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A short introduction to supersymmetry

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1 Lecture 1 (1h30min): Supersymmetry

1.1 What is here, what is not, what to read

Due to lack of time, only final results will be presented, no real derivations or explanations. I will use here only the component notation, the superspace is anyway only a technical detail. On top of that I will consider only perturbative physics, all the rich phisics of dynamical supersymmetry breaking will have to be skipped. Also, no collider physics or cosmology will be discussed.

There are many very good reviews and books on supersymmetry, for example, among many others: [1], [2] and [3] are the classical references, [4] is fastly becoming classical and it is continuously updated, [5] is a very useful introduction with all computational details, [6] is a clear overview of the main features, [7] and [8] are for those who like more formal approach, [9] and [10] are reviews on susy breaking, [11] is part of the Weinberg's famous course on quantum field theory, [12] is for fans of superspace.

1.2 Introduction

In field theory there are Lorentz (Poincare) symmetries and internal (gauge or global) symmetries. Any other possibility? Sohnius et al. showed that it is possible the so called supersymmetry, a symmetry that connects bosons with fermions. For this to happen the susy generator must have fermionic (Grassman) character. It predicts (super)multiplets in which different elements have different spins. All these elements should have the same mass. In the early days it was hoped that one could use known particles to form these multiplets (neutrino as the fermionic partner of the photon), but it was soon realized (Fayet) that this is impossible, and that one needs to really double all the known particles (a new boson for each known fermion and a new fermion for each known boson). Not only, two Higgses are needed in the supersymmetric version of the standard model, not just one.

Supersymmetric theories thus contain the same number of complex bosons and chiral fermions. This means that for each particle we need a partner with all the internal quantum number the same, but with a different spin:

$$\begin{aligned} FERMIONS\ (\psi_i) \quad &<-> \quad SFERMIONS\ (\phi_i)\ , \\ GAUGE\ \ BOSONS\ (W^a) \quad &<-> \quad GAUGINOS\ (\lambda^a)\ , \\ HIGGS\ (H) \quad &<-> \quad HIGGSINO\ (\tilde{H})\ . \end{aligned}$$

The particles on the left are our standard model particles, the ones on the right are their superpartners.

1.3 Some basics on supersymmetry

Let us shortly summarize some basic points in supersymmetry.

There are chiral multiplets, which members are a spin 1/2 Weyl fermion ψ_i (2 degrees of freedom on-shell) and a complex spin 0 boson ϕ_i (2 d.o.f.o.s.), both with the same quantum numbers.

The kinetic terms for the **chiral multiplets** are the same as in ordinary, non supersymmetric case, what changes is the potential, as well as the mass and Yukawa terms of the fermions. They are all described by a single holomorphic function, called superpotential, $W(\phi)$, such that the (F-term) potential is

$$V_F = \sum_i \left| \frac{\partial W}{\partial \phi_i} \right|^2 \ , \tag{1}$$

while the terms with fermion bilinears are given by

$$-\frac{1}{2} \frac{\partial^2 W}{\partial \phi_i \partial \phi_j} \psi_i \psi_j + h.c. \tag{2}$$

If the superpotential is invariant under some internal symmetries, so is the potential. If one considers renormalizable theories, then the superpotential is a polinomial up to the third power in fields.

There are also **vector** massless **multiplets**, which members are a real vector field A_μ (2 d.o.f.o.s.) and a Weyl fermion λ (2 d.o.f.o.s.). Such multiplets take care of the gauge interactions.

On top of changing all derivatives into covariant ones as usually, the introduction of gauge interactions adds two terms in the Lagrangian: the (D-term) potential

$$V_D = \frac{g^2}{2} \sum_a \left(\sum_{i,j} \phi_i^* T_{ij}^a \phi_j \right)^2 , \tag{3}$$

where T^a is the a^{th} generator of the gauge group in the representation of the ϕ chiral multiplets.

The other additions are new Yukawa interactions:

$$-g\sqrt{2}\phi_i^* T_{ij}^a \psi_j \lambda^a + h.c. , \tag{4}$$

1.4 The standard model supersymmetrized

Now we are well equipped to try to write down the superymmetric version of the standard model Lagrangian. From above it is clear, that the only unknown part is the superpotential W. Instead of the notation ψ_i let's use the usual Q for the quark weak doublet, u^c for the (charge conjugated) up quark weak singlet ($(u^c)_L \equiv C\bar{u}_R^T$), d^c for the down singlet, L for the leptonic weak doublet and e^c for the leptonic singlet, while the same notation with a tilde will denote their scalar (complex) partners.

