Electroweak Baryogenesis and Higgs Physics

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Abstract

We discuss the constraints on the Higgs sector coming from the requirement of the generation of the matter-antimatter asymmetry at the electroweak phase transition. These relate to both a strongly first-order transition, necessary for the preservation of the generated baryon asymmetry, and CP violation, necessary for its generation. This scenario may lead to exotic decays of the Standard Model like Higgs, a deviation of the di-Higgs production cross section, or CP violation in the Higgs sector. All these aspects are expected to be probed by the LHC as well as by electric dipole moment experiments, among others. Further phenomenological implications are discussed in this short review.

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1. INTRODUCTION

After the Higgs discovery at the LHC [1, 2], its properties including its mass and its couplings to the Standard Model (SM) particles have been studied in much detail [3, 4, 5, 6, 7]. Those studies have shown that the Higgs boson has properties similar to the ones expected in the Standard Model. However, only the couplings to the third-generation quarks and charged leptons, and weak gauge bosons are accurately known, with an uncertainty of the order of ten percent. In addition, we know very little about the Higgs potential. The SM potential is fully specified by its minimum and the Higgs mass, which defines the quartic coupling of the Higgs:

$$V_{\rm SM} = m_H^2 H^\dagger H + \frac{\lambda}{2}\left(H^\dagger H\right)^2. \quad (1)$$

This implies $v^2 = -\frac{m_H^2}{\lambda}$, with the neutral component ${\rm Re}[H^0] = \frac{1}{\sqrt{2}}(v+h)$, v is the Higgs vacuum expectation value (vev), and h is the normalized SM Higgs state, with mass $m_h^2 = \lambda v^2$.

Since the field h and the vacuum expectation value act in a coherent way on the fermion fields, the fermion couplings are both diagonal and proportional to the fermion masses, m_f, namely, $g_{f_i f_j h} = \frac{m_f}{v}\delta_{ij}$. The coupling to the massive gauge bosons, on the other hand, is proportional to their mass squared, namely, $g_{hhV} = \frac{m_V^2}{v}$. The tree-level Higgs sector is therefore CP-conserving, and all relevant CP-violating effects are in the charged gauge boson couplings to quarks, $g_{u_L^i d_L^j W^+} = \frac{g}{\sqrt{2}} V_{ij}^{\rm CKM}$, with $V^{\rm CKM}$ being the CKM Unitary matrix, which contains a phase, controlling CP violation, $\delta_{\rm CP}$. All the above properties are relevant to the question of baryogenesis (see, for example, [8, 9, 10, 11, 12]). Assuming CPT conservation, Sakharov defined the conditions for a successful baryogenesis scenario [13]:

(i) Baryon Number Violation,

(ii) C and CP Violation,

(iii) Non-equilibrium processes.

Baryon Number Violation is present in the SM and is induced by so-called Sphaleron processes [14], which are associated with the SM baryon and lepton number anomalous violating interaction, which at zero temperature are vanishingly small, but at high temperatures are suppressed by a Boltzmann suppression factor [15, 16]:

$$\Gamma_{\rm Sph} \propto T \exp\left(-E_{\rm Sph}/T\right). \quad (2)$$

Here, the Sphaleron energy $E_{\rm Sph} \sim 2\pi v B/g$ and $B \simeq 4$ [14], implying a Sphaleron energy of the order of 10 TeV at zero temperature.

As mentioned before, CP violation processes are present in the SM, but the corresponding processes are suppressed by Jarlskog invariant factors [17, 18], which, combined with the chiral suppression associated with the small first- and second-generation quark masses, makes it impossible to get the observed baryon number density:

$$\eta = \frac{n_B}{n_\gamma} \simeq 6 \times 10^{-10}, \quad (3)$$

where n_B and n_γ are the baryon and photon densities. Hence, new CP-violating sources beyond the SM ones are necessary for baryogenesis.

Finally, nonequilibrium processes occur if the electroweak phase transition is strongly first order. In such a case, bubbles of the real vacuum develop in the false, symmetry-preserving vacuum. If a given baryon number is generated at the electroweak phase transition at the critical temperature [8], after considering the proper Hubble expansion and the rate of baryon number violation [19], $\frac{n_B}{n_\gamma} \sim \frac{n_B}{n_\gamma}(T_c)\exp(-\frac{\Gamma_{\rm Sph}}{H}(T_c))$, where H is the Hubble expansion rate, $H \sim g_*^{1/2} T^2/M_{\rm Pl}$, and g_* are the number of relativistic degrees of freedom. For a transition temperature of the order of 100 GeV, this can only be fulfilled if

$$\frac{v(T_c)}{T_c} \gtrsim 1 \quad (4)$$

implying a strongly first-order electroweak phase transition (SFOEPT), which is schematically represented in Figure 1.

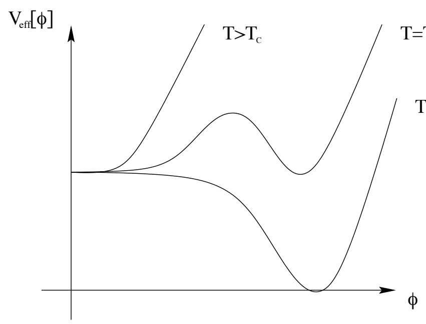

FIGURE 1: First-order phase transition. The Higgs vacuum expectation value $v(T)$ varies in a discontinous way at $T = T_c$, defined as the temperature at which the trivial and nontrivial minima are degenerate. From [9].

