O' Connell, T. C. Petersen and A. Tranberg, Phys. Rev. Lett. 111, 221804 (2013) [arXiv:1305.4362 [hep-ph]].

[14] G. C. Dorsch, S. J. Huber and J. M. No, JHEP 1310, 029 (2013) [arXiv:1305.6610 [hep-ph]].

[15] M. S. Carena, M. Quiros and C. E. M. Wagner, Phys. Lett. B 380, 81 (1996) [hep-ph/9603420].

[16] M. Carena, G. Nardini, M. Quiros and C. E. M. Wagner, JHEP 1302, 001 (2013) [arXiv:1207.6330 [hep-ph]].

[17] T. Cohen, D. E. Morrissey and A. Pierce, Phys. Rev. D 86, 013009 (2012) [arXiv:1203.2924 [hep-ph]].

[18] D. Curtin, P. Jaiswal and P. Meade, JHEP 1208, 005 (2012) [arXiv:1203.2932 [hep-ph]].

[19] D. J. Miller, R. Nevzorov and P. M. Zerwas, Nucl. Phys. B 681, 3 (2004) [hep-ph/0304049].

[20] J. F. Gunion, Y. Jiang and S. Kraml, Phys. Lett. B 710, 454 (2012) [arXiv:1201.0982 [hep-ph]].

[21] S. F. King, M. Muhlleitner and R. Nevzorov, Nucl. Phys. B 860, 207 (2012) [arXiv:1201.2671 [hep-ph]].

[22] Z. Kang, J. Li and T. Li, JHEP 1211, 024 (2012) [arXiv:1201.5305 [hep-ph]].

[23] J. Cao et al., JHEP 1203, 086 (2012) [arXiv:1202.5821 [hep-ph]]; JHEP 1210, 079 (2012) [arXiv:1207.3698 [hep-ph]].

[24] K. S. Jeong, Y. Shoji and M. Yamaguchi, JHEP 1209, 007 (2012) [arXiv:1205.2486 [hep-ph]].

[25] J. F. Gunion, Y. Jiang and S. Kraml, Phys. Rev. D 86, 071702 (2012) [arXiv:1207.1545 [hep-ph]].

[26] K. J. Bae, K. Choi, E. J. Chun, S. H. Im, C. B. Park and C. S. Shin, JHEP 1211, 118 (2012) [arXiv:1208.2555 [hep-ph]].

[27] K. Agashe, Y. Cui and R. Franceschini, JHEP 1302, 031 (2013) [arXiv:1209.2115 [hep-ph]].

[28] Z. Kang, J. Li, T. Li, D. Liu and J. Shu, Phys. Rev. D 88, 015006 (2013) [arXiv:1301.0453 [hep-ph]].

[29] N. D. Christensen, T. Han, Z. Liu and S. Su, JHEP 1308, 019 (2013) [arXiv:1303.2113, arXiv:1303.2113 [hep-ph]].

[30] S. F. King, M. Muhlleitner, R. Nevzorov and K. Walz, Nucl. Phys. B 870 (2013) 323 [arXiv:1211.5074 [hep-ph]].

[31] T. Gherghetta, B. von Harling, A. D. Medina and M. A. Schmidt, JHEP 02, 032 (2013) [arXiv:1212.5243 [hep-ph]].

[32] R. Barbieri, D. Buttazzo, K. Kannike, F. Sala and A. Tesi, Phys. Rev. D 87, no. 11, 115018 (2013) [arXiv:1304.3670 [hep-ph]].

[33] M. Badziak, M. Olechowski and S. Pokorski, JHEP 1306, 043 (2013) [arXiv:1304.5437 [hep-ph]].

[34] S. Moretti, S. Munir and P. Poulose, Phys. Rev. D 89, 015022 (2014) [arXiv:1305.0166 [hep-ph]].

[35] T. Cheng and T. Li, Phys. Rev. D 88, 015031 (2013) [arXiv:1305.3214 [hep-ph]].