The known Yukawas from the standard model make some terms a must:

$$W_Y = Q_i H_u \left(Y_u\right)_{ij} u_j^c + Q_i H_d \left(Y_d\right)_{ij} d_j^c + L_i H_d \left(Y_e\right)_{ij} e_j^c \tag{5}$$

where i, j run over generations. Notice that the standard model Higgs would be H_u (or $\tau_2 H_d^*$). This makes it clear, why we need two Higgses, since the superpotential is a holomorphic function of the fields.

The gauge quantum numbers of the fields are as in the standard model, with the hypercharge of the two Higgses determined by the invariance of the superpotential. This is the second reason for the necessity to double the Higgs sector: the chiral anomaly cancellations.

The higgsinos would be massless here, so one adds the term

$$W_\mu = \mu H_u H_d . \tag{6}$$

In the SM gauge interactions automatically conserve baryon and lepton numbers at the renormalizable level. Here this is not true anymore, and one could write down in principle

$$\Delta W = \epsilon_i L_i H_u + \lambda_{ijk} L_i L_j e^c_k + \lambda'_{ijk} L_i Q_j d^c_k + \lambda''_{ijk} u^c_i d^c_j d^c_k \tag{7}$$

It is easy to see that such terms would be catastrophic for the stability of the proton. Since proton longevity is extremely well measured, the above terms are bound to be very small. It is thus quite usual to assume that they are zero. In this limit, there is a new Z_2 discrete symmetry of the potential, under which the known SM particles and both Higgses are even, while the new partners are odd. We will assume it from now on, although from the SM point of view there is absolutely no reason for it. In some SO(10) grand unified theories this low-energy extra symmetry can be derived from reasonable assumptions, and are thus motivated.

We will call the Lagrangian described by the usual (gauge invariant) kinetic terms plus terms (1)-(6) the supersymmetric limit of the standard model. As we will see next time, this model is unrealistic. Before noticing it, let us first mention its virtues.

1.5 Hierarchy problem

The first one is the stabilization of the gauge hierarchy problem. This problem is connected with the fact, that usually in field theory scalar masses are not protected by quantum effects. Even if the tree order scalar mass is small, it can get arbitrary high corrections by loops with heavy fields running inside. This is not true for spin 1 and spin 1/2 fields, since their corrections are protected by gauge symmetry and chiral symmetry. Supersymmetry solves this problem by linking a scalar (spin 0) mass to its superpartner fermion mass.

Although these thoughts sound rather philosophical, there is actually also a prediction coming from them. In fact, once you break supersymmetry and give masses to the sfermions (we will see in the next lecture that this is unavoidable), supersymmetry still preserves some correlation among different masses. It can be shown, that the stop and Higgs mass are related, a fact that we will hopefully be able to check at LHC or at some next collider. In the limit of a large stop mass this reduces to the approximate relation between the lightest neutral Higgs mass m_{h^0} and the stop mass $m_{\tilde{t}}$ (see for example [13] and references therein)

$$m_{h^0}^2 \approx m_Z^2 \cos^2 2\beta + 6 \left(\frac{\alpha_2}{4\pi}\right) \frac{m_t^4}{m_W^2} \log \left(\frac{m_{\tilde{t}}^2}{m_t^2}\right) , \tag{8}$$

where $\tan\beta \equiv \langle H_u \rangle \,/\, \langle H_d \rangle$.

1.6 Radiative electroweak symmetry breaking

In the SM the potential is

$$V = m_H^2 \left|H\right|^2 + \lambda \left|H\right|^4 \ . \tag{9}$$

To have the electroweak symmetry breaking, one needs $m_H^2 < 0$. There is no explanation why this is so. Also, it is known that this mass cannot change sign with running in the SM.

In supersymmetric extensions of the SM there are also other bosons with masses $m_{\tilde{t}}^2$, $m_{\tilde{t}^c}^2$, etc. Clearly they must be positive (only H_u and H_d need negative square masses), otherwise we would for example break too much gauge symmetry, as for example the colour. So here one would like to understand why is the Higgs special with respect to the other bosons, for example the stop. In other words, one would like to see why the Higgs mass square is negative, while the masses square for the other bosons are positive. It is a remarkable fact that supersymmetry combined with gauge symmetry can give a possible explanation of this behaviour.