2. CONDITIONS FOR A SFOEPT

We will perform first an analysis of the Higgs potential at the one-loop level. In such a case,

$$V(h,T) = V_{\rm tree}(h,0) + \Delta V(h,0) + \Delta V(h,T). \tag{5}$$

Here, $\Delta V(h,0)$ describes the one-loop corrections at zero temperature, which include the renormalization of the masses and couplings and hence depends on the renormalization scale Q. The finite temperature corrections $\Delta V(h,T)$ are finite and differ for the case of bosons and fermions, due to their different statistics, and are given by [20]

$$\Delta V_{B,F}(h,T) = \pm g_{B,F} T^4 I_{B,F}\left(m_{B.F}(h)/T\right), \tag{6}$$

where $g_{B,F}$ is the number of degrees of freedom associated with the bosonic and fermionic fields (for instance, four for a Dirac fermion and one for a real scalar), $I_{B,F}(x)$ is an integral function, and $m_{B,F}(h)$ is the mass of the particle in the background of the Higgs field.

The integral functions may be expressed in terms of a high temperature expansion, which are valid for $m_{B,F}(h) \lesssim T$. Ignoring the field-independent terms, $I_{B,F}$ are given by

$$\frac{m_B(h)^2}{24T^2} - \frac{\left(m_B^2(h)^{3/2}\right)}{12\pi T^3} - \frac{m_B(h)^4}{64\pi^2 T^4}\ln\left[\frac{m_B(h)^2}{T^2}\right] + \ldots \tag{7}$$

and

$$\frac{m_F(h)^2}{48T^2} + \frac{m_F(h)^4}{64\pi^2 T^4}\ln\left[\frac{m_F(h)^2}{T^2}\right] + \ldots, \tag{8}$$

respectively. Observe the appearance of a nonanalytic dependence on $(m_B^2(h))^{3/2}$. In the case of boson fields with a linear dispersion relation on the Higgs field, this leads to a cubic term on h in the potential that is important for the generation of a first-order phase transition.

The logarithmic term is of the same form as the logarithmic corrections that appear in the renormalization of the Higgs potential in the $\overline{\rm MS}$ or $\overline{\rm DR}$ schemes [21]

$$\pm\frac{g_{B,F} m_{B,F}^4(h)}{64\pi^2}\left(\ln\left[\frac{m_{B,F}^2(h)}{Q^2}\right] + c\right), \tag{9}$$

where c is a constant. This implies that, at one loop, the logarithmic dependence on $m(h)^2$ cancels at high temperatures.

The Higgs and Goldstone masses become negative at low temperatures, leading to complex contributions to the potential. This can be fixed by a Daisy resummation of diagrams [20], which depends on the thermal mass $m_B^2(\phi, T)$:

$$m_B^2(h) + \Pi_B(T) = m_B^2(h) + a_B T^2, \tag{10}$$

where for any boson field a_B controls the temperature correction to its mass. After resummation, the bosonic contribution to the Higgs potential is corrected by

$$-T\sum_B g_B \frac{m_B(h,T)^3 - m_B(h,0)^3}{12\pi} \tag{11}$$

leading to a real potential at $T > T_c$. A systematic way of including higher loop effects was presented in [22]. An alternative resummation method, proposed by Parwani [23], is to replace $m_B^2(h)$ by $m_B^2(h,T)$ in all terms of the high-temperature effective potential. Away from the high-temperature regime, one should proceed with care. A self-consistent resummation, in which the Boltzmann suppression factors are included in a natural way, has been proposed to deal with this problem (see, for example, [24]).

Considering the high-temperature corrections, assuming that all relevant fields have a linear dispersion relation in h and absorbing the renormalization effects in to a redefinition of the parameters of the model, the Higgs potential then will have the form

$$V(h,T) = \frac{\left(m_H^2 + a_H T^2\right)}{2}h^2 - ETh^3 + \frac{\lambda(T)}{8}h^4, \tag{12}$$

where $E = \sum_B \frac{g_B y_B^3}{12\pi}$, the sum is performed on all bosons of mass $m_B^2 = y_B^2 h^2$, and we have redefined the potential in order to have zero value at $h = 0$. In the Standard Model, the most relevant fields are given by the transverse modes of the gauge bosons [25, 26]. The contribution to a_H, on the contrary, comes from all boson and fermion fields.

As shown in Figure 1, we shall define the critical temperature T_c as the temperature at which this new minimum is degenerate with the trivial one at $h = 0$. It is straightforward to show that this leads to

$$\frac{v\left(T_c\right)}{T_c} = \frac{4E}{\lambda\left(T_c\right)}. \tag{13}$$

Considering the mass of the W and Z gauge bosons given by $m_W^2(h) = g_2^2 h^2/4$ and $m_Z^2(h) = (g_2^2 + g_1^2)h^2/4$, respectively, and $g_V = 2$ for the transverse modes, we get

$$\frac{v\left(T_c\right)}{T_c} = \frac{2g_2^3 + \left(g_1^2 + g_2^2\right)^{3/2}}{12\pi\lambda\left(T_c\right)}. \tag{14}$$

The requirement that $v(T_c)/T_c \gtrsim 1$ can only be fulfilled for $\lambda(T_c) \lesssim 0.03$, which, since T_c is of the order of the weak scale, implies Higgs masses of the order or lower than 40 GeV, which is in conflict with the SM Higgs mass value. Moreover, a lattice gauge theory analysis shows that the line of first-order phase transitions ends at Higgs masses of the order of 70 GeV, and for $m_h = 125$ GeV, it becomes a crossover [27, 28].

