Degenerate states in the scalar boson spectrum
Is the Higgs Boson a Twin?

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The extension of the standard model to $SU(3)_L \times SU(3)_R \times SU(3)_C$ is considered. Spontaneous symmetry breaking requires two $(3^*,3,1)$ Higgs field multiplets with a strong hierarchical structure of their vacuum expectation values. An invariant potential is constructed to provide for these vacuum expectation values. This potential gives masses to all scalar fields apart from the 15 Goldstone bosons. In case there exists a one-to-one correspondence between the vacuum expectation values of the two field multiplets, the scalar boson spectrum contains degenerate eigenstates. The lowest eigenstate has a mass near 123 GeV close to the Higgs-like particle discovered at the LHC. In one class of solutions this lowest state is a nearly degenerate twin state. Each member is a superposition of fields from both multiplets with about equal strength. The twins are non identical twins, namely different combinations of a conventional Higgs and a Higgs field which is not coupled to fermions, only to gauge bosons. A second class of solutions leads again to degenerate states but in this case the state near 123 GeV remains a single state even for identical low scale vacuum expectation values in both multiplets.

I. INTRODUCTION AND SCOPE.

The author’s successful prediction [1] of the mass of the Higgs-like boson discovered at the LHC [2] may be fortuitous since it involved a speculation about potentials and their normalization. Also, in the literature, similar mass values in the correct energy range have been obtained using quite different speculations; see the compilation in [3]. Nevertheless, our own attempt may be worth to be discussed further. It differs strongly from others and needs little theoretical apparatus. The model is based on the extension of the Glashow-Weinberg-Salam group to $SU(3)_L \times SU(3)_R \times SU(3)_C$, the trinification group [4] [5], which is a subgroup of $E_6$ [6]. This group can only be unbroken at and above the scale where the gauge couplings $g_1$ and $g_2$ unite [7]. Thus, this model needs the presence of a very high mass scale $M$ of the order of $10^{13}$ GeV. This scale appears to be the right one for producing light Majorana neutrinos in the observed mass range. The appearance of such a high scale in the same particle representation of a symmetry group is challenging but not unwelcome. At low energies, the corresponding particle mixings still leave a trace in the limit $M \rightarrow \infty$. To break the trinification group down to the standard model group and finally to $U_e(1)$, two scalar matrix fields $H$ and $\tilde{H}$ transforming according to $(3^*,3)$ with respect to $SU(3)_L \times SU(3)_R$ are necessary [7] [8]. Thus, we have to deal with 36 real scalar fields. With respect to the $SU(2)_L$ subgroup they form 6 complex Higgs doublets and 6 complex singlets.

Our aim is to construct in a phenomenological way a potential for these scalar fields which provides spontaneous symmetry breaking and gives non-zero masses to all fields except for the 15 Goldstone bosons. Our input consists of the vacuum expectation values of these fields. Like in all present theories these fundamental quantities are not yet understood and thus have to be taken in accordance with known facts and restrictions. In this article we are particular interested in the existence of degenerate mass eigenstates. Some degeneracies are strictly valid due to to CP invariance. Others approximate ones arise from symmetry breaking parameters small compared to $M$. A particular interesting approximate degeneracy can occur if the vacuum expectation values (vevs) of the matrix field $H$ also appear in the vevs of $\tilde{H}$.

For the construction of the potential only vacuum expectation values will be used. It is taken to be of the Coleman-Weinberg type [8]. No explicit mass parameters, like in the tree potential of the standard model, are used except two tiny ones to account for the non-vanishing determinants of the two matrix fields $H$ and $\tilde{H}$. They break the discrete symmetry $H \rightarrow -H$ and $\tilde{H} \rightarrow -\tilde{H}$ and the symmetry under a pure phase transformation.

The vacuum expectation value of the first matrix ($H$) can be chosen diagonal. The vev of the field $H^1_1$, $v_1 = \langle H^1_1 \rangle$, is related to the top quark mass according to $m_t = g_t v_1$. $\langle H^2_2 \rangle$ is denoted by $b$ and connected to the bottom quark mass: $m_b = g_b b$. It is small compared to $v_1$. The element $\langle H^3_2 \rangle$, on the other hand, is huge and can be identified with $M$. The element $v_1$ is necessarily smaller or at most equal to $v = 174$ GeV, the vacuum expectation value of the mass: $m_t = g_t v_1$. With respect to the $SU(3)_C$ subgroup they form 6 complex Higgs doublets and 6 complex singlets.