The differences among different fields is coming from the different renormalization group equations. Consider the system of the Higgs and the two stop masses (similar equations can be written for H_d, $\tilde{b}$ and $\tilde{b}^c$, changing $y_t \to y_b$, which is relevant if $\tan\beta$ is large) :

$$\mu \frac{dm_{H_u}^2}{d\mu} = 3 \frac{y_t^2}{16\pi^2} \left(m_{H_u}^2 + m_{\tilde{t}^c}^2 + m_{\tilde{t}}^2\right) + \ldots , \tag{10}$$

$$\mu \frac{dm_{\tilde{t}^c}^2}{d\mu} = 2 \frac{y_t^2}{16\pi^2} \left(m_{H_u}^2 + m_{\tilde{t}^c}^2 + m_{\tilde{t}}^2\right) + \ldots , \tag{11}$$

$$\mu \frac{dm_{\tilde{t}}^2}{d\mu} = 1 \frac{y_t^2}{16\pi^2} \left(m_{H_u}^2 + m_{\tilde{t}^c}^2 + m_{\tilde{t}}^2\right) + \ldots , \tag{12}$$

where the dots denote other contributions not necessary for the understanding of the phenomenon. We see that the righthandsides are the same except for the numerical factor in front of them. These numerical factors are just consequences of gauge symmetry, i.e. of the representation for the Higgs and stops under the gauge groups, which is not difficult to read from the relevant 1-loop Feynman diagrams (different number of fields running in the loop).

So we see that all the three masses increase with the scale (the righthand-sides are positive) and that the Higgs mass increase faster than the other two. So one can think that the three masses are for example equal at a large scale, for example at the Planck scale, $\mu = M_{Planck}$. Then one evolves the three masses down using the RGE (10)-(12). Since $m^2_{H_u}$ decreases faster than the other two masses squared with decreasing scale μ, one could finish with a negative $m^2_{H_u} < 0$ at small scales, but positive $m^2_{\tilde{t}^c,\tilde{t}} > 0$. This is indeed what happens when one chooses realistic values for the parameters, giving a very appealing explanation for the electroweak symmetry breaking in the MSSM.

One can solve the system (10)-(12) with the result

$$\begin{aligned} m^2_{H_u}(m_W) &= m_0^2 \left[\frac{3}{2} \left(\frac{m_W}{\Lambda} \right)^{3/(8\pi^2)} - \frac{1}{2} \right] , \quad (13) \\ m^2_{\tilde{t}^c}(m_W) &= m_0^2 \left[\left(\frac{m_W}{\Lambda} \right)^{3/(8\pi^2)} \right] , \quad (14) \\ m^2_{\tilde{t}}(m_W) &= m_0^2 \left[\frac{1}{2} \left(\frac{m_W}{\Lambda} \right)^{3/(8\pi^2)} + \frac{1}{2} \right] , \quad (15) \end{aligned}$$

where we assumed a universal boundary condition at the large scale Λ:

$$m^2_{H_u}(\Lambda) = m^2_{\tilde{t}^c}(\Lambda) = m^2_{\tilde{t}}(\Lambda) = m_0^2 \ . \quad (16)$$

A reasonable value for Λ is for example the Planck scale M_{Pl}, which gives the mass term for $m^2_{H_u}$ negative, while the other two turn out to be safely positive.

1.7 Unification of couplings

The 1-loop RGE for the gauge couplings g_i are

$$\mu \frac{dg_i}{d\mu} = \frac{1}{16\pi^2} b_i g_i^3 \ , \quad (17)$$

where $i = 1, 2, 3$ denotes the gauge group and b is the so called β function coefficient. It can be straightforwardly calculated in any theory via (G,F,B stay for gauge bosons, fermions, bosons)

$$b = -\frac{11}{3}C_2(G) + \frac{2}{3}T_F + \frac{1}{3}T_B \ . \tag{18}$$

The Dynkin index

$$T(R)\delta_{ab} = Tr\left(T_a(R)T_b(R)\right) \tag{19}$$

and the second Casimir

$$C_2(R)\delta^{ij} = \sum_a \left(T_a(R)T_b(R)\right)^{ij} \tag{20}$$

depend on the choice of the gauge group and on the representation involved. The indices a, b run over the generators of the group ($N^2 - 1$ in SU(N)), while i, j run from 1 to the dimension of the representation. The normalization usually chosen is $T = 1/2$ for the fundamental representation (quarks, leptons). Then one has in the SU(N) group for the fundamental representation $C_2 = (N^2 - 1)/(2N)$, and for the adjoint $T = C_2 = N$. The dimension of the representation is N for fundamentals and $N^2 - 1$ for adjoint. To remember also that in SU(2) the generators in the fundamental are the Pauli matrices $T_a^{ij} = \tau_a^{ij}/2$, while in the adjoint representations are the Levi-Civita antisymmetric tensor $T_a^{ij} = -i\epsilon_{aij}$.