3. NEW PHYSICS

The corrections to the Higgs potential may be induced, for instance, by

(i) An enhanced cubic term E in h,

(ii) A barrier that persists at zero temperature,

(iii) Fields that couple strongly to the Higgs (see, for example, [29]).

The first option can only occur via particles that are present in the plasma and, therefore, have masses of the order of the weak scale. A particular example, which has been studied in much detail, was the case of top squarks, which couple with couplings of order one to the Higgs and come in three colors and hence induce additional corrections to the Higgs potential [30, 31, 32, 33, 34, 35]. However, after resummation, a large cubic term can only be obtained if the field-independent mass of the stop is negative, implying a very light stop. Such a light-colored particle leads to strong contributions to the gluon fusion production of the Higgs boson that have not been observed experimentally [36, 37]. Beyond this example, one can postulate other particles to give a finite temperature cubic correction, but in general strong couplings and high multiplicities are demanded, which lead to strong constraints on these models [38]. In the following, we shall concentrate on the second possibility, which leads to nontrivial modifications of the Higgs properties.

3.1. *Singlet Extension of the SM*

The simplest extension of the SM is just to add a real singlet field. Even this simple extension allows for several new dimensionless and dimensional couplings. A simplification occurs in the case of a discrete Z_2 symmetry leading to the addition of the following potential terms:

$$V_s = \mu_s^2 s^2 + \frac{\lambda_{sh}}{2} s^2 h^2 + \frac{\lambda_s}{4} s^4, \tag{15}$$

while the Z_2 breaking terms include linear and cubic terms in the singlet field. An interesting case is the one in which the Z_2 symmetry is broken spontaneously. In the simplest approximation in which only the quadratic dependence on the temperature is considered, one obtains that the transition may occur from a minimum with zero doublet and nonzero singlet vacuum expectation value to one in which the doublet expectation value is nonvanishing, which develops at lower temperatures. The critical temperature is then defined when these two minima become degenerate. It is straightforward to show that [39]

$$\frac{v_c}{T_c} = \frac{4E}{\tilde{\lambda}} = \frac{4E}{\lambda_{\rm SM}} \left[1 + \sin^2\theta \left(\frac{m_h^2}{m_s^2} - 1 \right) \right], \quad \tilde{\lambda} = \lambda - \frac{\lambda_{sh}^2}{2\lambda_s}, \tag{16}$$

where θ is the scalar mixing angle. A strongly first-order phase transition may be obtained for a proper choice of λ_{sh} and λ_s. Such a transition demands a light singlet and a sizable mixing, which lead to a modification of the couplings of the Higgs to the SM fermions by a factor $\cos\theta$. Large values of θ are hence restricted due to precision electroweak and Higgs measurements [40]. This opens the possibility of the SM-like Higgs to decay into two light singlets [39]. On the other hand, as it is shown in Figure 2 in which the full one-loop finite T corrections are included, in general small values of $\lambda_{sh} \lesssim \mathcal{O}(0.1)$ and values of λ_s which are even smaller are required in this scenario.

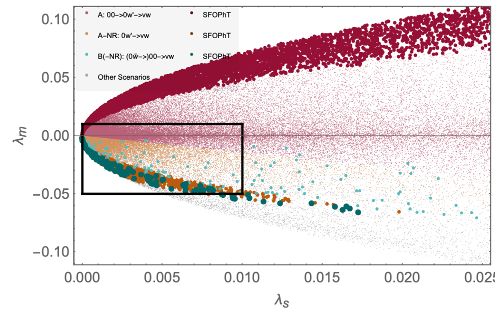

FIGURE 2: Parameters consistent with different transition patterns, where $\lambda_m \equiv \lambda_{sh}$. From [39].

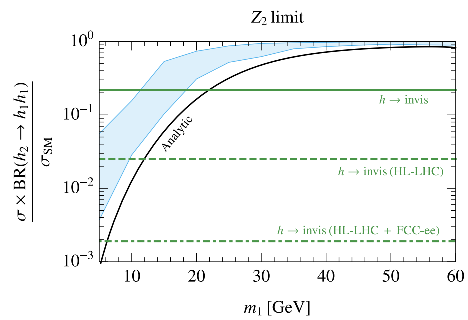

FIGURE 3: Branching ratio for the exotic decay of the SM Higgs going to two light singlets. From [41].

