The quantum vacuum of the minimal \( \text{SO}(10) \) GUT

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Abstract. We reexamine the longstanding no-go excluding all potentially viable \( \text{SO}(10) \to \text{SU}(3)_c \otimes \text{SU}(2)_L \otimes \text{U}(1)_Y \) symmetry breaking patterns within the minimal renormalizable non-supersymmetric \( \text{SO}(10) \) GUT framework featuring the 45-dimensional adjoint representation in the Higgs sector. A simple symmetry argument indicates that quantum effects do change the vacuum structure of the model dramatically. A thorough analysis of the one-loop effective potential reveals that the phenomenologically favoured symmetry breaking chains passing through the \( \text{SU}(4)_C \otimes \text{SU}(2)_L \otimes \text{U}(1)_Y \) or \( \text{SU}(3)_c \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{U}(1)_{B-L} \) intermediate stages are, indeed, supported at the quantum level. This brings the class of minimal non-supersymmetric \( \text{SO}(10) \) GUTs back from oblivion, providing a new ground for a potentially realistic model building.

1. Introduction

In spite of the enormous progress the field of grand unified theories (GUTs) experienced since the seminal works of Georgi and Glashow \([1]\) it is still not clear whether the basic scheme is by any means reflected by nature. Nevertheless, the inherent testability of the simplest GUT scenarios makes the idea of a unified description of the strong and electroweak interactions so popular that there has been a lot of efforts spent recently in attempts to pin down the simplest potentially viable scenarios, in particular those based on the \( \text{SU}(5) \) \([2]\) and/or \( \text{SO}(10) \) \([3, 4]\) gauge symmetry. Apart from the canonical test imposed by the required high-energy convergence of all three Standard Model (SM) gauge couplings, a successful candidate should obey also the constraints inflicted by the proton longevity (stretching over \( 10^{33} \) years) and the absolute neutrino mass scale, which are indeed closely related to the generic need for a compatibility with the observed SM flavour structure.

However, the situation is not very clear at the moment as even the very concept of minimality admits a certain number of interpretations. For instance, it can be understood as a mere “simplicity” related to e.g. a minimum rank of the gauge group, the Higgs sector dimensionality, the naturalness of the doublet-triplet splitting, or even the complexity of the gauge unification pattern. In this respect, the models based on \( \text{SU}(5) \) gauge group are frequently regarded to as more minimal than the \( \text{SO}(10) \) ones. Alternatively, predictivity (i.e., the number of independent couplings) is often invoked as a measure of minimality; what matters in such a case is namely the renormalizability of the theory, the number of Higgs multiplets, the complexity of the GUT-scale Yukawa sector, the presence or absence of supersymmetry (SUSY) etc. In this respect, the \( \text{SO}(10) \) GUTs (especially their SUSY variants) usually score better than the \( \text{SU}(5) \) models, see...
e.g. [4, 5] and references therein. To us, even though potentially troublesome on the naturalness side, this is a very physical option - besides giving generally better prospects for an ultimate testability the preference of $SO(10)$ also offers a natural relief from the persistent trouble with the simplest $SU(5)$ models, failing either on the unification side (like e.g. the original Georgi-Glashow model) or suffering from an overly fast proton decay when SUSY is invoked\(^1\).

Sticking to the $SO(10)$ case, minimality is essentially equivalent to the complexity of the Higgs sector. In the class of $SO(10)$ GUTs, the choice of the Higgs representation is subject to several basic requirements. First, there is the need to break consistently the rank=5 $SO(10)$ gauge symmetry down to the rank=4 $SU(5)$ of the SM. Remarkably, none of the basic candidate Higgs representations (namely, $16$ and/or $126$) capable of breaking the $U(1)$ subgroup of $SO(10)$, can account for the full $SU(5)$ subgroup breakdown on its own, since their SM singlets leave a full $SU(5)$ subgroup of $SO(10)$ intact. From the group theory point of view, the simplest solution one can devise is an extra adjoint $45$ in the Higgs sector whose additional pair of VEVs brings in enough freedom to provide the full $SO(10) ightarrow SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ breakdown. Unfortunately, this approach fails in the SUSY case where the extra $F$-flatness conditions force the VEVs in $45$ to align along the direction of $\langle 16 \rangle$ and thus $SU(5)$ still remains intact \cite{6}. Thus, $210$ is usually invoked in the SUSY $SO(10)$ context, giving rise to the minimal SUSY $SU(10)$ GUT with a Higgs sector spanning over the reducible representation $10 \oplus 126 \oplus 126 \oplus 210$, see e.g. \cite{7, 8} and also \cite{4}.

