Higher Dimensional Operators in the MSSM

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Abstract. The origin and the implications of higher dimensional effective operators in 4-dimensional theories are discussed in non-supersymmetric and supersymmetric cases. Particular attention is paid to the role of general, derivative-dependent field redefinitions which one can employ to obtain a simpler form of the effective Lagrangian. An application is provided for the Minimal Supersymmetric Standard Model extended with dimension-five R-parity conserving operators, to identify the minimal irreducible set of such operators after supersymmetry breaking. Among the physical consequences of this set of operators are the presence of corrections to the MSSM Higgs sector and the generation of “wrong”-Higgs Yukawa couplings and fermion-fermion-scalar-scalar interactions. These couplings have implications for supersymmetry searches at the LHC.

Keywords: Higher dimensional operators, supergravity and supersymmetry breaking, Higgs mass corrections, non-holomorphic couplings

PACS: 11.30.Pb, 12.60.Jv, 12.60.-i, 12.60.Fr, 11.15.Tk

Introduction

The Standard Model (SM) and its minimal supersymmetric version (MSSM) are the best models we currently have for describing the low-energy physics. Despite their success, there are many reasons to think that they are only a low-energy manifestation of a more fundamental theory (supergravity, strings, extra dimensions, etc), valid at higher scales. However, the exact details of such a fundamental theory are in many cases unknown (moduli problem, vacua degeneracy, etc) and making definite predictions for new physics is difficult. One possibility to investigate new physics beyond the SM/MSSM is to use instead an effective field theory approach. This approach is a fully consistent and useful framework for such study.

In effective field theories, operators of dimension larger than four are present, suppressed by the (high) scale of new physics $M_X \gg m_Z$. The origin and the effects of these operators are discussed in this talk, for the case of 4 dimensional non-supersymmetric and supersymmetric theories. We shall distinguish two classes of higher dimensional operators: class A of operators involving at most two space-time derivatives acting on physical fields (one derivative in the case of fermions); class B of operators which contain more than two such derivatives (one in the case of fermions). In general these classes of operators are not entirely independent of each other.

Regarding their origin, higher dimensional operators are generated by integrating out new physics at $M_X \gg m_Z$ ($M_X \sim \text{TeV}$ or higher). In compactification of higher dimensional theories such operators are usually generated, suppressed by the volume of compactification. Much more commonly, higher dimensional operators are generated in 4D renormalisable theories, after integrating out massive states. Therefore, although some interactions may look non-renormalisable in the effective formulation, they may actually be a low energy manifestation of a renormalisable theory valid up to a much higher scale. The familiar Fermi interaction is such an example.

The power of the effective approach resides in arranging these operators in series of powers of $1/M_X$, to which additional organising criteria, such as symmetry arguments inspired by low-energy phenomenology, are also considered. The effective Lagrangian has then the form

$$\mathcal{L} = \mathcal{L}_0 + \sum_{i,n} \frac{c_n^i}{M_X^n} \phi_n^i$$

where $\mathcal{L}_0$ is the SM or the MSSM Lagrangian; $\phi_n^i$ is an operator of dimension $d = n + 4$ with the index $i$ running over the set of operators of a given dimension, and $c_n^i \sim \mathcal{O}(1)$. This description is appropriate for scales $E$ which satisfy $E \ll M_X$. Constraints from phenomenology can then be used to set bounds on the scale of new physics $M_X$. The effects of $\phi_n^i$ on observables can be comparable to one-loop effects in the SM/MSSM, as we shall see in an example, and this shows the importance of their study.