The Z_2 symmetry may remain preserved at zero temperature via a phase transition leading to a breakdown of the Z_2 symmetry at high temperatures, recovered at zero temperature. Certain conditions on the couplings must be fulfilled and, in particular, the coupling λ_{sh} should be sizable. The results are then less dependent on E but strongly dependent on the structure of the tree-level potential. In the case of a light singlet, this may again lead to exotic decays of the SM-like Higgs. The decay rate is given by [41]

$$\Gamma(h \to ss) \simeq \frac{\lambda_{hs}^2 v^2}{8\pi^2 m_h} \sqrt{1 - \frac{4m_s^2}{m_h^2}}. \tag{17}$$

Due to the vanishing scalar mixing, the singlet states do not couple to the SM particles and hence these decays contribute to invisible Higgs decays, which are currently being looked for at the LHC. As it is shown in Figure 3, due to the correlation of the decay branching ratio with the couplings governing the singlet mass, only light singlets, with a mass smaller than about 20 GeV, are allowed in this case. The prediction for current bounds on the branching ratio of the decay of an SM-like Higgs boson into two singlets, for a more general case of a SFOEPT in singlet extensions of the SM, is depicted in Figure 4.

One can also consider singlet extensions, including other symmetry breaking patterns and phenomenological properties, like Dark Matter, and we refer the reader to [43, 44, 45, 46, 47], for more details on these subjects.

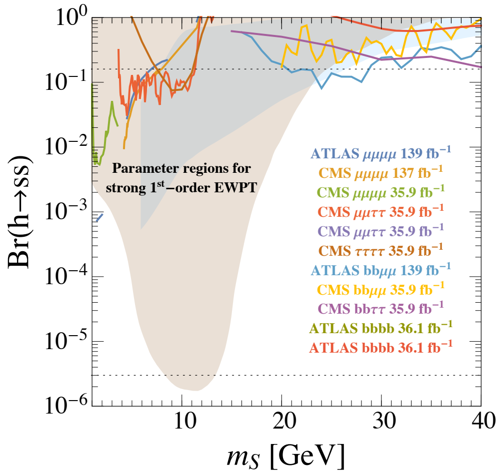

FIGURE 4: Region of parameters for a SFOEPT (shaded region) and the corresponding branching ratio of the decay of the SM-like Higgs into two singlets. Current experimental bounds are given. From [42].

3.2. *Heavy Singlet Field*

In this case, we will not demand a Z_2 symmetry, and we shall integrate out the singlet field. We shall include the potential terms

$$V(h,S) = \frac{m_H^2}{2}h^2 + \frac{\lambda}{8}h^4 + \frac{m_S^2}{2}s^2 + A_{hS}sh^2 + \frac{\lambda_{hS}}{2}h^2s^2. \quad (18)$$

The equation of motion of the field S is, at low momentum, given by their interaction with the Higgs field h [48, 49, 50]:

$$S = -\frac{A_{hS}h^2}{m_S^2 + \lambda_{Sh}h^2}. \quad (19)$$

Assuming that m_S is large compared to the weak scale, one can therefore obtain an effective potential by replacing s in the previous expression

$$V_{\text{eff}}(h) = \frac{m^2}{2}h^2 + \frac{\lambda}{8}h^4 - \frac{\left(A_{hS}h^2\right)^2}{2\left(m_S^2 + \lambda_{hS}h^2\right)}. \quad (20)$$

Expand $V_{\text{eff}}(h)$ in powers of $1/m_S^2$,

$$\frac{m^2}{2}h^2 + \left(\frac{\lambda}{8} - \frac{A_{hS}^2}{2m_S^2}\right)h^4 + \frac{A_{hS}^2\lambda_{hS}}{2m_S^4}h^6 + \ldots, \quad (21)$$

the result is then an effective potential in which the quartic term is modified and may become negative, and there are nonrenormalizable corrections which are parametrized by m_S^2. and the trilinear and quartic couplings A_{hS} and λ_{hS}. If one defines the correction to be proportional to $c_6h^6/(8\Lambda^2)$, however, this cutoff has a nontrivial dependence on these couplings and cannot be identified with the mass scales m_S. If we go to higher orders in $1/m_S^2$, one will obtain higher powers of h in the potential. We will provide a more general expression in terms of an effective theory in the next section.

3.2.1. *Higgs Potential at Higher Orders in h^2*

The simplest extension would be given by

$$V(\phi,T) = \frac{m^2 + a_0T^2}{2}h^2 + \frac{\lambda}{8}h^4 + \frac{c_6}{8\Lambda^2}h^6. \quad (22)$$

We do not include a cubic E term, assuming that its effects are small compared to the dominant temperature effects. This case has been studied in the literature in various contexts [38, 46, 50, 51, 52, 53, 54, 55, 56]. It is straightforward to show that the Higgs mass and the trilinear Higgs coupling are given by

$$\begin{aligned} m_h^2 &= \lambda v^2 + 3\frac{c_6v^4}{\Lambda^2}, \\ \lambda_3 &= \frac{3m_h^2}{v}\left(1 + \frac{2c_6v^4}{m_h^2\Lambda^2}\right). \end{aligned} \quad (23)$$

Therefore, there is a relevant modification of the trilinear Higgs coupling, proportional to c_6, which is one of the most relevant signatures of this scenario. We require $c_6 > 0$ for the stability of the potential. The requirement of a minimum of the potential at $h = v(T_c) = v_c$ degenerate with the extreme at $h = 0$ at $T = T_c$ leads to

$$v_c^2 = -\frac{\lambda\Lambda^2}{2c_6}, \quad (24)$$

what implies $\lambda < 0$.

