MSSM HIGGS SECTOR AT THE ONE-LOOP LEVEL

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Abstract

This work provides an elementary introduction to the Higgs sector
renormalisation within the Minimal Supersymmetric Standard Model
(MSSM) framework. The main aim of the paper is to clarify some tech-
nical details that are usually omitted in the existing literature. The MSSM
tree-level relation \( m_h^2 + m_H^2 = m_A^2 + m_Z^2 \) is renormalised using the stan-
dard technique of direct computation of the relevant one-loop Feynman
diagrams. The calculation is performed within the unitary gauge and
the definition of the renormalised parameters is briefly reviewed. The ex-
pected cancellation of ultraviolet divergences is explicitly checked and the
well-known leading-log term is recovered. All the necessary ingredients of
the computations are summarized in the appendices which makes the work
more self-contained.

1 Introduction

The Minimal Supersymmetric Standard Model (MSSM) provides one of the first
realistic attempts to describe physics beyond the Standard Model. Although
there is a huge amount of new fields and parameters in any SUSY-extended
model, the MSSM is able to reproduce all the successful predictions of the original
GWS theory with a very good accuracy. The MSSM Higgs sector is enlarged
due to the fact that supersymmetry forbids the usual mechanism of generation
of the up-type quark masses. This problem requires introduction of another
Higgs doublet with opposite charges, which produces all the necessary quark
mass-terms. Consequently there are five Higgs particles and three unphysical
Goldstone bosons in the theory.

The global \( N = 1 \) supersymmetry, being a new symmetry in addition to
the original \( SU(3)_c \otimes SU(2)_L \otimes U(1) \) local gauge invariance, puts constraints
on the Higgs sector, which is rather indefinite in nonsupersymmetric theo-
ries. In the case of MSSM this leads to the famous tree-level relation \( m_h \leq m_Z |\cos 2\beta| \) \footnote{to be published in Czech Journal of Physics}
which bounds the mass of the lightest Higgs scalar. Note that up to now there is no experimental evidence of such state; the present day
lower limit is \( m_h \geq 88.3 \text{ GeV} \), see \footnote{malinsky@hp02.troja.mff.cuni.cz}. However, radiative corrections modify
the upper bound substantially \footnote{\&}. This bound is closely related to the
tree-level sum-rule
\[ m_h^2 + m_H^2 - m_A^2 - m_Z^2 = 0 \]  

(1)

Renormalising this rule one can obtain the shift of the unsatisfactory bound descending from the radiative corrections. Up to now there are many papers performing the computation on the one-loop level [5, 6, 7, 8] and also some two-loop results were already obtained [9]. Most of these papers are concentrated on the evaluation of the finite part of the correction only without explicit discussion of the divergences. It is expected that the ultraviolet divergences in such relations originating from the additional symmetry cancel, but there is still no general proof of this based on the relevant Ward identities. From this point of view it may be useful to demonstrate explicitly the mechanism of the compensation in the particular case of relation (1); this, in fact, is one of the goals of this work. In this respect, this paper supplements the calculations presented in the cited literature.

The whole analysis is performed for the top-stop sector only. It is sufficient due to the observation that the leading term coming from any fermion-sfermion cluster is proportional to the fourth power of the fermionic mass [5, 6]. Note that the contributions coming from the chargino and neutralino sectors are negligible [7].

The paper is organized as follows: the definition of the renormalised parameters is briefly reviewed in section 1; the U-gauge allows us to simplify the matter essentially compared to the choice of [5]. Section 2 is devoted to discussion of the UV-divergences originating from the one-loop Feynman graphs renormalising the relevant 2-point Green functions. The leading logarithmic term is recovered in the third section. Most of the technical details are deferred to Appendices.

2 Definition of the renormalised parameters

As was stated before the whole computation will be performed in the unitary gauge; this particular choice reduces the number of diagrams to be considered and simplifies the renormalisation scheme. On the other hand, the presence of Goldstone bosons in \( R_\xi \) gauges causes for example the total cancellation of contributions descending from the tadpole diagrams (see [5]), which does not occur in U-gauge.

The diagrams to be considered are listed in Appendix A. The ultraviolet divergences coming from the loops are handled using the standard technique of dimensional regularisation, see [10].

2.1 Renormalized (pseudo)scalar masses:

The renormalized masses \( m_X \) of (pseudo)scalars \( h, H \) and \( A \) are defined generi-
cally by
\[(1 + \delta Z_X) m_{XB}^2 = m_X^2 + \delta m_X^2; \quad (2)\]
Next, let us denote the sum of all (relevant) graphs by \(-i\Pi_X(q) = \sum -i\Pi_X^I(q)\). Then the 2-point Green functions can be written in the form (see [11]).
\[i\Gamma^{(2)}_X(q, -q) = q^2 - m_X^2 - i\Pi_X(q) + i (\delta Z_X q^2 - \delta m_X^2) + \text{higher order}\]
We adopt the so-called on-shell renormalisation scheme in which all the external momenta (q’s) are taken to be on the mass-shell, i.e. \(q^2 = m_X^2\) where \(m_X^2\) denotes the squared mass of the considered particle. In this scheme we use the following renormalisation conditions
\[\Gamma^{(2)}_X(q, -q) = 0, \quad \frac{\partial \Gamma^{(2)}_X}{\partial q^2}(q, -q) = 1 \quad \text{at} \quad q^2 = m_X^2\]
This particular choice implies
\[\delta Z_X = 0 + \text{higher order}; \quad \delta m_X^2 = -\Pi_X(q^2 = m_X^2) + \text{higher order} \quad (3)\]
The physical mass \(m_X\) can then be expressed (using (2) and (3)) as
\[m_X^2 = m_{BX}^2 + \Pi_X(q^2 = m_X^2) + \text{higher order}\]
### 2.2 Renormalized Z-boson mass:

Let us denote the sum of all the one-loop (Z) ‘vacuum polarisation graphs’ by \(-i\Pi^\mu_\nu_Z(q)\). This quantity renormalizes the Z-boson mass to the new value
\[m_Z^2 = m_{ZB}^2 - A_Z(q^2 = m_Z^2) + \text{higher order}\]
(we have again used the on-shell conditions) where \(A_Z(q^2)\) is defined by
\[\Pi^\mu_\nu_Z(q^2) \equiv A_Z(q^2) g^{\mu\nu} + B_Z(q^2) q^\mu q^\nu\]
(i.e. corresponds to the coefficient of the transverse part of \(\Pi^\mu_\nu_Z\)).