However, though successful in accounting for some of the basic features of the lepton flavour mixing \cite{9}, the minimal renormalizable SUSY $SO(10)$ model fails due to a generic tension between the absolute neutrino mass scale and the gauge coupling unification constraints, \cite{10, 11} and there is a general need to go beyond the simplest setting, see e.g. \cite{12, 13, 14, 15} and references therein.

On the other hand, if predictivity rather than naturalness is at stakes, supersymmetry is not a mandatory ingredient of the $SO(10)$ GUTs because the unification constraints refuting the simplest non-SUSY $SU(5)$ settings can be satisfied with an intermediate energy scale $M_I$ below the scale of grand unification $M_G$. In particular, the recent analysis of the relevant symmetry breaking chains \cite{16} reveals an interesting picture associated to the two basic single-step chains

\begin{equation}
SO(10) \xrightarrow{M_G} SU(4)_C \otimes SU(2)_L \otimes U(1)_R \xrightarrow{M_I} SM,
\end{equation}

\begin{equation}
SO(10) \xrightarrow{M_G} SU(3)_C \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \xrightarrow{M_I} SM,
\end{equation}
driven by the simplest conceivable Higgs sector spanned over just $45 \oplus 16$ of $SO(10)$. Unfortunately, this scenario, requiring in practice a few orders of magnitude hierarchy between $M_I$ and $M_G$, has been long ago \cite{17, 18, 19} claimed incompatible with the dynamics of the scalar sector. The reason is an emergence of tachyonic masses in the Higgs spectrum unless $M_I$ essentially coincides with $M_G$, at odds with the running constraints.

All these arguments are, however, based on the tree-level analysis of the Higgs potential. Remarkably enough, the recent work \cite{20} reveals a dramatic qualitative change of the picture once quantum effects are properly taken into account, with a potential of bringing the class of minimal non-supersymmetric $SO(10)$ GUTs back to life.

2. The classical failure of the minimal non-SUSY $SO(10)$ GUT - a closer look

In order to fully appreciate this fact one has to inspect in detail the available vacua of the renormalizable non-SUSY scalar potential spanned over $45 \oplus 16$ in the Higgs sector. However, before entering the realm of quantum effects one has to understand the origin of the tree-level no-go implied by \cite{17, 18, 19}.

\(^1\) For a nice account for the proton decay issue in the minimal SUSY $SU(5)$ see e.g. \cite{2}.
The most generic classical renormalizable scalar potential one can write with only $45 \oplus 16$ at hand reads $V_0 = V_{45} + V_{16} + V_{\text{int}}$, where [21]

\begin{align}
V_{45} &= -\frac{\mu^2}{2} \text{Tr} 45^2 + \frac{a_1}{4} (\text{Tr} 45^2)^2 + \frac{a_2}{4} \text{Tr} 45^4, \\
V_{16} &= -\frac{\nu^2}{2} 16^1 16 + \frac{\lambda_1}{4} (16^1 16)^2 + \frac{\lambda_2}{4} (16^1 16)(16^1 16^+) \\
\end{align}

(3)

(4)

correspond to the mass and self-interaction terms for $45$ and $16$, respectively, while

\begin{align}
V_{\text{int}} &= \alpha(16^1 16) \text{Tr} 45^2 + \beta 16^1 45^2 16 + \tau 45 16 \\
\end{align}

(5)

provides the desired interaction between the two sectors. Employing an $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ notation the vacuum manifold is parametrized by the VEVs of the three different SM singlets

\begin{align}
\langle (1, 1, 1, 0)_{45} \rangle \equiv \omega_Y, \quad \langle (1, 1, 3, 0)_{45} \rangle \equiv \omega_Y, \quad \langle (1, 1, 2, +1/2)_{16} \rangle \equiv \chi_{B-L}. \\
\end{align}