Using equations (23) and (24), and the minimization relations, one obtains

$$\frac{c_6}{\Lambda^2} = \frac{m_h^2}{3v^2\left(v^2 - \frac{2}{3}v_c^2\right)}, \quad (25)$$

$$T_c^2 = \frac{3c_6}{4\Lambda^2 a_0}\left(v^2 - v_c^2\right)\left(v^2 - \frac{v_c^2}{3}\right). \quad (26)$$

Demanding both c_6 and T_c^2 to be positive, we get that trilinear coupling in the case of a first-order phase transition is bounded by

$$\frac{2}{3} \leq \delta \leq 2, \quad (27)$$

where $\delta = \frac{\lambda_3 - \lambda_3^{\text{SM}}}{\lambda_3^{\text{SM}}} = \kappa_\lambda - 1$.

Moreover, for $c_6 = 1$, we obtain a bound on the effective cutoff Λ, namely,

$$\frac{v^2}{m_h} < \Lambda < \frac{\sqrt{3}v^2}{m_h}, \quad (28)$$

which correspond to upper and lower bounds on Λ of approximately 500 GeV and 850 GeV, respectively, and as shown in equation (23), larger enhancement δ is obtained for the smaller values of the cutoff. The phase transition becomes stronger first-order for smaller values of the cutoff and becomes a weakly first-order one for values of Λ close to the upper bound in equation (28).

In Figure 5, we show the possible triple Higgs coupling enhancement factor δ as a function of the cutoff Λ for different extensions of the SM effective potential. The particular case of the potential of order h^6 is represented by the blue curve. The conditions to obtain a FOEPT up to order h^{10} are shown in this figure. The parameter space consistent with a SFOEPT is shown in Figure 6.

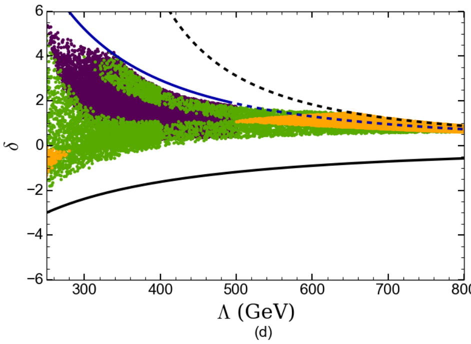

FIGURE 5: Triple Higgs coupling correction δ as a function of the cutoff Λ for the h^{10} case. The dashed dark blue shows the values consistent with a first-order electroweak phase transition FOEPT for the h^6 potential extension, for $c_6 = 1$. The different colors correspond to the different hierarchies of the effective potential h^{2n} coefficients in the case of a FOEPT. For other details, see [50].

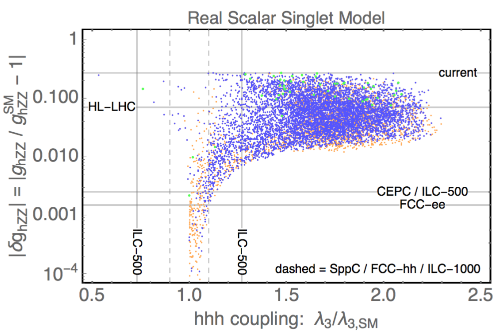

FIGURE 6: Modifications of λ_3 and of g_{hZZ} for the case of a (very) strongly first-order phase transition are shown by the (green) blue points. From [56].

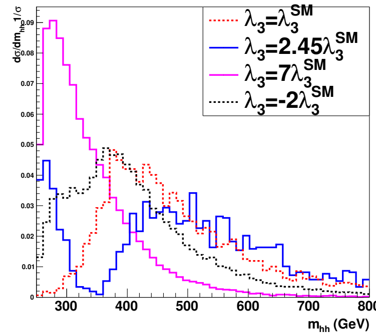

FIGURE 7: Normalized m_{hh} distributions for different values of λ_3. The cancellation between the box and triangle diagram is exact at $\lambda_3 = 2.45\lambda_3^{\mathrm{SM}}$ at $2m_t$ threshold, which explains the dip. Note that the distribution shifts to smaller values as $|\lambda_3|$ increases. From [71].

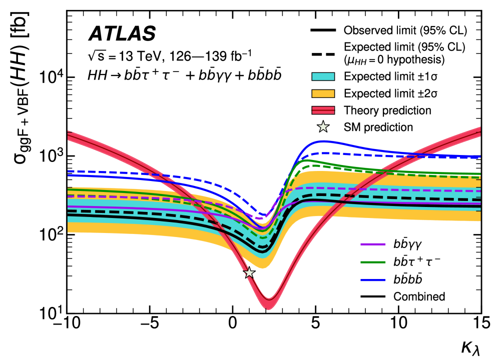

FIGURE 8: Current sensitivity in the determination of the triple Higgs coupling, assuming all other couplings to be fixed to the SM value. From [72].

3.3. *Experimental Tests of the Trilinear Higgs Couplings*

The Higgs trilinear coupling λ_3 can be probed by the double Higgs production at the LHC. At the leading order (LO), there are two diagrams contributing to the process. The triangle diagram is sensitive to λ_3 and the box diagram. The two diagrams interfere with each other destructively. The QCD NNLO differential cross sections are known [57, 58, 59, 60, 61, 62, 63, 64, 65, 66, 67]. For the Higgs decays, one can consider $\gamma\gamma$, W^+W^-, $\tau^+\tau^-$, and $b\bar{b}$ modes, which are measured in the single Higgs production at the LHC. As shown in Figure 7, the m_{hh} predicted range depends strongly on the coupling λ_3, and therefore, the precise experimental window must be fixed judiciously in order to increase the search sensitivity [50, 68, 69, 70].