### 2.3 Renormalised sum-rule:

Having defined renormalised quantities we can recast the relation (1) in the renormalised form
\[m_h^2 + m_H^2 - m_A^2 - m_Z^2 = \Delta + \text{higher order}\]
where
\[\Delta \equiv \Pi_h(q^2 = m_h^2) + \Pi_H(q^2 = m_H^2) - \Pi_A(q^2 = m_A^2) + A_Z(q^2 = m_Z^2) \quad (4)\]
This is the most important relation of this section. In the following part we attempt to evaluate the one-loop leading-log term of $\Delta$. As was already stated before the leading term descends from the graphs involving top and supertop loops so the rest of this computation will be performed for this sector only. Note that the full quantity includes contributions from almost all the particles in the theory, which would complicate the calculation essentially without any impact on the leading term so the other contributions are simply omitted.

3 Cancellation of UV-divergences

In this section we show that the UV-divergent parts of the diagrams listed in Appendix A cancel. To proceed we put the external momenta on-shell and substitute in (4). For the sake of brevity there will be no difference between the symbols used for the divergent parts of the considered expressions and the full contributions in this section; moreover the overall factors $C_{uv}$ and $N_c$ are suppressed too. For example $B_t^h$ (see Appendix A) corresponds here to $g_h^2 (4\pi)^{-1} (3m_t^2 - \frac{1}{2}m_h^2)$. To simplify the reader’s insight the definitions of the partial sums of divergences (denoted by UV with relevant sub- and superscripts) take care of the sign of the corresponding expressions in (4).

3.1 UV divergences in graphs involving top loops

Let us start with the divergences descending from the graphs involving one top-quark loop. The first three graphs in (A.1) give

$$B_t^h + B_t^H - B_t^A =$$

$$= \frac{g^2 m_t^2}{16\pi^2 m_W^2 \sin^2 \beta} \left[ 3m_t^2 - m_t^2 \cos^2 \beta + \frac{1}{2} \left( -m_t^2 \cos^2 \alpha - m_h^2 \sin^2 \alpha + m_A^2 \cos^2 \beta \right) \right]$$

Next, the divergent part of the fourth graph in (A.1) contributing to (4) is after some algebra

$$B_Z^t = -\frac{g^2 m_t^2 m_Z^2}{32\pi^2 m_W^2} + \frac{g^2 m_Z^4}{24\pi^2 m_W^2} \left( \epsilon_L^2 + \epsilon_R^2 \right)$$

Utilising relations (36) of Appendix B the tadpole graphs in (A.1) give

$$-T_A^{ht} - T_A^{Ht} = -\frac{g^2 m_t^4}{16\pi^2 m_W^2}$$

Summing up the partial results (5)-(6) one obtains the total divergence coming from the graphs involving one top-quark loop:

$$\text{UV}_{\text{top}} = \frac{g^2}{16\pi^2 m_W^2 \sin^2 \beta} \left[ 2m_t^4 - m_t^2 m_Z^2 \sin^2 \beta + \frac{2}{3} m_Z^2 \sin^2 \beta \left( \epsilon_L^2 + \epsilon_R^2 \right) \right]$$
3.2 UV divergences in graphs involving supertop loops

The same brief list of divergences will be now built up for the diagrams with one supertop loop. The first type graphs in (A.2) give

\[
\text{UV}^{(1)}_{\text{stop}} \equiv B^{\tilde{t}_1 \tilde{t}_1} + B^{\tilde{t}_2 \tilde{t}_2} + B^{\tilde{t}_1 \tilde{t}_1} + B^{\tilde{t}_2 \tilde{t}_2} + B^{\tilde{t}_1 \tilde{t}_2} - B_A^{\text{tadpole}} =
\]

\[
\frac{1}{16\pi^2} \left( g_{\text{h}_1 \text{h}_1} + g_{\text{h}_2 \text{h}_2} + 2g_{\text{h}_1 \text{h}_1} + g_{\text{H}_1 \text{H}_1} + 2g_{\text{H}_2 \text{H}_2} + 2g_{\text{A}_1 \text{A}_1} \right)
\]

The result of the computation is

\[
\text{UV}^{(1)}_{\text{stop}} = \frac{g^2}{16\pi^2 m_W^2 \sin^2 \beta} \left[ -2m_t^4 - \frac{1}{2} m_t^2 (A_t m_6 \sin \beta + \mu \cos \beta)^2 + m_t^2 m_Z^2 \sin^2 \beta - m_2^4 \sin^2 \beta \left( \varepsilon_{L}^2 + \varepsilon_{R}^2 \right) \right]
\]

(8)