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Higgs field of the standard model. Thus, the matrix field $H$ has Yukawa couplings to the top and to the $b$ quark and also to a very heavy down type quark $D$. By imbedding the scalar multiplet fields in the $E_6$ model (or its trinification subgroup) the vev’s $v_1$, $b$ and $M$ are independent of the generations. The generation matrices, which are vev’s of flavon fields, occur as factors in front of $H$ and $\tilde{H}$.

The second matrix of scalar fields $\tilde{H}$ has no Yukawa coupling to fermions (except via a heavy singlet state $\tilde{b}$). The vev’s of $\tilde{H}$ are unknown but have to obey important restrictions. This matrix field needs to have a sizeable vacuum expectation value of the element $\tilde{H}^2_2$ which we denote by $M_2$. This vev breaks the left-right symmetry and thus determines the masses of the right handed vector bosons. The experimental limit on right handed currents provides for a lower limit for $M_2$ of the order of 10 TeV. Already the presence of $M$ and $M_2$ is sufficient to break the trinification group down to the Glashow-Weinberg-Salam group ($\times U(1)$). Like $v_1$, the vev of $H^1_1$ denoted by $v_2$ is restricted by the known value of $v$. The vev’s of the second row of $\tilde{H}$, $b_2 = \langle \tilde{H}^2_2 \rangle$ and $b_3 = \langle \tilde{H}^3_2 \rangle$ are expected to be small like $b$. In fact, $v_1$, $v_2$ together with $b$, $b_2$, $b_3$ have to obey $v_1^2 + v_2^2 + b^2 + b_2^2 + b_3^2 = v^2$. For small $b$’s this relation reduces to $v_1^2 + v_2^2 = v^2 = (174 \text{ GeV})^2$. These vev’s break the Glashow-Weinberg-Salam group down to the electromagnetic $U_e(1)$ symmetry. The vev $\langle \tilde{H}^3_2 \rangle = M_3$ is unknown and not restricted by known facts. It can be large.

Defining $H = h + if$ and $\tilde{H} = \tilde{h} + if$ where $h, \tilde{h}, f, \tilde{f}$ are matrices of real fields we have in general

$$ \langle h \rangle = \begin{pmatrix} v_1 & 0 & 0 \\ 0 & b & 0 \\ 0 & 0 & M \end{pmatrix}, \quad \langle \tilde{h} \rangle = \begin{pmatrix} v_2 & 0 & 0 \\ 0 & b_2 & b_3 \\ 0 & M_2 & M_3 \end{pmatrix}. \tag{I.1} $$

Here the phases of the vev’s of $\langle H \rangle$ are absorbed by the transformation matrices and the fermion fields. $\langle \tilde{H} \rangle$ cannot be diagonalized anymore. We took all its vev’s to be real.

The potential will have its minimum at specific vacuum field configurations chosen from the general form $\langle H \rangle$. All first derivatives of this potential have to vanish at this point. The $36 \times 36$ mass matrix obtained from the second derivatives of the potential produces 36 mass eigenvalues, 15 of them belonging to massless Goldstone bosons. The latter will become the longitudinal components of $W^\pm$, $Z^0$ and 12 heavier vector bosons. We will discuss the mass values and properties of the lightest of the massive scalars.

The vev’s $v_1$ and $v_2$ in $H$ and $\tilde{H}$ are responsible for the low scale spontaneous symmetry breaking. If they have the same value - for instance as a consequence of the same origin of the vev’s of $H$ and $\tilde{H}$ - the corresponding two scalar bosons can be nearly degenerate. They are then twins. These twins consists of field components from $H$ and also field components from $\tilde{H}$ which are not directly coupled to quarks and leptons. A suggested normalization of the potential gives them a mass of $\approx 123 \text{ GeV}$ close to the Higgs boson observed at the LHC [2]. However, the near mass degeneracy of the twins will only be present if the influence of the remaining states is not severe.