1 This is based on the talks of the authors (I. A. and D.M.G.), presented at the Planck 2008 conference (19-23 May, Barcelona), and SUSY 2008 (16-21 June, Seoul), and it will appear in the proceedings of the SUSY 2008 conference.
The non-supersymmetric case

Let us see some examples of the origin of these operators. Consider first the case of a tree level exchange of a massive $Z'$ gauge boson beyond the SM or MSSM:

$$\mathcal{L} \ni \left| (\partial_\mu - iZ'_\mu) H \right|^2 - \frac{M^2}{2} Z'_\mu Z^\mu \quad (2)$$

After integrating out $Z'$, a higher dimensional operator of class A is generated for instance for the Higgs field $H$, which we denote $\Delta \mathcal{L}'$:

$$\Delta \mathcal{L} = \frac{1}{M^2} (H^\dagger \partial_\mu H)^2 \quad (3)$$

Similarly, for fermions charged under $Z'$, one finds

$$\mathcal{L} \ni i \bar{\psi} \gamma^\mu D_\mu \psi - \frac{M^2}{2} Z'_\mu Z^\mu$$

$$\Rightarrow \Delta \mathcal{L}' = \frac{1}{2M^2} (\bar{\psi} \gamma_\mu \psi)^2 \quad (4)$$

There are also operators of class B which can be generated, by the kinetic mixing of light with massive states, upon integrating out the latter. For example, from

$$\mathcal{L} = \frac{1}{2} (\partial_\mu \phi)^2 + \frac{1}{2} (\partial_\mu \chi)^2 + c \partial_\mu \phi \partial_\mu \chi - \frac{\lambda}{4} \phi^4 - \frac{1}{2} \frac{M^2}{2} \chi^2 - \frac{1}{2} \frac{\lambda'}{2} \phi^2 \chi^2 \quad (5)$$

one finds upon integrating out the massive field $\phi'$:

$$\mathcal{L} = \frac{1}{2} (\partial_\mu \phi)^2 - \frac{\lambda}{4} \phi^4 + \frac{c^2}{2} \frac{M^2}{2} + \frac{1}{2} \lambda' \phi^2$$

$$\Rightarrow \frac{1}{2} (\partial_\mu \phi)^2 - \frac{\lambda}{4} \phi^4 + \frac{c^2}{2} \frac{M^2}{2} (\Box \phi)^2 + \cdots \quad (6)$$

This contains higher dimensional operators of class B (more than two derivatives). If one truncated the series of derivatives to the last term in the second line in (6), after a field redefinition one obtains a formulation of $\mathcal{L}$ which involves only two space-time derivatives but negative metric (ghost) fields [1]. Obviously, this is an artifact of the truncation: there are no ghosts present in the theory, as long as one retains the whole series of expansion in derivatives and provided that the original theory was ghosts-free. We shall generalise this to the supersymmetric case (see [2]).

Higher dimensional operators of class B are also present in the low energy effective action of string theory. One can expand the Dirac-Born-Infeld action, to find an infinite series of such operators. Also $\alpha'$ and loop corrections in string theory generate higher derivative operators. In effective field theories of compactification such operators are generated dynamically at one-loop [3] [4] [5] [6] [7] after integrating out momentum modes.

The supersymmetric case

Consider a general 2-derivative supersymmetric Lagrangian, with the following functions of the chiral superfields $\Phi_i$: the Kähler potential $K$, the superpotential $\mathcal{W}$ and the gauge kinetic function $f$:

$$\mathcal{L} = \int d^4 \theta K(\Phi_i^\dagger e^\lambda, \Phi^\lambda)$$

$$\Rightarrow \int d^2 \theta \left[ \mathcal{W}(\Phi_i) + f_{ab}(\Phi_i) \mathcal{W}^a \mathcal{W}^b \right] + h.c. \quad (7)$$

where $\mathcal{W}^a$ is the supersymmetric gauge field strength associated to the vector superfield $\mathcal{V}^a$. The presence of higher dimensional operators is hidden in the power expansion (in fields) of these functions:

$$K = \Phi_i^\dagger e^\lambda, \Phi^\lambda + \left[ \frac{c_{ijk}}{M_e} \Phi_i^\dagger e^\lambda \Phi^\lambda \Phi^k + h.c. \right] + \cdots$$

$$\mathcal{W} = \lambda_{ijk} \Phi_i^\dagger \Phi^\lambda \Phi^k + \left[ \frac{c_{ijkl}}{M_e} \Phi_i^\dagger \Phi^\lambda \Phi^k \Phi^l + \cdots \right]$$