Since the gauge unification in this minimal scenario typically requires the $B-L$ breaking scale $\chi_{B-L}$ to be lower than the higher from the pair $\omega_R, \omega_Y$ (to be identified with $M_G$), the heavy Higgs spectrum is governed namely by the VEVs of $45$. Apart from the generic setting with non-zero $\omega_R \neq \omega_Y$ corresponding to an intermediate $SU(3)_c \otimes SU(2)_L \otimes U(1)_{R} \otimes U(1)_{B-L}$ stage, there are several important limiting cases exhibiting an enhanced intermediate symmetry, in particular

\begin{align}
\omega_R \neq 0, \omega_Y = 0 : \quad &SU(4)_C \otimes SU(2)_L \otimes U(1)_R, \\
\omega_R = 0, \omega_Y \neq 0 : &SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}, \\
\omega_R = \pm \omega_Y \neq 0 : &\text{standard (+) or flipped (−) } SU(5) \otimes U(1). \\
\end{align}

(7)

The point is that while the first two cases (corresponding to the simplest conceivable chains $(1, (2))$ are the only ones from list (7) which are compatible with the non-SUSY unification constraints [20], they seem to be disfavoured by the detailed Higgs sector dynamics [17, 18, 19] due to the emergence of tachyons within the corresponding vacua.

2.1. The emergence of tachyons in the non-$SU(5) \otimes U(1)$ vacua

This is due to the fact that a full-fledged analysis of the classical Higgs potential (3)-(5) reveals a pair of scalar multiplets with the classical SM-level masses given by the formulae

\begin{align}
M^2(1, 3, 0)_{45} &= 2a_2(\omega_Y - \omega_R)(\omega_Y + 2\omega_R), \\
M^2(8, 1, 0)_{45} &= 2a_2(\omega_R - \omega_Y)(\omega_R + 2\omega_Y). \\
\end{align}

(8)

It is clear that in the majority of the parametric space, in particular, everywhere outside the domain

\begin{align}
-2 \leq \omega_Y/\omega_R \leq -1/2, \\
\end{align}

(9)

at least one of them develops a negative mass-squared, indicating an instability of the corresponding vacuum configuration. Consequently, all stable settings exhibit an intermediate flipped $SU(5) \otimes U(1)$ stage, at odds with unification constraints. This, indeed, is the core of the classical argument against the potential viability of the minimal non-SUSY $SO(10)$ GUT framework.

Since the simplicity of the formulae (8) is crucial for the above reasoning it is important to understand their structure in as much detail as possible. In this respect, it is particularly illuminating to discuss the global symmetries of the scalar potential $V_0$. 

3
2.2. Understanding the structure of the key mass formulae in terms of symmetries

When only trivial invariants (i.e. moduli) of both 45 and 16 are considered, the global symmetry of $V_0$ is $O(45) \otimes O(32)$. It is spontaneously broken down to $O(44) \otimes O(31)$ by the 45 and 16 VEVs yielding $44+31=75$ Goldstone bosons in the scalar spectrum. Since in this case also the gauge $SO(10)$ symmetry is broken to the SM gauge group, $45 - 12 = 33$ would-be Goldstone bosons, with the quantum numbers of the coset $SO(10)/$SM algebra, decouple from the physical spectrum, while $75 - 33 = 42$ pseudo-Goldstone bosons remain. Their masses are generally expected to receive contributions from the explicitly breaking terms $a_2$, $\lambda_2$, $\beta$, and $\tau$.

However, the tree-level masses of the crucial $(1,3,0)_{45}$ and $(8,1,0)_{45}$ multiplets, which belong among these pseudo-Goldstone bosons, turn out to depend only on the parameter $a_2$ but not on the other parameters expected, namely $\lambda_2$, $\beta$, and $\tau$. While the $\lambda_2$ and $\tau$ terms cannot obviously contribute at the tree level to 45 mass terms, one would generally expect a contribution from the $\beta$ term proportional to $\chi_{B-L}^2$ to pop up in relations (8). This, however, is not the case. One can understand that on general grounds by observing that the relevant scalar interaction in (5) resembles the the covariant-derivative interaction responsible for the contribution of $(16)$ to the gauge boson masses. As a consequence, no tree-level mass contribution from the $\beta$ coupling can be generated for the scalars in 45 which carry the SM quantum numbers of the $W^\pm$ and $Z^0$ bosons, i.e., $(1,3,0)_{45}$, and gluons $(8,1,0)_{45}$ that are protected by the residual $SU(5)$ symmetry left unbroken even with $\langle 16 \rangle \neq 0$.