The current sensitivity of di-Higgs production at the ATLAS experiment in different channels is shown in Figure 8. Due to the sensitivity of the di-Higgs production on the top Yukawa coupling (see, for instance, [71]), the determination of κ_λ has to be made by a fit to the single and di-Higgs production, but the allowed range of κ_λ is similar to the one obtained in the $\kappa_t = y_t/y_t^{\mathrm{SM}} = 1$ case [72].

The prospects for the high luminosity LHC run with $3000\,\mathrm{fb}^{-1}$ are excellent. Assuming an improvement of a factor 2 on both theoretical and systematic errors, something not unexpected, the sensitivity of the experiments for nonresonant Higgs production will be of about half of the SM cross section, implying a $3.4\,\sigma$ evidence for double Higgs production in the case of an SM cross section [73], and without taking into account a necessary combination of ATLAS and CMS analyses. Therefore, in combination with the single Higgs production cross section, which will strongly constrain the Higgs coupling to top and gauge bosons, the di-Higgs production will become an excellent probe of trilinear Higgs coupling modifications and therefore of electroweak baryogenesis.

Let us say, in addition, that resonant decays of non-standard Higgs bosons into pairs of SM-like Higgs bosons are also an important probe of baryogensis scenarios. We refer to [74, 75] for a further discussion of this subject.

3.4. *Extensions with Extra Higgs Doublets*

Two Higgs doublet models have also been studied in detail, and the parameter space consistent with a strongly first-order phase transition has been determined by both analytical and numerical methods; see, for example, [76, 77, 78]. A relevant parameter in these models is the ratio of the two Higgs vacuum expectation value, $\tan\beta$. Deviations of the SM-like Higgs boson couplings with respect to the SM value in the CP-conserving case are controlled by the misalignment of the direction of the Higgs expectation value, controlled by β, with the mixing in the CP-even sector, α. Values of $\cos(\beta - \alpha) \sim 0$ lead to the presence of a light Higgs with SM-like couplings [79, 80], and this condition tends to be fixed in most simulations. In the CP-violating case, additional mixing angles appear and the condition of alignment can be easily generalized to this case.

For nonstandard Higgs bosons with masses which are larger than the weak scale, preferred from the point of view of LHC searches, large quartic couplings are necessary in order to get a SFOEPT. These large couplings tend to induce a splitting between the CP-even and CP-odd Higgs bosons, which are assumed to be approximately good eigenstates, leading to the possibility of the decay of $A \to HZ$, which can be searched for at the LHC [81]. Recently, there has been also an intriguing claim of the possibility of more complicated phase transition patterns in the two Higgs doublet models, including the possibility of symmetry non-restoration, which would be very interesting and must be elucidated by further work [82].

The addition of doublets and singlets has also been studied, for instance, in the context of the Next To Minimal extension of the SM. This leads to a very complicated pattern of phases, which includes the possibility of a SFOEPT for particular values of the trilinear couplings A_λ of the singlet-Higgs fields $A_\lambda S H_u H_d$, even in the case of small quartic couplings [83, 84, 85, 86, 87]. The abundance of Higgs field directions implies the existence of several minima and transitions from one to the other may be obstructed by large barriers which demand a careful analysis of the transition rates in order to determine the possibility of baryogenesis [87]. Appropriate conditions may be found in regions of parameter space consistent with Dark Matter and an SM-like Higgs boson. Exotic decays of the heavy Higgs bosons into light singlets and SM Higgs and gauge boson states are expected in this case [88].

Further interesting scenarios include the possibility of inert doublets and composite Higgs models. We refer to [89, 90, 91] for further details on these subjects.

4. CP VIOLATION

4.1. Experimental Constraints

As stressed above, new CP-violating sources are necessary in order to generate the observed baryon asymmetry. These CP-violating phases tend to lead to contributions to the SM fermion electric dipole moments. These contributions depend on the interactions of these particles with the SM particles, but parametrically, for the electric dipole moments, they are governed by [11]

$$\begin{aligned} d_f &\sim \sin(\theta)\frac{\alpha}{4\pi}\frac{m_f}{M^2}, \\ d_f &\sim \sin(\theta)\frac{m_f}{\mathrm{MeV}}\left(\frac{1\,\mathrm{TeV}}{M}\right)^2 \times 10^{-26}\, e\cdot\mathrm{cm}, \end{aligned} \tag{29}$$

where θ is a new CP-violating phase, and M are the masses of the new particles. Observe that the constant of proportionality hides the details of the particle interactions. However, the comparison of this prediction with the current experimental bound on the electric dipole moment of the electron [92, 93]

$$d_e < 4.1 \times 10^{-30}\, e\cdot\mathrm{cm} \tag{30}$$

implies a strong bound on the new CP-violating sector.

The LHC can also put bounds on the CP-violating effects, by measuring the Higgs rates, which are affected by the inclusion of CP-violating couplings or directly by looking by CP-violating effects in the production and decay processes of the SM Higgs (see, for instance, [94, 95, 96, 97, 98, 99]).