$$f_{ab} = \delta_{ab} + \frac{f_{ab}}{M_e} \Phi^\lambda + \cdots \quad (8)$$

The first term in therhs would lead to a renormalisable theory. These functions can contain not only operators of class A, but also operators of class B. For example one can have: higher derivative operators in the superpotential $(a)$ and in the Kähler potential $(b)$:

$$\frac{\lambda_{ij}}{M_e} \int d^2 \theta \Phi_i \Box \Phi_j \sim \frac{\lambda_{ij}}{M_e} \int d^4 \theta \Phi_i D^2 \Phi_j$$

$$\frac{k_{ij}}{M_e^2} \int d^2 \theta \Phi_i \Box \Phi_j, \quad \frac{k_{ij}}{M_e^2} \int d^4 \theta \Phi_i \Phi_j D^2 \Phi_j, \cdots \quad (9)$$

where $D$ is the chiral supercovariant derivative. In $(a)$, terms like $\Phi \Box \Phi$ and $F \Box F$ are generated, where $\Phi = \Phi + \sqrt{2} \theta \psi + \partial^2 F$. In $(b)$ one finds terms like $\Box \phi^2$, $\Box \partial_\mu \psi$, $F \Box F$. Therefore for class B operators, auxiliary fields become dynamical degrees of freedom.

Let us present the origin of such operators in a simple case of a 4D supersymmetric renormalisable theory. Consider the Lagrangian

$$\mathcal{L} = \int d^4 \theta \left[ \Phi^\dagger \Phi + \chi \Box \chi \right]$$

$$\Rightarrow \int d^2 \theta \left[ m \Phi \chi + \frac{M^2}{2} \chi^2 + \frac{\lambda}{3} \Phi^3 \right] + h.c. \quad (10)$$

Using the eqs of motion of the massive field $\chi$ and some field redefinitions [8], one obtains

$$\mathcal{L} = \int d^4 \theta \left[ (1 + \frac{m^2}{M^2}) \Phi^\dagger \Phi + \frac{M^2}{2} \Phi^\dagger \Box \Phi + \cdots \right]$$

$$\Rightarrow \frac{m^2}{2M_e} \Phi^\dagger \Phi + \frac{\lambda}{3} \Phi^3 + \frac{m^2}{2M_e^2} \Phi^\dagger \Box \Phi + h.c. \quad (11)$$
If one keeps all terms in the series expansion above, the theory is ghost-free; the effective field theory \([11]\) is valid only below \(M_s\).

From this discussion the following question emerges: is it possible to reformulate a supersymmetric theory with higher dimensional operators of class B, in terms of a theory with operators of class A only (i.e. with at most two derivatives)? As we shall see shortly, the answer is in many cases affirmative \([2]\). Such a reformulation has interesting advantages. The coupling to gravity would become much simpler, and, as a result supersymmetry breaking is easier to study. In particular, the coupling of the supersymmetry breaking sector to the visible sector can be analysed by the usual standard methods. Given the presence of ghost superfields in the action, one could also ask whether such a theory makes sense. The answer is affirmative, as long as one treats the low energy theory as an effective theory, valid at energies \(E \ll M_s\), where \(M_s\) is essentially the mass of the ghost(s).

**Higher dimensional operators in the superpotential**

Let us give an example with operators of class B in the superpotential. Similar considerations apply to when these are present in the Kähler potential \([2]\). Consider

\[
\mathcal{L} = \int d^4 \theta \, \Phi \Phi + \int d^2 \theta \left[ \frac{s}{M_s} \Phi \Phi + \frac{m}{2} \Phi^2 + \frac{\lambda}{3} \Phi^3 \right] + h.c. \tag{12}
\]

with \(s = \pm 1\). A field redefinition of \(\mathcal{L}\), which treats \(\Phi\) and \(\Phi^\dagger \equiv \bar{D} \Phi\) as two superfields of a Lagrangian with constraints (since \(\Phi, \Phi^\dagger\) are not independent), brings \(\mathcal{L}\) to the form \([2]\)