Next, the total contribution descending from the second type graphs in (A.2) reads

\[
\text{UV}^{(2)}_{\text{stop}} \equiv B^{\tilde{t}_1 \tilde{t}_2} + B^{\tilde{t}_2 \tilde{t}_2} + B^{\tilde{t}_1 \tilde{t}_2} = \frac{g^2 m_Z^2}{16\pi^2 m_W^2} \left( \frac{1}{3} m_Z^2 \left( \varepsilon_{L}^2 + \varepsilon_{R}^2 \right) - \right.
\]

\[-2 \left[ m_t^2 \left( \varepsilon_{L}^2 \cos^2 \theta_t + \varepsilon_{R}^2 \sin^2 \theta_t \right) + m_t^2 \left( \varepsilon_{L}^2 \sin^2 \theta_t + \varepsilon_{R}^2 \cos^2 \theta_t \right) \right] \]

(9)

The divergence coming from the tadpole sector of (A.2) can be written in the form (note that the minus sign corresponds to the sign of \( \Pi_A \) in (4))

\[
\text{UV}^{(3)}_{\text{stop}} \equiv -T_{h_1}^2 - T_{A}^2 - T_{A}^2 - T_{A}^2
\]

with the result

\[
\text{UV}^{(3)}_{\text{stop}} = \frac{g^2}{32\pi^2 m_W^2} \left\{ m_t^2 \left( m_t^2 + m_t^2 \right) + m_t^2 \left( A_t m_6 + \mu \cot \beta \right)^2 + m_2^2 \cos 2\beta \left[ m_t^2 \left( \varepsilon_{L}^2 \cos^2 \theta_t - \varepsilon_{R}^2 \sin^2 \theta_t \right) + m_t^2 \left( \varepsilon_{L}^2 \sin^2 \theta_t - \varepsilon_{R}^2 \cos^2 \theta_t \right) \right] \right\}
\]

(10)

The last part of the total UV divergence originates from the seagull-type diagrams in (A.2):

\[
\text{UV}^{(4)}_{\text{stop}} \equiv S_{h_1}^2 + S_{h}^2 + S_{A}^2 + S_{A}^2 - S_A^2 - S_A^2 + S_Z^2 + S_Z^2
\]

After some algebra one gets

\[
\text{UV}^{(4)}_{\text{stop}} = - \frac{g^2 m_Z^2}{32\pi^2 m_W^2} \left\{ \left( m_t^2 + m_t^2 \right) + m_2^2 \cos 2\beta \left[ m_t^2 \left( \varepsilon_{L}^2 \cos^2 \theta_t - \varepsilon_{R}^2 \sin^2 \theta_t \right) + m_t^2 \left( \varepsilon_{L}^2 \sin^2 \theta_t - \varepsilon_{R}^2 \cos^2 \theta_t \right) \right] \right\}
\]

5
- $4m_Z^2 \left[ m_{\tilde{t}_1}^2 \left( \varepsilon_L^2 \cos^2 \theta_t + \varepsilon_R^2 \sin^2 \theta_t \right) + m_{\tilde{t}_2}^2 \left( \varepsilon_L^2 \sin^2 \theta_t + \varepsilon_R^2 \cos^2 \theta_t \right) \right] \right\} \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) \right) 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4.1 Contributions of magnitude $m_{t_i}^2$

Looking at the coupling constants in Appendix A and taking into account the relation (39) from Appendix B one can check that the only contributions proportional to $m_{t_1}^2$, $m_{t_2}^2$ come from the second, third and fourth type graphs in (A.2). The relevant expression is defined by

\[
F_{m_{t_i}^2} = \left[ B_{Z}^{i_2} + B_{A}^{i_2} - T_{A}^{h_1} - T_{A}^{h_2} - T_{A}^{H_1} - T_{A}^{H_2} + + S_{h_1}^{i_2} + S_{h_2}^{i_2} - S_{A}^{i_2} - S_{A}^{i_2} + S_{Z}^{i_2} + S_{Z}^{i_2} \right]_{m_{t_i}^2 \text{ only}} (12)
\]

After some manipulations one obtains

\[
F_{m_{t_i}^2} = \frac{N_c g^2 m_Z^2}{16 \pi^2 m_W^2} \left\{ 2 \left( \varepsilon_L^2 \cos^2 \theta_t + \varepsilon_R^2 \sin^2 \theta_t \right) m_{t_1}^2 \left( 1 - \ln m_{t_1}^2 \right) + + 2 \left( \varepsilon_L^2 \sin^2 \theta_t + \varepsilon_R^2 \cos^2 \theta_t \right) m_{t_2}^2 \left( 1 - \ln m_{t_2}^2 \right) - - (\varepsilon_L' \cos^2 \theta_t + \varepsilon_R' \sin^2 \theta_t) \right\} \left( 2m_{t_1}^2 - 2 \int_0^1 dx D_{m_{t_1}^2}^{i_1} (x) \ln D_{m_{t_1}^2}^{i_1} (x) \right) - - (\varepsilon_L' \sin^2 \theta_t + \varepsilon_R' \cos^2 \theta_t) \left( 2m_{t_2}^2 - 2 \int_0^1 dx D_{m_{t_2}^2}^{i_2} (x) \ln D_{m_{t_2}^2}^{i_2} (x) \right) - - \frac{1}{2} \sin^2 \theta_t \left( \varepsilon_L - \varepsilon_R \right)^2 \left( m_{t_1}^2 + m_{t_2}^2 - 2 \int_0^1 dx D_{m_{t_1}^2}^{i_1} (x) \ln D_{m_{t_1}^2}^{i_1} (x) \right) \right\} - - \frac{N_c g^2 m_Z^2}{32 \pi^2 m_W^2 \sin^2 \beta} (A_t m_6 \sin \beta + \mu \cos \beta) \int_0^1 dx \ln D_{m_{t_2}^2}^{i_2} (x) (13)
\]

Note that the last term must indeed be taken into account here because of (39). Fortunately it is strongly suppressed by the assumption of a small mixing in the supertop sector, see Appendix B.