4.2. CP-Violating Sources

The way the new CP-violating sources impact the baryon asymmetry is through the generation of particle CP-violating densities at the walls of bubbles of real vacuum state that expand at a certain velocity v_w. The particle densities diffuse into the bubbles while the left-handed baryons interact with the sphaleron states that provide the necessary baryon (and lepton) number violating processes. This is depicted in Figure 9. One should therefore solve the system of diffusion equations for the number densities or chemical potentials μ_i:

$$D_i \mu_i'' - v_w \mu_i + \Gamma_{ij}\mu_j = S_i, \tag{31}$$

where D_i is the diffusion constant, Γ_{ij} is the rate of interactions of the particle i with other particles j , the primes indicate derivatives with respect to the direction orthogonal to the bubble wall, and S_i are the CP-violating sources (see, for instance, [100]). Although the sphaleron processes in the symmetric phase should be included among the Γ_{ij}, they tend to be much slower than the other interactions $\Gamma_{\mathrm{Sph}} \sim \alpha_w^5 T$ (α_w is the weak coupling constant), and it is convenient to solve first for the CP-violating sources and finally solve for the baryon asymmetry while considering their sphaleron interactions. The final baryon asymmetry is then given by

$$n_B \sim -\frac{\Gamma_{\mathrm{Sph}}}{v_w}\int_{-\infty}^{0} n_{B_L}(z)\exp\left(zR/v_w\right), \tag{32}$$

where $R \propto \Gamma_{\mathrm{Sph}}$ is a relaxation coefficient and n_B is the final constant value in the broken phase. It is clear from this equation that the bubble wall velocity (and the bubble wall width) plays a relevant role in these predictions and that S_i must be properly computed in order to predict the correct baryon number. Both calculations are theoretically challenging (see, for instance, [85, 100, 101, 102, 103, 104, 105, 106, 107, 108, 109, 110] and [111, 112, 113]) and have been done in different approximations where for physics at the weak scale and phases of order one tend to predict values of the baryon asymmetry of the order of the observed values. Let us stress that there are large theoretical uncertainties in the computation of these quantities, and therefore, baryogenesis results are indicative of the region of parameters preferred in a given scenario but should not be taken as precise predictions of the models under analysis.

4.3. Modified Higgs Couplings

Since the mechanism of baryogenesis is connected to the electroweak phase transition, it is natural to expect that the Higgs may inherit CP-violating effects, including CP-violating couplings. If only the SM particles are in the high-temperature plasma at low energies, one can interpret the baryogenesis as being induced by the difference of the particle and antiparticle transmission and reflection coefficients at the bubble wall, induced by the CP-violating effects in the Higgs sector [114, 115, 116, 117, 118].

Following [118], one can define the Lagrangian

$$\mathcal{L} = y_f\left(1 + \frac{2}{v^2}\left(T_R^f + iT_I^f\right)H^\dagger H\right)\bar{f}_L f_R H + h.c., \tag{33}$$

where T_R^f and T_I^f characterize the Higgs fermion coupling modifications. In particular,

$$\frac{y_f}{y_f^{\mathrm{SM}}} = \frac{1}{\sqrt{\left(1+T_R^f\right)^2 + \left(T_I^f\right)^2}}. \tag{34}$$

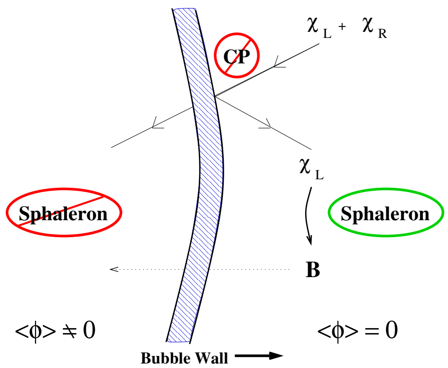

FIGURE 9: Particle transmission at the bubble wall. Inside the bubble, the sphaleron processes are suppressed. From [11].

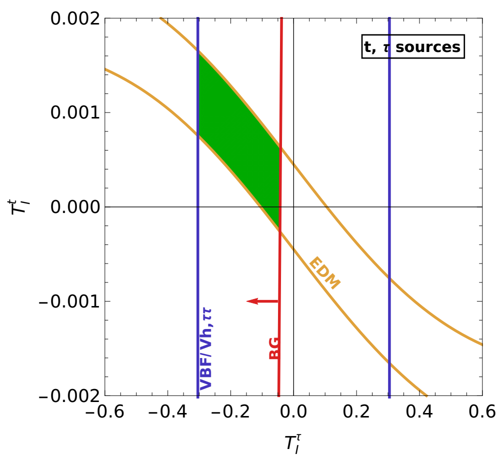

FIGURE 10: Constraints coming from electric dipole moments, Higgs precision measurements and the requirement of Baryogenesis. From [118].

The CP-violating sources depend on the alignment of the fermion mass with respect to its derivative with respect to the coordinate perpendicular to the bubble wall. The final result is given approximately by [118, 119]

$$Y_B = 8.6 \times 10^{-11} \times \left(51 T_I^t - 23 T_I^\tau - 0.44 T_I^b\right). \tag{35}$$

where $Y_B = n_B/s$ and s is the entropy density, $s \sim 7n_\gamma$. Hence, it is very important to determine the bounds on the imaginary component T_I^f. These bounds originate mainly from the electron dipole moment, which receives a contribution at two loops from the third-generation Higgs fermion couplings, $d_e \approx 1.1 \times 10^{-29}\, e\,$cm times approximately

$$2200 \left(\frac{y_t}{y_t^{\rm SM}}\right)^2 T_I^t + 10 \left(\frac{y_\tau}{y_\tau^{\rm SM}}\right)^2 T_I^\tau + 12 \left(\frac{y_b}{y_b^{\rm SM}}\right)^2 T_I^b. \tag{36}$$

It is clear from equations (35) and (36) that the tau coupling modification is the most likely to lead to a nonvanishing baryogenesis contribution. Single top or bottom Yukawa coupling effects can lead to only a fraction of the total baryon number but can add relevantly to the tau contribution. An example is given in Figure 10.