\[
\mathcal{L} = \int d^4 \theta \left[ \Phi \Phi - \Phi^2 \Phi^2 \right] \\
+ \left\{ \int d^2 \theta \left[ \frac{m}{16 \sqrt{\eta}} \left( (1 - \sqrt{\eta}) \Phi - (1 + \sqrt{\eta}) \Phi^2 \right) \right)^2 \\
+ \frac{m}{2 \sqrt{\eta}} (\Phi - \Phi^2)^2 \right\} + h.c. \right\} \tag{13}
\]

where \(\eta = 1 + (17/16) m^2/M^2_s\). For \(m \ll M_s\), \(\eta \rightarrow 1\) and then \([13]\) simplifies considerably. The relation between initial fields and new \(\Phi, \Phi^\dagger\) can be found in \([2]\). This result is easily extended for a general (derivative-free) contribution \(W(\Phi, \chi)\) to the superpotential, present on top of the above class B operator:

\[
\mathcal{L} = \int d^4 \theta \left[ \Phi \Phi + \chi \chi \right] \\
+ \int d^2 \theta \left[ \frac{s}{M_s} \Phi \Phi + W(\Phi, \chi) \right] + h.c. \tag{14}
\]

This can be re-written, if \(m \ll M_s\) \([3]\):

\[
\mathcal{L} = \int d^4 \theta \left[ \Phi \Phi - \Phi^2 \Phi^2 \right] \\
+ \int d^2 \theta \left[ \frac{M_s}{4 s} \Phi^2 + W(\Phi, \chi) \right] + h.c. \tag{15}
\]

These examples show how to “unfold” the original Lagrangian with operators of class B into a form with operators of class A only and additional superfields (\(\Phi^\dagger\)).

In these examples the scalar potential takes the form

\[
V = \sum_{\text{particles}} |F|^2 - \sum_{\text{ghosts}} |F|^2 \tag{16}
\]

The first contribution comes from particles and the second from the ghost degrees of freedom. One can then have \(V > 0\), or \(V < 0\), or even \(V = 0\) with broken supersymmetry. The breaking can be done by a non-trivial auxiliary field expectation value of particle-like \(F\), ghost-like \(F\), or of both types of fields. Consider for example a toy model with explicit soft breaking

\[
\mathcal{L} = \int d^4 \theta \left[ \Phi \Phi - \Phi^2 \Phi^2 \right] \\
+ \int d^2 \theta \left[ \Phi \Phi - \Phi^2 \Phi^2 \right] + h.c. \tag{17}
\]

where \(\phi_{1,2}\) are the scalar fields components of \(\Phi_{1,2}\). The two auxiliary fields have identical eqs of motion, so \(V(\phi_{1,2}) = V_{\text{soft}}(\phi_{1,2})\) and \(V\) has a minimum at \(\phi_1 = \phi_2\). Assuming \(W' \neq 0\), possible if \(W\) contains a linear term \(g(\Phi - \Phi^2)\), then \(F_1 = F_2 = g \neq 0\), so supersymmetry is broken, yet the overall scalar potential is vanishing.

**MSSM with higher dimensional operators**

We consider the extension of the MSSM by higher dimensional operators of class A and B and examine their physical effects \([3]\). Operators of class A and/or class B are generated by integrating out massive superfields which have interactions with light superfields or which mix with them. For example a superpotential with a massive gauge-singlet superfield \(\sigma\), \(W = \lambda \sigma H_1 H_2 + M_s \sigma^2\) gives upon integrating out \(\sigma\), an effective \(W\):

\[
W = \frac{\lambda^2}{M_s^2} (H_1 H_2)^2. \tag{18}
\]

Another possibility is to have two massive SU(2) doublet superfields \(H_{3,4}\) which couple to the two MSSM Higgs doublets \(H_{1,2}\). Ignoring for a moment any gauge interactions, then from

\[
\mathcal{L} = \int d^4 \theta \left[ \sum_{i=1}^4 H_i H_i + (v_1 H_1 H_3 + v_2 H_2 H_4 + h.c.) \right] \\
+ \int d^2 \theta \left[ \mu H_1 H_2 + M_s H_3 H_4 \right] + h.c. \tag{19}
\]
one finds after integrating out $H_{3,4}$ (with $\mu \ll M_s$):

$$L = \int d^4 \theta \left[ H_1^\dagger H_1 + H_2^\dagger H_2 + \sum_{i=1}^{10} \frac{v_i^2}{M_s^2} H_i^\dagger H_i \right]$$

$$+ \int d^4 \theta \left[ \mu H_1 H_2 + \frac{V_1 V_2}{M_s} H_1 H_2 \right] + \text{h.c.} \quad (20)$$