First, it can be checked immediately that the terms of the form const. $\times m_{t_i}^2$ exactly cancel. The remaining structure is already not so easy to handle. The assumption of a small mixing allows us to approximate $\sin \theta_t \sim 0$, which drops out of the penultimate term in (13). The key role of this observation consists in the fact that the rest of (13) already does not involve any mixed $D_{m_{t_1}^2}^{i_1} (x)$ term. Thus the expressions proportional to $m_{t_1}^2$ and $m_{t_2}^2$ split into two independent clusters. To proceed one can use the expansion

\[
\ln D_{m_{t_2}^2}^{i_2} (x) = \ln m_{t_1}^2 + \ln \left[ 1 - \frac{m_{t_2}^2}{m_{t_1}^2} x (1 - x) \right] = \ln m_{t_1}^2 - \frac{m_{t_2}^2}{m_{t_1}^2} x (1 - x) + O \left( \frac{m_{t_2}^4}{m_{t_1}^4} \right) (14)
\]
which originates from the definition of $D_{m_Z}^{ii}$, see Appendix A. Neglecting the contributions proportional to $m_Z^2$ one can check that the terms of the type $m^2_{ti} \ln m^2_{ti}$ cancel too. The previous discussion leads to the following result:

- In the case of no significant mixing within the supertop sector the contribution of the magnitude $m^2_{ti}$ turns out to be negligible compared to the correction proportional to $m^2_{ti} m^4_Z$, which is investigated in the next subsection.

If the mixing in the supertop sector is not negligible one can obtain large negative contribution proportional to $m^2_{ti}$ from this cluster of diagrams; for more comprehensive discussion see [5].

4.2 Contribution proportional to $m^4_t m^{-2}_Z$, leading log-term

Looking into (A.1) and (A.2), one can immediately write down the sum of relevant terms:

$$ \mathcal{F}_{m^4_t \times M^{-2}} = \left[ B_h^t + B_H^t - B_A^t - T_A^{tt} - T_H^{tt} + B_{h1}^{t1} + B_{h2}^{t2} + B_{H1}^{t1} + B_{H2}^{t2} \right] m^4_t \times M^{-2} \text{ only} \quad (15) $$

Note that all the irrelevant parts of orders $m^2_{ti}$, $m^2_Z m^2_h$ and $m^2_{th}$ were neglected; the possibly large factor $\sim m^2_A m^2_t \times M^{-2}$ is suppressed by $\cos \beta$. The other factors like $\sim m^2_h m^2_t \times M^{-2}$ or $\sim m^2_H m^2_t \times M^{-2}$ are assumed not to be high above $\sim m^2_Z m^2_t \times M^{-2}$; moreover they are put down by an overall factor coming
from the integration over $x$. The magnitude of the total error is approximately 10%. In addition it is easy to see that the non-logarithmic factors cancel.

The situation in the supertop-loop cluster (corresponding to the last four terms in (15)) is again quite complicated due to the structure of the appropriate squares of the couplings (28). Extracting only the relevant parts involving $m_t^4$ one gets

$$
\frac{\Gamma^{\text{stop}}_{m_t^4 \times M^2}}{m_t^4} = \frac{N_c g^2 m_t^4}{16 \pi^2 m_W^2 \sin^2 \beta} \left\{ \cos^2 \alpha \left[ \int_0^1 dx \ln D_{\tilde{t}_1 \tilde{t}_1}^{\tilde{t}_1 \tilde{t}_1}(x) + \int_0^1 dx \ln D_{\tilde{t}_2 \tilde{t}_2}^{\tilde{t}_2 \tilde{t}_2}(x) \right] + \\
+ \sin^2 \alpha \left[ \int_0^1 dx \ln D_{H}^{\tilde{t}_1 \tilde{t}_1}(x) + \int_0^1 dx \ln D_{H}^{\tilde{t}_2 \tilde{t}_2}(x) \right] \right\} \quad (17)
$$

The last thing to be done is to apply expansion similar to (14) on the $D$’s in the integrals. Summing up the particular results above it is straightforward to obtain the leading logarithmic term ($N_c = 3$)

$$
\Delta = \frac{3 g^2 m_t^4}{16 \pi^2 m_W^2 \sin^2 \beta} \ln \left( \frac{m_{\tilde{t}_1}^2 m_{\tilde{t}_2}^2}{m_t^4} \right) + \ldots \quad (18)
$$

This is the main result of the whole computation. It shows that including the leading logarithmic one-loop correction the original tree-level relation (1) can be recast in the form

$$
m_h^2 + m_H^2 - m_A^2 - m_Z^2 = \frac{3 g^2 m_t^4}{16 \pi^2 m_W^2 \sin^2 \beta} \ln \left( \frac{m_{\tilde{t}_1}^2 m_{\tilde{t}_2}^2}{m_t^4} \right) + \ldots
$$

This relation agrees with the results presented in the literature ([5], [7]).

Note that the relative error of the approximations used to derive the previous relation does not exceed 10%. It is mainly due to neglecting all the terms of order $m_t^4$ and lower. Next, the form of the leading term (18) is invalid in case that any significant mixing in the stop sector occurs.

**Conclusion**

The paper is devoted to one of the most important features of the SUSY-theories – the total cancellation of ultraviolet divergences in the relations originating from the supersymmetry. Although it is expected to be so in general, there is still no explicit proof based on the Ward identities. Therefore it is convenient to demonstrate this mechanism at least in a particular case of the MSSM tree-level relation $m_h^2 + m_H^2 - m_A^2 - m_Z^2 = 0$. The rule is renormalised using the usual diagrammatic technique. The only considered one-loop graphs are those involving top and supertop loops because the expected magnitude of the correction is the
largest (taking into account the observation of \cite{7} that no significant contribution descends from the chargino-neutralino sector). The original results of \cite{5} and \cite{6} are recovered performing the whole computation in the unitary gauge; the explicit mechanism of the divergence cancellation is shown. The finite part is discussed in detail in the case of no mixing in the supertop sector.