3. The quantum structure of the minimal non-SUSY $SO(10)$ GUT vacuum

On the other hand, we should expect the $\tau$ and $\beta$ interactions to contribute to the masses of $(1,3,0)_{45}$ and $(8,1,0)_{45}$ at the quantum level. While for the former the loops simply admit to contract the two 16’s sticking out of the $\tau$-vertex in (5), a similar loop due to a pair of $\beta$-vertices brings into play the $SU(5)$-breaking VEVs of 45. Similar contributions should also arise from the gauge interactions which break explicitly the independent global transformations on the 45 and 16 discussed in the previous section.

In the diagrammatic language, the relevant one-loop self-energies are depicted in Figure 1 where the individual components of the scalars transforming as 45 and 16 of $SO(10)$ have been denoted by generic symbols $\phi$ and $\chi$, respectively. While the exchange of the components of 16 is crucial, the $\chi_{B-L}$ does not need to enter; in the phenomenologically allowed unification patterns it gives actually negligible contributions.

![Figure 1. The diagrammatic structure of the three important types of quantum corrections which stabilise the vacua supporting the desired $SU(4)_C \otimes SU(2)_L \otimes U(1)_R$ (case A) or $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ (case B) symmetry breaking chains at the one-loop level. The dangerous tachyonic mass parameters of the $(1,3,0)_{45}$ and/or $(8,1,0)_{45}$ Higgs multiplets in 45 are lifted provided $8\pi^2|a_2| < \tau^2/\omega_R^2 + \beta^2 + 13g^4$ in case A, or $8\pi^2|a_2| < \tau^2/\omega^2 + 2\beta^2 + 19g^4$ in case B. The interested reader can find further details in the original work [20].](image)

One is thus lead to the conclusion that any result based on the particular shape of the tree-level 45 vacuum is dramatically altered at the quantum level. Let us emphasise that although one may in principle avoid the $\tau$-term by means of e.g. an extra $Z_2$ symmetry, no symmetry can forbid the $\beta$-term and the gauge loop contributions.
3.1. The one-loop effective potential analysis

The quantum shifts in the masses of the two critical multiplets read

\[ \Delta M_{ab} = \frac{1}{16\pi^2} \text{Tr} \left[ \mathcal{M}_s^2(\phi, \chi) \left( \log \frac{\mathcal{M}_s^2(\phi, \chi)}{\mu^2} - \frac{3}{2} \right) + 3 \mathcal{M}_g^4(\phi, \chi) \left( \log \frac{\mathcal{M}_g^2(\phi, \chi)}{\mu^2} - \frac{5}{6} \right) \right] . \]

(10)

Here the terms in the bracket correspond to the contribution from the (real) scalars and gauge bosons, respectively, and \( \mathcal{M}_s \) and \( \mathcal{M}_g \) denote the relevant field-dependent tree-level mass matrices. Note also that the contribution due to the SM fermions has been neglected because of the hierarchical suppression of their masses with respect to the GUT-scale Higgs and gauge bosons. For completeness, let us also mention that the specific form of \( \Delta M \) corresponds to the Landau gauge and the \( \overline{\text{MS}} \) renormalization scheme.

Using a generic symbol \( \psi \) for any of the components of \( 45 \oplus 16 \) the first derivative of the one-loop part of the effective potential \( \Delta V \) is given by

\[ \frac{\partial \Delta V}{\partial \psi} = \frac{1}{16\pi^2} \text{Tr} \left\{ \frac{\partial^2 \mathcal{M}_s^2}{\partial \psi_a \partial \psi_b}, \mathcal{M}_s^2 \right\} \left( \log \frac{\mathcal{M}_s^2}{\mu^2} - \frac{3}{2} \right) + \mathcal{M}_s^2 \frac{\partial^2 \mathcal{M}_s^2}{\partial \psi_a \partial \psi_b} + \text{gauge part} , \]

where, for the sake of simplicity, only the scalar contribution due to the first term in eq. (10) has been written explicitly; the gauge part derived from the \( \mathcal{M}_g \)-piece of (10) is completely analogous. Let us also note that the formula above holds regardless of the commutation properties of \( \mathcal{M} \) and its first derivatives.