If gauge interactions are present, the last term becomes

$$\frac{V_1 V_2}{4 M_s} \int d^4 \theta \left[ H_2 e^{-V_1} D^2 \frac{V_1}{M_s} H_1 + \text{h.c.} \right] \quad (21)$$

where $V_1 = g_2 \tilde{V}_\sigma - g_3 \tilde{V}_Y$. This Lagrangian contains operators with more than two derivatives and can be unfolded into one with two space-time derivatives only, as seen above. This ends our discussion on the origin of these operators (for details see the Appendix of [25]).

Let us now examine the physical implications of such operators. We consider an extension of the MSSM (hereafter called MSSM$_5$), by all possible $d = 5$ operators which respect the R-parity symmetry. These are similar to the operators in [18], [20], [21]. We ignore the $d = 6$ ones since they are sub-leading. The new Lagrangian is

$$L = L_0 + L_0^{(S)} \quad (22)$$

where

$$L_0 = \int d^4 \theta \left[ 2 \gamma_1 H_1^\dagger e^{V_1} H_1 + 2 \gamma_2 H_2^\dagger e^{V_2} H_2 \right] + \cdots$$

$$+ \left\{ \int d^4 \theta \left[ Q \lambda U U^c - Q \lambda D D^c H_1 - L \lambda E E^c H_1 \right. \right.$$

$$+ \mu H_1 H_2 \left. \right] + \text{h.c.} \right\} \quad (23)$$

The dots in (23) stand for Higgs-independent terms and

$$L_{0}^{(S)} = \frac{1}{M_s} \int d^4 \theta \left[ H_1^\dagger e^{V_1} Q Y_U U^c + H_2^\dagger e^{V_2} Q Y_D D^c \right.$$

$$+ H_2^\dagger e^{V_2} \lambda \gamma_1 H_1 + \gamma_2 H_1^\dagger \lambda \gamma_2 H_1$$

$$+ \delta(\Sigma) \left. \left[ Q U^c T_O Q D^c + Q U^c T_L L E^c + \lambda \lambda \gamma_1 H_1^\dagger H_2 \right] \right]$$

$$+ \text{h.c.} \quad (24)$$

with $Q, U^c, D^c, L, E^c$ the quark and lepton superfields in a self-explanatory notation. Above we introduced the following spurion dependent function-coefficients:

$$A = A(S, S^T), \quad B = B(S, S^T), \quad \Gamma = \Gamma(S, S^T), \quad \Sigma_{1,2} = \Sigma_{1,2}(S, S^T),$$

$$T_Q = T_Q(S), \quad T_L = T_L(S), \quad \lambda_H = \lambda_H(S), \quad Y_F = Y_F(S, S^T), \quad F = U, D, E \quad (25)$$

where $S = M_s \theta^2$ is the spurion superfield and $M_s$ the supersymmetry breaking scale. Any supersymmetry breaking associated with the presence of the above interactions is included using the spurion field technique.

Not all operators in (24) are independent [8]. To remove this operator “redundancy” we introduce the field re-definitions

$$H_1 \rightarrow H_1 - \frac{1}{M_s} D^2 \left[ \Delta_1 H_1^\dagger e^{V_1} (i \sigma_2) \right]^T + \frac{1}{M_s} Q \rho U U^c$$

$$H_2 \rightarrow H_2 + \frac{1}{M_s} D^2 \left[ \Delta_2 H_1^\dagger e^{V_1} (i \sigma_2) \right]^T + \frac{1}{M_s} Q \rho D D^c$$

$$+ \frac{1}{M_s} L \rho E E^c \quad (26)$$

where

$$\rho_F = \rho_F(S); \quad F : U, D, E, \quad \Delta_i = \Delta_i(S, S^T) \quad i = 1, 2 \quad (27)$$