A Relevant Feynman graphs

This Appendix contains all the graphs discussed in previous sections. They are divided into two main subgroups - diagrams with quarks in loops and diagrams with the corresponding SUSY-partners. Each group consists of several types of graphs; the notation is usual and (perhaps) self-explanatory. The special symbols are defined as follows:

- $C_{UV} \equiv \varepsilon^{-1} - \gamma_e + \ln 4\pi$ denotes the "divergent" part of a graph; here $2\varepsilon = 4 - d$; $d$ is the noninteger dimension used in the dimensional regularisation procedure; $\gamma_e$ is the Euler-Mascheroni constant.
- $B^{f_1 f_2 \ldots}_X$, $T^{f_1 f_2 \ldots}_X$ and $S^{f_1 f_2 \ldots}_X$ denote self-energies descending from the Feynman graphs (usually called blobs, tadpoles and seagulls) with external lines $X$ and internal $f_1, f_2, \ldots$
- $D^{f_1 f_2}(x) \equiv m^2_{f_1} (1 - x) + m^2_{f_2} x - q^2 x (1 - x)$ is the common factor arising from the regularisation prescription (see \cite{10}) for UV-divergent graphs; $q$ is the momentum of the incoming (and outgoing) particle; in the on-shell scheme $q^2 = m^2_X$.
- The constants $g_{f_1 f_2 f_3 \ldots}$ denote the numerical parts of the corresponding vertices. The non-number parts are contained in the structure of integrands.
- For the sake of brevity the overall colour factor $N_c = 3$ is suppressed but must be included to obtain the correct results.

A.1 Graphs with top quark loops

I. Scalar and pseudoscalar self-energies:

\[
B^t_X(q) = -i g^2_{\chi}\mu^{2\varepsilon} \int \frac{d^d k}{(2\pi)^d} \text{Tr} \frac{i}{k - m_t} \frac{i}{\not k - \not q - m_t}; \quad X = h, H
\]

\[
B^{t}_A(q) = -i g^2_{\chi}\mu^{2\varepsilon} \int \frac{d^d k}{(2\pi)^d} \text{Tr} \gamma_5 \frac{i}{k - m_t} \gamma_5 \frac{i}{\not k - \not q - m_t}
\]
Using the routine calculational procedure (see for instance [10]) the results are

\[
B_X(q) = g_{Xtt}^2 \frac{1}{4\pi^2} \left\{ C_{uv} \left(3m_t^2 - \frac{q^2}{2}\right) + \left(m_t^2 - \frac{q^2}{6}\right) + \right.
= \int_0^1 dx \ln \frac{D^{st}(x)}{\mu^2} \left[-3m_t^2 + 3q^2 x(1-x)\right] \right\}; \quad X = h, H
\]

\[
B_A(q) = g_{Att}^2 \frac{1}{4\pi^2} \left\{ C_{uv} \left(-m_t^2 + \frac{q^2}{2}\right) - \left(m_t^2 - \frac{q^2}{6}\right) + \right.
\int_0^1 dx \ln \frac{D^{st}(x)}{\mu^2} \left[m_t^2 - 3q^2 x(1-x)\right] \right\}
\]

The corresponding couplings are

\[
g_{htt} = -ig_{htt} \frac{m_t}{2m_W \sin^\beta}; \quad g_{htt} = -ig_{htt} \frac{m_t}{2m_W \sin^\beta}; \quad g_{Att} = \frac{g_{Att}}{2m_W \cot^\beta}
\]

II. Z-boson self-energy graph (vacuum polarisation tensor):

\[
\begin{array}{c}
Z \\
\overline{\gamma} \\
_t \\
\gamma
\end{array}
\begin{array}{c}
\overline{t} \\
\rightarrow \\
Z
\end{array}
\begin{array}{c}
\gamma \\
\overline{\gamma}
\end{array}
\equiv B_Z(q)^{\mu\nu}
\]

\[
B_Z(q)^{\mu\nu} = -ig_{Ztt}^2 2\pi \int \frac{d^4k}{(2\pi)^4} \frac{i\gamma^\mu K(\gamma)}{k - m_t} \frac{i\gamma^\nu K(\gamma)}{-q - m_t};
\]

where

\[K(\gamma) \equiv (\epsilon_L^t + \epsilon_R^t)I_4 - (\epsilon_L^t - \epsilon_R^t)\gamma_5\]

The constants \(\epsilon_L^t, \epsilon_R^t\) are defined as follows (in general \(\epsilon^t = T_3 - Q_f \sin^2 \theta_W\))

\[
\epsilon_L^t \equiv \frac{1}{2} - q_t \sin^2 \theta_W; \quad \epsilon_R^t \equiv -q_t \sin^2 \theta_W
\]

The coupling \(g_{Ztt}\) is

\[
g_{Ztt} = -ig \frac{2}{\cos \theta_W}
\]