Similarly, the second derivatives of \( \Delta V \) can be written in the form

\[
\frac{\partial^2 \Delta V}{\partial \psi_a \partial \psi_b} = \frac{1}{16\pi^2} \text{Tr} \left\{ \left[ \frac{\partial^2 \mathcal{M}_s^2}{\partial \psi_a \partial \psi_b}, \mathcal{M}_s^2 \right] + \left[ \frac{\partial \mathcal{M}_s^2}{\partial \psi_a} \frac{\partial \mathcal{M}_s^2}{\partial \psi_b} \right] \left( \log \frac{\mathcal{M}_s^2}{\mu^2} - \frac{3}{2} \right) + \frac{\partial \mathcal{M}_s^2}{\partial \psi_a} \frac{\partial^2 \mathcal{M}_s^2}{\partial \psi_b} + \mathcal{M}_s^2 \frac{\partial^2 \mathcal{M}_s^2}{\partial \psi_a \partial \psi_b} \right. \\
+ \left. \sum_{m=1}^{\infty} (-1)^{m+1} \frac{1}{m} \sum_{k=1}^{m} \left( m \atop k \right) \mathcal{M}_s^{2k} \frac{\partial^k \mathcal{M}_s^2}{\partial \psi_a} \mathcal{M}_s^{2m-k} \frac{\partial^{m-k} \mathcal{M}_s^2}{\partial \psi_b} \left( \mathcal{M}_s^2 - 1 \right)^{m-k} \right\} + \text{gauge part} ,
\]

(12)

where the infinite series of nested commutators (taken \( k - 1 \) times for each \( k \)) looks after the general non-commutativity of \( \mathcal{M}_s \) and \( \partial \mathcal{M}_s^2/\partial \psi \) which becomes relevant here.

Note also that such terms are also crucial for a consistent cancellation of the infrared divergencies emerging in the one-loop mass matrix \( \mathcal{m}_{\alpha\beta} = \partial^2 V_{\text{eff}}/\partial \psi_\alpha \partial \psi_\beta \) when evaluated for the solutions of the one-loop stationary equations \( \partial V_{\text{eff}}/\partial \psi_\alpha = 0 \). These are, of course, due to the presence of Goldstone bosons in the Landau gauge inherent to formula (10). Let us recall that the effective potential, however, is defined at zero external momenta, and the physical Higgs spectrum is derived from the “shifted” mass matrix \( \mathcal{M}_{\alpha\beta}(p^2) = \mathcal{m}_{\alpha\beta} + \Delta \Sigma_{\alpha\beta}(p^2) \) for \( p^2 = M_s^2 \) where \( \Delta \Sigma_{\alpha\beta}(p^2) \equiv \Sigma_{\alpha\beta}(p^2) - \Sigma_{\alpha\beta}(0) \) is the difference of the relevant one-loop self-energies evaluated at two different four-momenta.

3.2. The one-loop corrections to the masses of \((1, 3, 0)_{45}\) and \((8, 1, 0)_{45}\)

Putting all this together, the quantum shifts in the masses of the two critical multiplets read

\[
\Delta M^2(1, 3, 0)_{45} = \frac{1}{4\pi^2} \left[ \tau^2 + \beta^2 (2\omega_R^2 - \omega_R \omega_Y + 2\omega_Y^2) + g^4 \left( 16\omega_R^2 + \omega_Y \omega_R + 19\omega_Y^2 \right) \right] + \text{logs} ,
\]

\[
\Delta M^2(8, 1, 0)_{45} = \frac{1}{4\pi^2} \left[ \tau^2 + \beta^2 (\omega_R^2 - \omega_R \omega_Y + 3\omega_Y^2) + g^4 \left( 13\omega_R^2 + \omega_Y \omega_R + 22\omega_Y^2 \right) \right] + \text{logs} ,
\]

(13)
where the symbols “logs” denote the sub-leading logarithmic contributions. It is easy to see that with these terms added to the tree-level structure (8) both (1,3,0)$_{45}$ and (8,1,0)$_{45}$ mass-squares can be made simultaneously positive for wide ranges of the $\omega_R$ and $\omega_Y$ parameters, well outside the classical domain (9). This holds, especially, for the configurations with either $\omega_Y \ll \omega_R$ or $\omega_R \ll \omega_Y$ corresponding to the classically forbidden $SU(4)_C \otimes SU(2)_L \otimes U(1)_R$ (case A) or $SU(3)_C \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ (case B) intermediate symmetries. In particular, the dangerous tachyonic mass parameters of the (1,3,0)$_{45}$ and/or (8,1,0)$_{45}$ Higgs multiplets in 45 are lifted for $8\pi^2|a_2| < \tau^2/\omega_R^2 + \beta^2 + 13g^4$ in case A, or $8\pi^2|a_2| < \tau^2/\omega_Y^2 + 2\beta^2 + 19g^4$ in case B. Thus, the quantum effects restore the full symmetry breaking potential of the minimal non-SUSY SO(10) Higgs sector spanned over $45 \oplus 16$, bringing back into play the symmetry breaking chains favoured by the gauge coupling unification constraints! Let us also stress that the fact that $U(1)_{B-L}$ has been broken by the spinorial 16 does not play a major role in the argument so an alternative Higgs model based on $45 \oplus 126$ representation exhibits the same qualitative behaviour. Similarly, an extra 10 added to the Higgs sector in order to get potentially realistic Yukawa couplings makes no harm because it can not develop any GUT-scale VEVs.