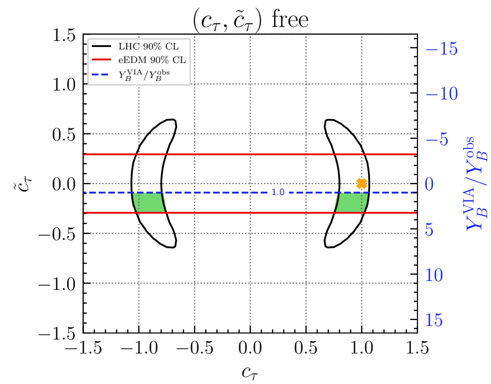

FIGURE 11: Constraints on the tau coupling from baryogenesis, LHC Higgs data and electric dipole moments. Here $c_\tau = T_R^\tau$ and $\tilde{c}_\tau = T_I^\tau$. From [94].

As it is shown in Figure 11, these coupling modifications are not yet constrained by the LHC [94] but may be probed in the near future at this collider [95, 96, 97, 98, 99].

4.4. CP Violation beyond the SM Higgs Sector

Let us stress that the above analysis is only valid if only the third-generation SM particles participate in the process of baryogenesis and in the contribution to electric dipole moments. This is generically not the case, and therefore, a more complete analysis is in order. In two Higgs doublet models [78, 120, 121, 122], for instance, the electric dipole moments receive contributions from both the standard and nonstandard Higgs bosons (for a complete analysis in the case of a Z_2 symmetry, see [123]). There are also specific correlations between the different CP-violating couplings that must be taken into account in different kinds of two Higgs doublet models, and cancellations between different contributions to the electric dipole moments [124, 125] may occur, which can lead to the possibility of baryogenesis in these models. However, as emphasized in [126], cancellation effects on the electric dipole moment of the electron may not exclude the constraints from the neutron or Hg electric dipole moments, and hence, the parameter space should be carefully studied to evaluate the possibility of electroweak baryogenesis.

In supersymmetric extensions, the CP-violating sources include relevant contributions from the electroweak fermion sector, as has been shown in many relevant works [85, 100, 101, 102, 103, 104, 105, 106, 107, 108, 109, 110], and this scenario is consistent with the inclusion of a Dark Matter candidate. The CP-violating currents originate from contributions coming from the relative phases of the Gaugino and Higgsino mass parameters and are proportional to derivatives of the Higgs fields, which imply nonstandard Higgs bosons at the reach of the LHC. Similar to the case of two Higgs doublet models, the electric dipole moments lead to strong constraints on these models, unless specific cancellations between different contributions occur [127, 128].

Let us add in closing that the strong constraints imposed by electric dipole moments have motivated the study of models in which CP violation occurs in a Dark sector (see [129, 130, 131]). These scenarios are associated with reduced, but still nonvanishing electric dipole moments that may be probed in future searches. A related idea is the possibility of breaking CP at temperatures close to the phase transition, enabling the baryogenesis processes, while recovering the CP symmetry at zero temperature, leading to no observable electric dipole moments. This idea has been explored, for instance, in [132, 133], and the conditions for baryogenesis were further studied in [132].

5. CONCLUSIONS

Baryogenesis at the electroweak phase transition remains an intriguing possibility and demands new physics at the weak scale. This scenario can be probed by precision Higgs physics, exotic Higgs decays and double Higgs production. Since the new physics inducing the modification should not decouple, one must expect to find new physics at the reach of the LHC, including a Dark Matter candidate, although its nature is not specified by the requirement of baryogenesis. In this short review, focused on Higgs physics, we have not discussed the production of gravitational waves, which is a relevant feature of a SFOEPT (see, for instance, [39, 56, 134, 135, 136, 137]). However, there tends to be tension between the bubble wall velocities associated with the very strong phase transition demanded for gravitational wave observations and the moderate ones necessary for proper diffusion of the baryon density in electroweak baryogenesis scenarios.

The requirement of new sources of CP violation which couple to the Higgs also leads to the expectation of nontrivial CP-violating couplings, inducing new collider signatures as well as electric dipole moments. Present constraints of these couplings imply that the tau-lepton modifications are the most likely to be associated with the generation of the baryon asymmetry. Let us stress, however, that these conclusions are based on a simplified scenario where only the third-generation Higgs anomalous couplings lead to the generation of the baryon asymmetry and to the electric dipole moment contributions and if there is new physics at the weak scale, some of these conclusions may be modified.

Overall, the coming years will lead to a further probe of the Higgs sector and of new physics at the weak scale and hence of this exciting scenario for the generation of the matter—antimatter asymmetry.

CONFLICTS OF INTEREST

The author declares that there are no conflicts of interest regarding the publication of this paper.

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