Performing the usual steps the result becomes

\[
B_Z(q)^{\mu\nu} = g_{Ztt}^2 \frac{1}{4\pi^2} \left\{ \int_0^1 dx \ln \frac{D^{tt}(x)}{\mu^2} \left(-\frac{m_t^2}{2}\right) + \right.
\int_0^1 dx \ln \frac{D^{tt}(x)}{\mu^2} \left[\frac{m_t^2}{2}g^{\mu\nu} + \frac{2}{3}(\epsilon_L^t \epsilon_R^t + \epsilon_L^t \epsilon_R^t)(g^{\mu\nu} q^2 - q^\mu q^\nu)\right] + \left.
\int_0^1 dx \ln \frac{D^{tt}(x)}{\mu^2} \left[\frac{m_t^2}{2}g^{\mu\nu} - \frac{2}{3}(\epsilon_L^t \epsilon_R^t - \epsilon_L^t \epsilon_R^t)(q^\mu q^\nu - g^{\mu\nu} q^2)x(1-x)\right] \right\}
\]

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III. Tadpoles involving top quark loop:

In general there are eight graphs to be considered in this paragraph. They are the following (here $X = h, H, A$ and $S = h, H$):

\[
\begin{align*}
XX \equiv T_X^{St}(q) \\
ZZ \equiv T_Z^{St}(q)^{\mu \nu}
\end{align*}
\]

Fortunately many of them cancel because of a nice property of the corresponding couplings "sitting" in the upper vertex

\[
ghhh + gHHh + gZZh = 0 \\
ghhH + gHHH + gZZH = 0
\] (24)

The remaining graphs are evaluated in the usual way:

\[
T_A^{St}(q) = ig_{AAS}g_{ltS} \frac{i}{m_S^2} \mu^{2\varepsilon} \int \frac{d^d \ell}{(2\pi)^d} \text{Tr} \frac{i}{\ell - m_t};
\]

This d-dimensional integration is already easy to handle; we get

\[
T_A^{St}(q) = 4g_{AAS}g_{ltS} \frac{m_t^3}{m_S^2} \frac{1}{16\pi^2} \left( C_{UV} + 1 - \ln \frac{m_t^2}{\mu^2} \right); \quad (25)
\]

The couplings $g_{htt}$ and $g_{hht}$ are already written in (24); the remaining constants are

\[
\begin{align*}
g_{AAh} &= -ig_{Z} m_Z \cos2\beta \sin(\alpha + \beta) ; & g_{AAH} &= ig_{Z} m_Z \cos2\beta \cos(\alpha + \beta)
\end{align*}
\] (26)

At the end of this subsection note that the coupling constants used in this article can be found for example in [1] and [7]. In the case of supertops one must transform the rules in [1] from the $L - R$ basis to the supertop mass-diagonal basis $1 - 2$; this procedure is well described in the cited paper.
A.2 Graphs with supertop loops

I. Diagrams of the first type:

There are 7 graphs to be investigated in this category, namely

\[ B^X_{ij}(q) = g_{X_i,j} g_{X_j,i} \mu^{2\epsilon} \int \frac{d^d k}{(2\pi)^d} \frac{i i}{k^2 - m_i^2} \frac{i j}{(k - q)^2 - m_j^2} (2 - \delta_{ij}) \]

Using again (10), the last expression can be simplified to the final form

\[ B^X_{ij}(q) = g_{X_i,j} g_{X_j,i} (2 - \delta_{ij}) \frac{1}{16\pi^2} \left[ C_{UV} - \int_0^1 dx \ln \frac{D^{X,ij}_q(x)}{\mu^2} \right] \quad (27) \]

Note that the factor \((2 - \delta_{ij})\) counts the number of nonequivalent contractions. To finish this paragraph it is necessary to specify the couplings; the vertices involving scalars are symmetric with respect to \(i \leftrightarrow j\)

\[
\begin{align*}
\frac{g_{hi1i1}}{\cos \theta_W} &= \frac{igm_z}{\cos \theta_W} \sin(\alpha + \beta) \left( \epsilon_i^L \cos^2 \theta_t - \epsilon_i^R \sin^2 \theta_t \right) - \frac{igm_z^2 \cos \alpha}{m_W \sin \beta} \\
&\quad - \frac{igm_z \sin 2\theta_t}{2m_W \sin \beta} \left( A_i m_\epsilon \cos \alpha - \mu \sin \alpha \right) \\
\frac{g_{hi2i2}}{\cos \theta_W} &= \frac{igm_z}{\cos \theta_W} \sin(\alpha + \beta) \left( \epsilon_i^L \sin^2 \theta_t - \epsilon_i^R \cos^2 \theta_t \right) - \frac{igm_z^2 \cos \alpha}{m_W \sin \beta} \\
&\quad + \frac{igm_z \sin 2\theta_t}{2m_W \sin \beta} \left( A_i m_\epsilon \cos \alpha - \mu \sin \alpha \right) \\
\frac{g_{hi1i2}}{\cos \theta_W} &= -\frac{igm_z}{\cos \theta_W} \sin(\alpha + \beta) \left( \epsilon_i^L + \epsilon_i^R \right) \sin \theta_t \cos \theta_t \\
&\quad - \frac{igm_z \cos 2\theta_t}{2m_W \sin \beta} \left( A_i m_\epsilon \cos \alpha - \mu \sin \alpha \right) \quad (28) \\
\frac{g_{Hi1i1}}{\cos \theta_W} &= -\frac{igm_z}{\cos \theta_W} \cos(\alpha + \beta) \left( \epsilon_i^L \cos^2 \theta_t - \epsilon_i^R \sin^2 \theta_t \right) - \frac{igm_z^2 \sin \alpha}{m_W \sin \beta} \\
&\quad - \frac{igm_z \sin 2\theta_t}{2m_W \sin \beta} \left( A_i m_\epsilon \sin \alpha + \mu \cos \alpha \right) \\
\frac{g_{Hi2i2}}{\cos \theta_W} &= -\frac{igm_z}{\cos \theta_W} \cos(\alpha + \beta) \left( \epsilon_i^L \sin^2 \theta_t - \epsilon_i^R \cos^2 \theta_t \right) - \frac{igm_z^2 \sin \alpha}{m_W \sin \beta} \\
&\quad + \frac{igm_z \sin 2\theta_t}{2m_W \sin \beta} \left( A_i m_\epsilon \sin \alpha + \mu \cos \alpha \right)
\end{align*}
\]
\[ g_{HT_i t_2} = \frac{igm_Z}{\cos \theta_W} \cos(\alpha + \beta) (\varepsilon^L_L + \varepsilon^R_R) \sin \theta t \cos \theta t \]
\[ -\frac{igm_t \sin 2\theta_t}{2m_W \sin \beta} (A_t m_6 \sin \alpha + \mu \cos \alpha) \]