(27) can be chosen arbitrarily. To avoid the presence of flavour changing neutral currents (FCNC), the following simple ansatz is made

$$T_Q(S) = c_Q(S) \lambda_U(0) \otimes \lambda_D(0)$$

$$T_L(S) = c_L(S) \lambda_U(0) \otimes \lambda_E(0)$$

$$\rho_F(S) = c_F(S) \lambda_F(0)$$

$$Y_F(S, S^T) = y_F(S, S^T) \lambda_F(0), \quad F : U, D, E. \quad (28)$$

and, as usual

$$\lambda_F(S) = \lambda_F(0)(1 + A F S), \quad F : U, D, E. \quad (29)$$

With these and a suitable choice for the coefficients of the spurion entering in $\Delta_2$ one can set $T_Q = T_L = 0$ also $A = B = \Gamma = 0$ and $Y_F \rightarrow y_F(S^T) \lambda_F(0), \quad F = U, D, E$. Then $L_{0}^{(S)}$ becomes

$$L_{0}^{(S)} = \frac{1}{M_s} \int d^4 \theta \left[ H_1^\dagger e^{V_1} Q Y_U S^T \right] U^c$$

$$+ H_2^\dagger e^{V_2} \lambda \gamma_1 H_1$$

$$+ \delta(\Sigma) \left[ Q U^c T_O Q D^c + Q U^c T_L L E^c + \lambda \lambda \gamma_1 H_1^\dagger H_2 \right]$$

$$+ \frac{1}{M_s} \int d^2 \theta \left( \lambda \gamma_1(S) H_1 H_2 \right)^2 + \text{h.c.} \quad (30)$$

with

$$\Sigma_2 = g_2 \tilde{V}_\sigma + g_3 \tilde{V}_Y.$$ 

The new Yukawa couplings $Y_F^\prime(S^T), \quad F : U, D, E$ have now a dependence on $S^T$ only:

$$Y_F^\prime(S^T) = \lambda_F(0)(x_0^T + x_1^T S^T) \quad (31)$$

Following (26), the couplings of $L_0$ (23) and $\Sigma_{1,2}$ have acquired, at the classical level, threshold corrections which depend on the scale $M_s$ [25]. The new form of $L_{0}^{(S)}$ in (30) gives the minimal irreducible set of R-parity conserving dimension-five operators that can be present beyond MSSM.
Let us address some of the physical consequences of the higher dimensional operators in $\mathcal{L}^{(5)}$ of (30). For related studies see [9, 10, 11, 12]. The Higgs scalar potential $V_H$ obtained from (30) is:

$$V_H = \tilde{m}_1^2 |h_1|^2 + \tilde{m}_2^2 |h_2|^2 + (B \mu h_1 h_2 + h.c.)$$

$$+ \frac{g^2}{8} (|h_1|^2 - |h_2|^2)^2 + \frac{1}{2} (\eta_3 (h_1 h_2)^2 + h.c.)$$

$$+ (|h_1|^2 + |h_2|^2) (\eta_1 h_1 h_2 + h.c.)$$

$$+ (|h_1|^2 + |h_2|^2) (\eta_2 h_1 h_2 + h.c.)$$

(32)

where $g^2 = g_2^2 + g_1^2$. Here $\eta_1 \propto g^2 M_s/M_s$, $\eta_2 \propto 2 \mu/M_s$, $\eta_3 \propto M_i/M_s$. $\eta_1$ is due to the derivative term in (24). Although its contribution to $V_H$ can be removed by redefinitions (25), up to a finite renormalisation of the soft masses, one can however keep it, in order to see the effects of such renormalisation.