4. Conclusions and outlook

Typically, only the minimal GUTs can be subject to a thorough scrutiny by complementary methods such as the proton lifetime constraints, the absolute neutrino mass scale or the match between the effective SM flavour structure and the underlying GUT-scale Yukawa sector. In this respect, the recent failure of the minimal renormalizable SUSY SO(10) GUT re-opened the question of the potential viability of the simplest non-SUSY SO(10) scenarios. Remarkably, the generic no-go obtained about 30 years ago, which plagued the simplest model with $45 \oplus 16$ in the Higgs sector, has been recently identified as a mere artefact of the classical approach that is lifted at the quantum level. This reiterates the intriguing question whether a potentially viable and intrinsically consistent non-SUSY SO(10) theory can be spanned over a renormalizable Higgs sector with only $45 \oplus 16$ (or, alternatively, $45 \oplus 126$) participating at the GUT symmetry breakdown. The final answer, however, is left to a future analysis.

References

[1] Georgi H and Glashow S L 1974 Phys. Rev. Lett. 32 438–441
[2] Bajc B, Fileviez Perez P and Senjanovic G 2002 (Preprint hep-ph/0210374)
[3] Albright C H and Barr S M 2000 Phys. Rev. D62 093008 (Preprint hep-ph/0003251)
[4] Aulakh C S, Bajc B, Melfo A, Senjanovic G and Vissani F 2004 Phys. Lett. B588 196–202 (Preprint hep-ph/0306242)
[5] Fukuyama T, Ilakovac A, Kikuchi T, Meljanac S and Okada N 2004 JHEP 0905 126 (Preprint hep-ph/0406068)
[6] Aulakh C S, Bajc B, Melfo A, Rasin A and Senjanovic G 2001 Nucl. Phys. B597 89–109 (Preprint hep-ph/0004031)
[7] Clark T E, Kuo T K and Nakagawa N 1982 Phys. Lett. B115 26
[8] Aulakh C S and Mohapatra R N 1983 Phys. Rev. D28 217
[9] Bajc B, Senjanovic G and Vissani F 2003 Phys. Rev. Lett. 90 051802 (Preprint hep-ph/0210207)
[10] Aulakh C S and Garg S K 2006 Nucl. Phys. B757 47 (Preprint hep-ph/0512224)
[11] Bertolini S, Schwetz T and Malinsky M 2006 Phys. Rev. D73 115012 (Preprint hep-ph/0605006)
[12] Bertolini S, Frigerio M and Malinsky M 2004 Phys. Rev. D70 095002 (Preprint hep-ph/0406117)
[13] Aulakh C S and Garg S K 2006 (Preprint hep-ph/0612021)
[14] Bajc B, Dorner I and Nemevsek M 2008 JHEP 11 007 (Preprint 0809.1069)
[15] Melfo A, Ramirez A and Senjanovic G 2010 (Preprint arXiv:1005.0834)
[16] Bertolini S, Di Luzzio L and Malinsky M 2009 Phys. Rev. D80 015013 (Preprint arXiv:0903.4049 [hep-ph])
[17] Yasue M 1981 Phys. Rev. D24 1005; Yasue M 1981 Phys. Lett. B103 33
[18] Anastaze G, Derendinger J P and Buccella F 1983 Z. Phys. C20 269–273
[19] Babu K S and Ma E 1985 Phys. Rev. D31 2316
[20] Bertolini S, Di Luzzio L and Malinsky M 2010 Phys. Rev. D81 035015 (Preprint arXiv:0912.1796 [hep-ph])
[21] Buccella F, Ruegg H and Savoy C A 1980 Phys. Lett. B94 491
[22] Coleman S R and Weinberg E J 1973 Phys. Rev. D7 1888–1910