[36] C. Beskidt, W. de Boer and D. I. Kazakov, Phys. Lett. B 726, 758 (2013) [arXiv:1308.1333 [hep-ph]].

[37] J. Cao, F. Ding, C. Han, J. M. Yang and J. Zhu, JHEP 1311, 018 (2013) [arXiv:1309.4939 [hep-ph]].

[38] U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. 496, 1 (2010).

[39] H. H. Patel and M. J. Ramsey-Musolf, JHEP 1107, 029 (2011) [arXiv:1101.4665 [hep-ph]].

[40] C. L. Wainwright, S. Profumo and M. J. Ramsey-Musolf, Phys. Rev. D 86, 083537 (2012) [arXiv:1204.5464 [hep-ph]].

[41] M. Quiros, hep-ph/9901312.

[42] R. R. Parwani, Phys. Rev. D 45, 4695 (1992) [Erratum-ibid. D 48, 5965 (1993)] [hep-ph/9204216].

[43] M. E. Carrington, Phys. Rev. D 45, 2933 (1992).

[44] P. B. Arnold and O. Espinosa, Phys. Rev. D 47, 3546 (1993) [Erratum-ibid. D 50, 6662 (1994)] [hep-ph/9212235].

[45] M. Pietroni, Nucl. Phys. B 402 (1993) 27 [arXiv:hep-ph/9207227].

[46] A. T. Davies, C. D. Froggatt and R. G. Moorhouse, Phys. Lett. B 372, 88 (1996) [hep-ph/9603388].

[47] S. J. Huber and M. G. Schmidt, Nucl. Phys. B 606, 183 (2001) [hep-ph/0003122].

[48] A. Menon, D. E. Morrissey and C. E. M. Wagner, Phys. Rev. D 70, 035005 (2004) [hep-ph/0404184].

[49] K. Funakubo, S. Tao and F. Toyoda, Prog. Theor. Phys. 114, 369 (2005) [hep-ph/0501052].

[50] M. Carena, N. R. Shah and C. E. M. Wagner, Phys. Rev. D 85, 036003 (2012) [arXiv:1110.4378 [hep-ph]].

[51] C. Balazs, A. Mazumdar, E. Pukartas and G. White, arXiv:1309.5091 [hep-ph].

[52] S. Profumo, M. J. Ramsey-Musolf and G. Shaughnessy, JHEP 0708, 010 (2007) [arXiv:0705.2425 [hep-ph]].

[53] J. R. Espinosa, T. Konstandin and F. Riva, Nucl. Phys. B 854, 592 (2012) [arXiv:1107.5441 [hep-ph]].

[54] C. Cheung and Y. Zhang, JHEP 1309, 002 (2013) [arXiv:1306.4321 [hep-ph]].

[55] U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0502, 066 (2005) [hep-ph/0406215].

[56] U. Ellwanger and C. Hugonie, Comput. Phys. Commun. 175, 290 (2006) [hep-ph/0508022].

[57] U. Ellwanger and C. Hugonie, Comput. Phys. Commun. 177, 399 (2007) [hep-ph/0612134].

[58] https://twiki.cern.ch/twiki/bin/view/CMSPublic/PhysicsResultsHIG

[59] P. A. R. Ade et al. [Planck Collaboration], arXiv:1303.5076 [astro-ph.CO].

[60] G. Hinshaw et al. [WMAP Collaboration], Astrophys. J. Suppl. 208, 19 (2013).

[61] P. Draper, T. Liu, C. E. M. Wagner, L. -T. Wang and H. Zhang, Phys. Rev. Lett. 106, 121805 (2011) [arXiv:1009.3963 [hep-ph]].

[62] J. Huang, T. Liu, L. -T. Wang and F. Yu, arXiv:1309.6633 [hep-ph].

Note: Figure translations are in progress. See original paper for figures.