### II. POTENTIAL FORMED FROM INVARIANTS OF $H$ AND $\tilde{H}$ SEPARATELY.

The potential for the 36 fields has to be formed from $SU(3)_L \times SU(3)_R$ invariants. In a first step we consider separately invariants of $H$ and invariants of $\tilde{H}$. Invariants which combine $H$ with $\tilde{H}$ fields will be added later. As mentioned in the introduction tree type mass terms are not included in the construction. Our construction is based on the vevs of invariants only. They fix the position of the minimum. Up to order 4 in the fields one remains with the invariants (the determinants of $H$ and $\tilde{H}$ play no role at this stage)

$$ J_1 = (Tr[H^\dagger \cdot H])^2, \quad J_2 = Tr[H^\dagger \cdot H \cdot H^\dagger \cdot H], \quad J_3 = (Tr[\tilde{H}^\dagger \cdot \tilde{H}])^2, \quad J_4 = Tr[\tilde{H}^\dagger \cdot \tilde{H} \cdot \tilde{H}^\dagger \cdot \tilde{H}]. \tag{II.2} $$

The potential $V(H, \tilde{H})$ is a superposition of the 4 invariants for which all 36 first derivatives have to vanish at the proposed minimum. As in [1] the superposition of the invariants must necessarily contain logarithmic functions of the invariants. In order to have only one logarithmic term, [II.1] will be restricted to contain a single large scale only. Therefore, we take $M_2$ and $M_3$ proportional to the large scale $M$.

For the potential we choose a naive phenomenological ansatz with coupling constants $c_i$ for the quartic interactions.

$$ V_0 = \frac{1}{t} (c_1 J_1 + c_2 J_2 + c_3 J_3 + c_4 J_4) \tag{II.3} $$

with

$$ t = \kappa + \log\left(\frac{J_1 J_2 J_3 J_4}{\langle J_1 \rangle \langle J_2 \rangle \langle J_3 \rangle \langle J_4 \rangle}\right). $$
The constant $\kappa$ introduced here allows to take a convenient value for the denominator in the logarithmic term: the product of vevs of the four invariants.

To get the stationary point at the wanted position one can perform the shift

$$
\begin{align*}
h_1^0 &\to v_1 + h_1^0, \quad h_2^0 \to b + h_2^0, \quad h_3^0 \to M + h_3^0, \quad \tilde{h}_1^0 \to v_2 + \tilde{h}_1^0, \\
\tilde{h}_2^0 &\to b_3 + \tilde{h}_2^0, \quad \tilde{h}_3^0 \to M_3 + \tilde{h}_3^0.
\end{align*}
$$

(II.4)

Now the vanishing of all derivatives of $V_0$ is required for the point $H = \tilde{H} = 0$ of the shifted fields.

The condition for the vanishing of all first derivatives at the point $H = \tilde{H} = 0$ for the shifted fields has a straightforward solution. It determines the parameter $\kappa$ and the coefficients $c_2$, $c_3$, $c_4$ in terms of $c_1$:

$$
\begin{align*}
\kappa &= 4, \quad c_2 = c_1 \left(\frac{b^2 + M^2 + v_1^2}{b^4 + M^4 + v_1^4}\right), \\
c_3 &= c_1 \left(\frac{(b^2 + M^2 + v_1^2)^2}{(b_2^2 + b_3^2 + v_2^2 + M_2^2 + M_3^2)^2}\right), \\
c_4 &= c_1 \left(\frac{(b^2 + M^2 + v_1^2)^2 + v_2^2 + 2(b_2 M_2 + b_3 M_3)^2 + (M_2^2 + M_3^2)^2}{(b_2^2 + b_3^2 + v_2^2 + M_2^2 + M_3^2)^2}\right).
\end{align*}
$$

(II.5)

Obviously, for $M_2^2 + M_3^2 = M^2$, one has $c_2 = c_3 = c_4 = c_1$ in the large $M$ limit.

The second derivatives of $\frac{1}{2}V_0$ at the point $H = \tilde{H} = 0$ of the shifted fields provide now for the $36 \times 36$ mass matrix. Degenerate states can be expected in case there is a relation between the vevs of $\tilde{H}$ with the vevs of $H$. Such a situation may arise by a high scale spontaneous symmetry breaking fixing simultaneously the vacuum expectation values of both matrix fields. The closest connection is a one-to-one correspondence of all vevs of $\tilde{H}$ with the vevs of $H$ combined with the requirement $\langle \det \tilde{H} \rangle = \pm \langle \det H \rangle$. One gets

$$
\langle h \rangle = \begin{pmatrix}
v_1 & 0 & 0 \\
0 & b & 0 \\
0 & 0 & M
\end{pmatrix}, \quad \langle \tilde{h} \rangle = \begin{pmatrix}
v_1 & 0 & 0 \\
0 & 0 & \pm b \\
0 & M & 0
\end{pmatrix},
$$