while vertices with pseudoscalar \( A \) are antisymmetric:
\[ g_{At_i t_1} = -g_{At_2 t_1} = \frac{gm_t}{2m_W \sin \beta} (A_t m_6 \cos \beta - \mu \sin \beta) \]
\[ g_{At_i t_2} = -g_{At_2 t_2} = 0 \quad (29) \]

II. Z-boson self-energy graphs with looping superquarks:

The graphs relevant to this paragraph are all of the type
\[
\begin{array}{c}
\text{Z} \quad \tilde{t}_i \\
\text{Z} \quad \tilde{t}_j
\end{array}
\]

The graphs relevant to this paragraph are all of the type
\[
B_{Z}^{\tilde{t}_i \tilde{t}_j}(q)^{\mu \nu}; \quad j = 1, 2
\]

As in the previous cases, after some algebra one obtains
\[
B_{Z}^{\tilde{t}_i \tilde{t}_j}(q)^{\mu \nu} = ig_{Z}^{2} m_{t_1} \mu^{2 \epsilon} \int \frac{d^d k}{(2 \pi)^d} \frac{i (2k - p)^{\mu}}{k^2 - m_{t_1}^2} \frac{i (2k - p)^{\nu}}{(k - q)^2 - m_{t_2}^2}
\]

As in the previous cases, after some algebra one obtains
\[
B_{Z}^{\tilde{t}_i \tilde{t}_j}(q)^{\mu \nu} = -g_{Z}^{2} m_{t_1} \mu^{2 \epsilon} \left\{ C_{UV} \left[ \frac{1}{3} (q^{\mu} q^{\nu} - g^{\mu \nu} q^{2}) + (m_{t_1}^2 + m_{t_2}^2) g^{\mu \nu} \right] + \right.
\]
\[+ g^{\mu \nu} \left( m_{t_1}^2 + m_{t_2}^2 - \frac{q^2}{3} \right) - \int_{0}^{1} dx \left[ q^{\mu} q^{\nu} (1 - 2x)^2 + 2g^{\mu \nu} D_{q}^{\tilde{t}_i \tilde{t}_j}(x) \right] \ln \left( \frac{D_{q}^{\tilde{t}_i \tilde{t}_j}(x)}{\mu^2} \right) \left\} \right.
\]

III. Tadpoles involving supertop loop:

In general there is again many diagrams belonging to this paragraph; as in the previous section the relations (24)-(24) ensure that most of the graphs cancel.

The remaining are (here \( S = h, H \) and \( i = 1, 2 \)):
The contributions coming from these graphs are
\[ T_A^{S\bar{t}_i}(q) = -ig_AAS\bar{g}_i\bar{t}_iS \frac{i}{m_S^2} \mu^{2\varepsilon} \int \frac{d^d k}{(2\pi)^d} \frac{i}{k^2 - m_{\tilde{t}_i}^2}; \]
which gives after regularisation
\[ T_A^{S\bar{t}_i}(q) = -g_AAS\bar{g}_i\bar{t}_iS \frac{m_{\tilde{t}_i}^2}{m_S^2} \frac{1}{16\pi^2} \left( C_{\text{UV}} + 1 - \ln \frac{m_{\tilde{t}_i}^2}{\mu^2} \right); \quad (30) \]
The necessary coupling constants are written in (26) and (28).

**IV. Seagull graphs:**

Due to the presence of quadrilinear vertices involving two Higgses and two superquarks there is an additional sort of graphs in this section – the so-called seagull graphs which look as \((S = h,H,A \text{ and } i = 1,2)\)
\[ S_{\tilde{t}_i} \equiv S_{\tilde{t}_i}^{S}(q); \quad Z_{\tilde{t}_i} \equiv S_{\tilde{t}_i}^{Z}(q)^{\mu\nu} \]

Note that similar graphs can not appear in the fermion sector because the quadrilinear vertices involving fermions have mass dimensions > 4 and they would cause nonrenormalisability of the theory. The contributions originating from these graphs can be written as
\[ S_{\tilde{t}_i}^{S}(q) = ig_{SS\tilde{t}_i\tilde{t}_i} \mu^{2\varepsilon} \int \frac{d^d k}{(2\pi)^d} \frac{i}{k^2 - m_{\tilde{t}_i}^2}; \]
\[ S_{\tilde{t}_i}^{Z}(q)^{\mu\nu} = ig_{ZZ\tilde{t}_i\tilde{t}_i} \mu^{2\varepsilon} g^{\mu\nu} \int \frac{d^d k}{(2\pi)^d} \frac{i}{k^2 - m_{\tilde{t}_i}^2}; \]
which after regularisation gives
\[ S_{\tilde{t}_i}^{S}(q) = -g_{SS\tilde{t}_i\tilde{t}_i} m_{\tilde{t}_i}^2 \frac{1}{16\pi^2} \left( C_{\text{UV}} + 1 - \ln \frac{m_{\tilde{t}_i}^2}{\mu^2} \right); \quad (31) \]
\[ S_{\tilde{t}_i}^{Z}(q)^{\mu\nu} = -g_{ZZ\tilde{t}_i\tilde{t}_i} m_{\tilde{t}_i}^2 g^{\mu\nu} \frac{1}{16\pi^2} \left( C_{\text{UV}} + 1 - \ln \frac{m_{\tilde{t}_i}^2}{\mu^2} \right); \]