In the limit of large $\tan \beta = v_2/v_1$ with the pseudoscalar mass parameter $m_A$ fixed at a value $m_A > m_Z$, one finds:

$$m_H^2 = m_a^2 + \frac{4m_s^2 v^2}{m_a - m_Z} (\eta_2 - \eta_1) \cot \beta + O(\cot^2 \beta)$$

(33)

where $v^2 = v_1^2 + v_2^2$. This would suggest that an increase of the mass of the lightest Higgs above $m_Z$ would be possible, thus lifting the tree level bound we have in the MSSM. However, this expansion valid at large $\tan \beta$ only, breaks perturbative expansion in $1/M_s$, since then dimension-six operators and higher are relevant. Note that $\eta_1$ plays no role in the relation among physical masses since

$$m_H^2 + m_A^2 = m_a^2 + m_2^2 + 2 \eta_2 v^2 \sin 2 \beta + \eta_3 v^2$$

(34)

A numerical analysis shows that the lightest Higgs mass can be increased mildly relative to $m_Z$ by up to $\epsilon = 16\%$ ($m_h \leq 105$ GeV) for $m_A$ close to $m_Z$; however if $m_A$ increases above $m_Z$, $\epsilon$ is very small. In conclusion, quantum corrections are still needed for a value of $m_h$ above the bound of 114 GeV, like in the MSSM; however, in the MSSM the amount of stop mixing needed to achieve this can be relaxed relative to the MSSM case. In conclusion the MSSM Higgs sector is rather stable under the addition of higher dimensional operators, in our approximation of including only $d = 5$ operators.

**Physical consequences: new couplings from $\mathcal{L}^{(5)}$.**

Another consequence of the presence of the irreducible set of dimension-five operators in (30) is the generation of MSSM at the tree level. One new coupling is a "wrong"-Higgs Yukawa coupling, which exchanges usual holomorphic dependence on one Higgs by the dependence on the hermitian conjugate of the other ($H_i \leftrightarrow H^*_j$) [13, 14]. Such couplings can also arise in the MSSM at one loop, upon integrating out massive squarks, where they are suppressed by $M^2 / M_c \times (\text{loop} \times \text{factor})$. Here they are suppressed by $M_t / M_s$ only, as seen below:

$$\frac{M_s}{M_s} x_2^U (\lambda_0^U)_{ij} (h^* \tilde{q}_L) u^c_{R f} + h.c.$$

$$\frac{M_s}{M_s} x_2^D (\lambda_0^D)_{ij} (h^* \tilde{q}_L) d^c_{R f} + h.c.$$ 

$$\frac{M_s}{M_s} x_2^E (\lambda_0^E)_{ij} (h^* \tilde{L}_L) e^c_{R f} + h.c.$$ 

(35)

with the notation: $\lambda_0^F \equiv \lambda_F(0)$, $F : U, D, E$, and where $x_0^F$ can be read from (31). These couplings can bring a $\tan \beta$ enhancement of a prediction for a physical observable, such as the bottom quark mass, relative to bottom quark Yukawa coupling:

$$m_b = \frac{v \cos \beta}{\sqrt{2}} \left( \lambda_0^b + \delta \lambda_0^b + \Delta \lambda_0^b \tan \beta \right)$$

(36)

Here $\lambda_0^b$ is the ordinary bottom quark Yukawa coupling, $\delta \lambda_0^b$ is its one-loop correction in the MSSM and $\Delta \lambda_0^b$ is a "wrong"-Higgs coupling contribution, obtained from integrating our massive squarks at one loop in the MSSM, which in our case receives an additional correction from (35). This last correction can be larger than its one-loop-generated MSSM counterpart [13, 15, 16, 17]. This can bring a $\tan \beta$ enhancement of the Higgs decay rate into bottom squarks pairs.

Note that in the MSSMs defined by eq. (30), couplings proportional to $M_i$ involving "wrong"-Higgs A-terms are not present, given our FCNC ansatz (28) leading to (31). If this ansatz is not imposed on the third generation, then one could have such terms:

$$\frac{M^2}{M^*} \left[ y_{u,3} \tilde{q}_{L,3} \tilde{u}_{R,3} + y_{d,3} \tilde{q}_{L,3} \tilde{d}_{R,3} + y_{e,3} \tilde{e}_{L,3} \tilde{e}_{R,3} \right]$$

where $y_{f,3}$, $f = u, d, e$ are the coefficients of the component $SS^f$ of $Y'(S, S')$ of the third generation.