(II.6)

i.e. $v_2 = v_1$, $b_2 = 0$, $b_4 = \pm b$, $M_2 = M$, $M_3 = 0$. It follows

$$
\langle J_3 \rangle = \langle J_1 \rangle \quad \text{and} \quad \langle J_4 \rangle = \langle J_2 \rangle
$$

(II.7)

suggesting already that mass degeneration will occur. In the present section the sign of $b$ in $\langle \tilde{h} \rangle$ does play no role. But for latter purposes we will denote the form with the $+$ sign as case 1, the form with the $-$ sign as case 2. In the latter $\langle \tilde{h} \rangle$ differs from $\langle h \rangle$ only by an $SU(2)_R$ $U$-spin rotation with angle $\frac{\pi}{2}$.

This correspondence between the vevs of both matrix fields has besides (II.7) the consequence $c_3 = c_1$, $c_4 = c_2$. The mass matrix and mass eigenvalues resulting from the potential (II.3) can now easily be obtained. The masses are shown here to second order in $v_1$ for large $M$.

$$
\begin{align*}
m_1^2 &= c_1 (v_1^2 + b^2), \quad m_2^2 = c_1 (v_1^2 + b^2), \quad m_3^2 = c_1 (4M^2 + 5v_1^2 + 5b^2), \quad m_4^2 = c_1 (4M^2 + 5v_1^2 + 5b^2), \\
& \quad m_i^2 = 0 \quad i = 5, \ldots, 36.
\end{align*}
$$

(II.8)

There is a clear degeneracy of the two low and the two high mass states. The equality of the low masses is entirely due to the identification $v_2 \to v_1$. It holds (approximately) also in the more general form (II.1) provided $M_2^2 + M_3^2 = M^2$ and the $b$’s are small compared to $v_1 = v_2$. The 32 mass zero states can be separated into 17 states of mass zero and 15 mass zero Goldstone states. One zero mass state is due to the so far unbroken general phase transformation of $H$ and $\tilde{H}$. The remaining 16 zero mass states are due to our provisional neglection of invariants which connect the fields in $H$ with the fields in $\tilde{H}$. Without them $H$ or $\tilde{H}$ can be independently transformed by $SU(3)_L \times SU(3)_R$ matrices with no change of the invariants.

Since $b$ is sufficiently small (the b-quark mass is around $3 - 4$ GeV), it follows from $v_1^2 + v_2^2 = v^2$ the numerical value $v_1 = v_2 = \frac{\sqrt{2}}{2} = 123$ GeV, a value quite close to the mass of the Higgs boson found at LHC.

We identify now these degenerate states, these twins, with the Higgs particle of about 125 GeV found at the LHC. The twins are non-identical twins. They are combinations of the field $h_1^0$, the usual Higgs field which couples to fermions, and the field $\tilde{h}_1^0$ which has no Yukawa coupling to fermions and only couples to quarks and leptons via its coupling to gauge bosons. The mass value obtained at LHC fixes the coefficient $c_1$ in (II.3) to be very close two 1: $m_{Higgs} = v_1\sqrt{c_1} \simeq 125$ GeV, i.e. $c_1 \simeq 1.03$. 
The requirement of vanishing first derivatives of \( V \) parameter of the new invariants are kept of order 1. For this reason we put the factor \( c \) the positivity of the mass matrix is important. The contributions of \( J \) provide non-zero masses for all fields except the Goldstone ones. These two also break the symmetry under pure phase transformations. According to our vacuum structure the input vacuum expectation value \( M \) and in particular the influence of not yet considered invariants combining \( H \) and \( \tilde{H} \) can modify this value and lift the degeneracy.

### III. COMBINING \( H \) AND \( \tilde{H} \) FIELDS AND THE MASS SPECTRUM.

The inclusion of invariants neglected up to now is necessary to get an acceptable mass spectrum: in order not to be in conflict with experiments, all of the so far 17 massless states should become heavier than the standard model like Higgs. This cannot be achieved without problems and caveats.