In general there are 8 graphs to deal with. Fortunately the coupling constants can be nicely summed up so that
\[ g_{hh\tilde{t}_i\tilde{t}_i} + g_{HH\tilde{t}_i\tilde{t}_i} + g_{hh\tilde{t}_2\tilde{t}_2} + g_{HH\tilde{t}_2\tilde{t}_2} - g_{AA\tilde{t}_i\tilde{t}_i} - g_{AA\tilde{t}_2\tilde{t}_2} = g_{GG\tilde{t}_2\tilde{t}_2} + g_{GG\tilde{t}_1\tilde{t}_1}; \quad (32) \]
where

\[ g_{GG\tilde{t}_1\tilde{t}_1} = -\frac{ig^2}{2\cos^2\theta_W} \cos 2\beta \left( \varepsilon^L_1 \cos^2\theta_t - \varepsilon^R_1 \sin^2\theta_t \right) - \frac{ig^2 m^2_t}{2m^2_W} \]

\[ g_{GG\tilde{t}_2\tilde{t}_2} = -\frac{ig^2}{2\cos^2\theta_W} \cos 2\beta \left( \varepsilon^L_1 \sin^2\theta_t - \varepsilon^R_1 \cos^2\theta_t \right) - \frac{ig^2 m^2_t}{2m^2_W} \] (33)

Note that these constants are exactly the couplings of the would-be Goldstone boson (which is within U-gauge absent). The last unspecified parameters are the couplings \( g_{ZZ\tilde{t}_1\tilde{t}_1} \):

\[ g_{ZZ\tilde{t}_1\tilde{t}_1} = \frac{2ig^2}{\cos^2\theta_W} \left( \varepsilon^L_1 \cos^2\theta_t + \varepsilon^R_1 \sin^2\theta_t \right) \]

\[ g_{ZZ\tilde{t}_2\tilde{t}_2} = \frac{2ig^2}{\cos^2\theta_W} \left( \varepsilon^L_1 \sin^2\theta_t + \varepsilon^R_1 \cos^2\theta_t \right) \] (34)

B Some useful relations and comments

This Appendix is devoted to several clarifications necessary to make the text more self-consistent.

The first note refers to the MSSM itself. The full-range discussion of the relevant part of the MSSM classical lagrangian is obviously out of the scope of this paper. There are several comprehensive works in the literature that can be used for this purpose, namely [1], [2], [11] or [12]. The notation is similar to that used in [1].

The rest of the Appendix contains some comments on technical details of the computation. To enable the reader follow the steps described above it is necessary to write down several not so well known tree-relations often used during the computation. First of them,

\[ m^2_h \cos^2\alpha + m^2_H \sin^2\alpha - m^2_A \cos^2\beta - m^2_Z \sin^2\beta = 0 \] (35)

can be derived by utilising (1) and relations

\[ \sin 2\alpha = \left( \frac{m^2_h + m^2_H}{m^2_h - m^2_H} \right) \sin 2\beta; \quad \cos 2\alpha = \left( \frac{m^2_A - m^2_Z}{m^2_h - m^2_H} \right) \cos 2\beta \]

that can be found for example in [2]. Note only that the parameters \( \alpha \) and \( \beta \) are the mixing angles in the scalar and pseudoscalar parts of the Higgs sector. These relations are also very handy if we want to express (tree) Higgs masses in terms of \( m^2_h \), \( \alpha \) and \( \beta \) dealing with the factors \( m^2_h \) or \( m^2_H \) coming from the tadpoles in (A.1) and (A.2).

The next thing to be clarified is the role of the parameters \( A_t m_6 \) and \( \mu \) in the supertop mass-squared matrix. This matrix in the \( L - R \) basis looks (see
\[
M^2_{t_{L,R}} = \begin{pmatrix}
A & B \\
B & C
\end{pmatrix}
\]  \hspace{1cm} (37)

where

\[
A = M_3^2 + m_2^2 \cos 2\beta \epsilon_L^t + m_t^2 \\
B = m_t (A_t m_6 + \mu \cot \beta) \\
C = M_3^2 + m_2^2 \cos 2\beta \epsilon_R^t + m_t^2
\]

is the usual parametrisation of its entries. (Note that the constants \(M_3^2\) and \(M_3^2\) are the so-called soft-SUSY breaking terms which in general split the masses of supertops and shift them high above \(m_t\).) The eigenvalues of this matrix can be easily derived in the form

\[
m^2_{\tilde{t}_{1,2}} = \frac{1}{2} \left[ A + C \pm \sqrt{(A+C)^2 - 4(AC-B^2)} \right]  \hspace{1cm} (38)
\]

The mixing (diagonalising) angle \(\theta_t\) is then defined by

\[
\tan \theta_t = \frac{2B}{A-C}
\]

which can be rewritten in terms of the eigenvalues (38) as follows

\[
\sin 2\theta_t = \frac{2m_t (A_t m_6 + \mu \cot \beta)}{m^2_{\tilde{t}_1} - m^2_{\tilde{t}_2}}  \hspace{1cm} (39)
\]

This relation connects the off-diagonal entries in the supertop mass-squared matrix with the magnitude of the supertop mass-split and the mixing angle \(\theta_t\). Assuming now that the supertop mass-squared does not exceed the top-scale too much the assumption of a small mixing can be recast in the form \((A_t m_6 + \mu \cot \beta) \ll m_t\).

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