There are also new, important supersymmetric couplings that are generated, which affect the amplitude of processes like quark + quark $\rightarrow$ squark + squark, or involving (s)leptons as well. These are

$$\alpha_0^U (\lambda_0^U)_{ij} (\lambda_0^U)_{kl} \tilde{q}_{L i} \tilde{d}_{R j} q_{L k} \tilde{u}_{R l} + h.c.$$ 

$$\alpha_0^D (\lambda_0^D)_{ij} (\lambda_0^D)_{kl} \tilde{q}_{L i} \tilde{u}_{R j} q_{L k} \tilde{d}_{R l} + h.c.$$ 

$$\alpha_0^E (\lambda_0^E)_{ij} (\lambda_0^E)_{kl} \tilde{L}_i \tilde{e}_{R j} l_{L k} \tilde{e}_{R l} + (L \leftrightarrow Q, E \leftrightarrow U) + h.c.$$ 

(37)
These couplings can be important particularly for the third generation. The largest effect would be for squarks pair production from a pair of quarks; the corresponding amplitude can be comparable to the MSSM tree level contribution [18, 19]. Consider for example $qq \rightarrow \tilde{q}\tilde{q}'$, generated by a tree-level gluon exchange. The MSSM amplitude behaves as the first term in $A_{qq\rightarrow \tilde{q}\tilde{q}'}^{\text{total}}$:

$$A_{qq\rightarrow \tilde{q}\tilde{q}'}^{\text{total}} \sim \frac{g^2}{\sqrt{s}} + \frac{\lambda^U \lambda^D}{M_*}$$  \hspace{1cm} (38)$$

where $\sqrt{s}$ is the center of mass energy. The second term is generated by the additional couplings in (37). While the MSSM contribution decreases with $s$, the second term is constant with potentially significant effects.

### Conclusions

Effective field theories provide an appropriate framework for the study of new physics beyond the SM and the MSSM. In these theories our ignorance of high energy physics is parametrised in terms of higher dimensional operators, which are organised in terms of inverse powers of the scale of new physics $M_*$. To further restrict the exact form of the effective action, other organising criteria are used, such as symmetry principles inspired by the low energy phenomenology. Using these criteria, one is then able to make testable predictions for the low energy observables. This is important since often the exact details of the high-scale, fundamental theory are not known in detail (moduli problem, vacua degeneracy, etc). It is difficult to make testable predictions in this, and the use of effective theories can provide a successful approach.

There are two classes of higher dimensional operators, with up to two space-time derivatives (class A) and with more than two such derivatives (class B). While the former class is more studied, class B is also a common presence. Class B operators are generated in 4D renormalisable theories, by integrating massive fields, with the result of generating an infinite series of derivatives. Truncating this series can generate ghosts fields, which signals that the theory is only valid below the scale of these states. Using general field redefinitions one can reformulate a theory with both classes of operators in terms of a second-order one with class A operators only. This can have applications when coupling such theories to gravity.

We considered the study of the R-parity conserving, dimension-five operators and their generalisation to the supersymmetry breaking case, that extend the MSSM Lagrangian. Not all these operators are independent. Using general, spurious dependent field redefinitions, we removed the redundancy and identified the minimal irreducible set of dimension-five operators that can exist beyond the MSSM. The phenomenological implications of this MSSM extension were studied. It turns out that the MSSM Higgs sector is rather stable, in the approximation used, under the presence of these operators, although a mild increase of the lightest neutral Higgs scalar may be present, up to $\approx 105$ GeV, before taking into account quantum corrections. Additional couplings are also generated, and can dominate their counterparts generated in the MSSM alone at the loop-level. For example squark production and Higgs decays into b-quarks are enhanced by the presence of dimension-five operators. The method to identify the irreducible set of higher dimensional operators is general and can be applied to other models, too.

### ACKNOWLEDGMENTS

This work was partially supported by ANR (CNRS-USAR) contract 05-BLAN-007901, INTAS grant 03-51-6346, EC contracts MRTN-CT-2004-005104, MRTN-CT-2004-503369 and MRTN-CT-2006-035863, CNRS PICS # 2530, 3059, 3747, 4172, and European Union Excellence Grant MEXT-CT-2003-509661.

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