There are quite a number of different invariants containing the fields of both multiplets \( H \) and \( \tilde{H} \). Most combinations for which all first derivatives vanish lead to some negative eigenvalues of the mass matrix. According to our vacuum structure all of the so far 17 massless states should become heavier than the standard model Higgs twin states, mix them with other states and change their masses or have troublesome discontinuities by a change of parameters. We take for the additional potential the combination

\[
V_S = \frac{v_1^2}{M^2} (r_1J_1 + r_2J_2 + r_3J_3 + r_4J_4 + r_5J_5) + r_6J_6 + r_7J_7 + r_8J_8 + r_9J_9
\] (III.10)

with the new invariants

\[
J_5 = Tr[H^\dagger \cdot \tilde{H} \cdot \tilde{H}^\dagger \cdot H], \quad J_6 = Tr[H^\dagger \cdot H \cdot \tilde{H}^\dagger \cdot \tilde{H}], \quad J_7 = Tr[H^\dagger \cdot \tilde{H} \cdot H^\dagger \cdot \tilde{H} + Tr[\tilde{H}^\dagger \cdot H \cdot \tilde{H}^\dagger \cdot H], \quad (III.11)
\]

\[
J_8 = b (\det H + \det H^\dagger), \quad J_9 = b (\det \tilde{H} + \det \tilde{H}^\dagger).
\]

The contributions of \( J_1...J_5 \) turn out to be extremely small compared to the ones for \( V_0 \) as long as the other \( r \) parameter of the new invariants are kept of order 1. For this reason we put the factor \( \frac{v_1^2}{M^2} \) in front of these invariants. The first seven invariants used here respect the symmetry \( H \rightarrow -H \) and \( \tilde{H} \rightarrow -\tilde{H} \), while the last two break this symmetry. These two also break the symmetry under pure phase transformations. According to our vacuum structure both will contribute because of the presence of the \( b \) terms in (III.6). The total potential \( V = V_0 + V_S \) should now provide non-zero masses for all fields except the Goldstone ones.

As a first try one could set the small \( b \)'s to zero and needs then only the first seven invariants in (III.10). A very appealing mass spectrum is achieved with all other masses lying above the (almost) degenerate twins. However, even the tiniest change of \( b \) in (11) away from zero destroys this picture. A fatal discontinuity is present changing abruptly the possibility of the mass matrix \( V \).

Let us then use the vacuum structure (III.6) where the continuous parameter \( b \) appears in both matrix fields. In fact, this choice appears to be the only one for which the potential formed from the 9 invariants allows for a positive definite mass matrix with masses above the (almost) degenerate twins. As free parameters one can use \( r_1, r_2 \) and \( r_7 \). The requirement of vanishing first derivatives of \( V \) - after shifting the fields - fixes then the remaining parameters \( r_3, r_4, r_5, r_6, r_8 \) and \( r_9 \). For the case 1 (the plus value in (III.6)) one obtains

\[
\begin{align*}
    r_3 &= r_1, \quad r_4 = r_2, \quad r_9 = -r_8, \quad r_5 = -2(r_1 + r_2) - \frac{2r_1v_1^2}{b^2 + M^2}, \\
    r_6 &= -2(b^2M^4(b^2 + M^2)r_7 + b^2M^2(b^2 + M^2)r_1v_1^2 - (b^4r_1 + b^2M^2(r_1 + r_7) + M^4(r_1 + r_7)))v_1^4 + r_1v_1^8, \\
    r_8 &= -\frac{2r_1v_1^2(b^2 + M^2 - 2v_1^2)(b^2 + M^2 + v_1^2)}{M(b^2 + M^2)(b^2M^2 - v_1^4)}.
\end{align*}
\] (III.12)

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1 I am very much indebted to my colleague Werner Wetzel who advised me about the multiple occurrence of these shrill discontinuities when proposing the vev of invariants.
The relations between the excluded because its production cross-section is lower than the one of a standard Higgs boson.

discontinuities it is necessary to say that the twin picture can only be maintained if there is a very strict relation between the vevs of the two multiplets like the one suggested in (II.6 ).

If we increase the parameter $r_1$ by a factor 2 the mass of the heavier of the twins is increased making the mass difference of the twin states about 2 GeV. A change of the numerical value of $b$ is of little influence. Therefore, the model appears to be a good example for a model with a twin structure of the lowest scalar states with properties which are not easily distinguishable from a pure standard model Higgs. One has a smooth behavior when changing the parameters within a large region. However, if one changes the vacuum structure by adding even a tiny new element, for instance $\langle \tilde{h}_0^3 \rangle$, the mass matrix changes abruptly to an unphysical, not positive definite form. Because of these discontinuities it is necessary to say that the twin picture can only be maintained if there is a very strict relation between the vevs of the two multiplets like the one suggested in (II.6).

If the Higgs boson is not a twin, one can still have the superposition of the two types of fields and near degeneracies of states of higher mass. The sign of $b$ in the vacuum structure of $\tilde{H}$ could be negative, an option we called case 2. The relations between the $r$ values change and another range of input parameters $r_1$, $r_2$ and $r_7$ can be used. In spite of $v_2 = v_1$ the lowest state is then non-degenerate and the next state much higher in mass. Still this boson is an almost equal superposition of the fields $h_1^1$ and $\tilde{h}_1^1$. Thus it is not a usual Higgs boson. It can soon be experimentally excluded because its production cross-section is lower than the one of a standard Higgs boson.

In a modified form of case 2 one may take $v_2$ different from $v_1$, close to zero. As an example we take $M = 10^{13}$, $b = 4$ GeV, but $v_1 \approx v$ and $v_2 \approx 5$ GeV to be slightly bigger than $b$. The corresponding value for $c_1$ is $c_1 \approx \frac{1}{2}$. The lowest boson has then again a mass of $\approx 123$ GeV. Choosing $r_1 = r_2 = 4$ and $r_7 = 2 \cdot 10^{-3}$ this state is not degenerate and the corresponding field is to 99% the field $h_1^1$. In this example the next 4 states have almost the same mass of 984 GeV (the two charged ones with identical masses) and are all equal weight superpositions with regard to the two fields of different couplings to fermions. In case 1 and in case 2 the bosons 3 and 4 have opposite electromagnetic charges and are strictly degenerate by CP invariance.

The clarification whether the Higgs boson is a single resonance or a twin, or perhaps a degenerate state of a two-Higgs-doublet model \[10\], requires more experimental data and a very detailed analysis by experts.

IV. CONCLUSIONS

By extending the Glashow-Weinberg-Salam group to $SU(3)_L \times SU(3)_R \times SU(3)_C$ as in [7,8] the scalar sector consists of two $(3^*, 3, 1)$ multiplets $H$ and $\tilde{H}$ where only the fields in $H$ have a Yukawa coupling to quarks and leptons. The vacuum expectation values of $H$ and $\tilde{H}$ contain, besides the low mass elements responsible for the electro-weak symmetry breaking, also elements having a very high scale. A phenomenological potential has been constructed which reproduces these vevs. This invariant potential is taken to be of the Coleman-Weinberg type \[9\] and thus has no tree-level mass term. It is suggestive to assume a relation between the vev’s of $H$ and $\tilde{H}$, in particular, to take the low scale elements $v_1$ in $H$ and $v_2$ in $\tilde{H}$ to be equal. Using at first only very simple $SU(3)_L \times SU(3)_R$ invariants, this assumption has the consequence that the two lowest scalar states are degenerate mass eigenstates with masses near $\sqrt{\frac{\lambda}{\mu}} = 123$ GeV. The two states are mixtures of a Higgs field with conventional properties and a Higgs field which can couple to fermions only via gauge vector bosons. Problematic for maintaining the degeneracy is the necessity of including more invariants in order to get appropriate masses for the remaining fields. We considered the possibility for a one-to-one correspondence of the vevs of both multiplets and of their determinants. In the main case considered the twin structure can survive and the mass spectrum can be chosen to be in accord with experimental constraints. We pointed out that by constructing potentials from vacuum expectation values discontinuities can appear which can only be tolerated if very strict relations between the vevs of the two multiplets hold. Nevertheless, it may be worthwhile to look for a twin structure of the Higgs boson found at the LHC.
Let me add a word on the quadratic divergencies. In the standard model - at least to one loop - the quadratic divergence can be viewed as being solely a problem for the vacuum expectation value of the Higgs field. The quadratic divergence of this quantity does not affect other particle properties and only indirectly the particle masses due to their couplings to the Higgs field. Vacuum expectation values (like the cosmological constant) are not yet understood. But it is well known that by fine tuning the vacuum expectation value to its experimental value, or by subtracting the relevant tadpole graphs, or by assuming their cancellation at a very high scale [11, 12], these divergencies have no further consequences for the particle spectrum. In the present phenomenological approach the only input are vev’s - which themselves are certainly affected by quadratic divergencies - but are taken to be fixed. Thus one can expect, that in attempts of this type, the particle spectrum and other particle properties are not influenced by our non-understanding of vacuum expectation values and their quadratic divergencies.

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