For supersymmetric theories we know that for each fermion (boson) there is a boson (fermion) in the same group representation, so (18) can be written more compactly as

$$b = -3C_2(G) + T \ . \tag{21}$$

A simple exercise shows that in MSSM $b_i = (33/5, 1, -3)$, while in SM $b_i = (41/10, -19/6, -7)$. Negative coefficients here mean asymptotic freedom. One knows the experimental values of g_i at M_Z and can evolve them towards larger scales μ using (17). It is then surprising that in the supersymmetric case the three couplings unify in a single point at $\mu \approx 10^{16}$ GeV. To appreciate this fact one should notice that this unification fails badly in the nonsupersymmetric case. So, if we have supersymmetric partners at M_Z or close to 1 TeV as required by naturalness (hierarchy problem), then we have unification of gauge couplings for free!

1.8 Candidate for dark matter

Once R-parity is preserved, one has at his disposal a stable massive particle (the lightest supersymmetric partner - LSP), i.e. the lightest particle with negative R-parity. If this particle is electrically and colour neutral, but feels weak interactions (the so-called neutralino), like for example the Higgsino, bino or neutral wino, then its mass of the order of 100 GeV or 1 TeV is automatically of the right order of magnitude to be a possible dark matter candidate. This is in contrast with the standard model, where no such candidate exists.

2 Lecture 2 (1h30min): Susy breaking

2.1 Soft terms

So far we described only the good qualities of the supersymmetric version of the standard model. Obviously there are many drawbacks, the main is that the theory described last time is completely unrealistic. In fact, it predicts the doubling of all the particles we know, i.e. for any known fermion there is a new scalar with the same quantum numbers. This is clearly impossible, there is for example no scalar particle with the quantum numbers and the mass of the electron. In other words, supersymmetry must be broken. The most pragmatic way to obtain it is to add by hand a new set of terms, which break supersymmetry softly. In this way one maintain some of the good features encountered before, but still make the theory possible. In general this amounts to adding the following dimension two and three terms to the Lagrangian:

$$\mathcal{L}_{soft} = m_i^2 \phi_i^* \phi_i + \Big(B_{ij} \phi_i \phi_j + A_{ijk} \phi_i \phi_j \phi_k + m_A \lambda^A \lambda^A + h.c. \Big) \ , \qquad (22)$$

where i, j, k run over all different representations and A over all different gauge groups. The new parameters are undetermined, but can be sometimes restricted by experimental data. For example, from the fact that we have never found the sfermions, we conclude that their masses must be larger than around 100 GeV. Other stringent restrictions come from the rare flavour changing neutral currents: in the standard model these processes are suppressed by the one-loop factor ($\approx 1/(4\pi)^2$), the GIM mechanism ($\approx m_i/m_W$)

and by the small off-diagonal elements of the CKM mixing matrix. The new terms in (22) bring new contributions to the FCNC (in $K-\bar{K}$, $B-\bar{B}$ mixing, $b \to s\gamma$, etc). These are automatically compatible with the experimentally measured or bounded values only in some cases, for example among others:

- if the soft masses in (22) are much higher than naively expected form naturalness, i.e. on the order of 10^{3-4} TeV or more;

- if the coefficients A in (22) are very small, generation independent or proportional to the relevant Yukawa matrices in the superpotential.

As you can see, the bad side of such a model (which we call the minimal supersymmetric standard model - MSSM) is a big proliferation of unknown parameters, something like 100 or so. The desire to simplify things adding a new symmetry (supersymmetry) clashed with the facts of nature, forcing us to add the new terms (22). Theoretically it would be thus much better to have a model of calculating or deriving the unknown coefficients. This will be studied in detail in the remaining part of this lecture. We will present two different scenarios.

2.2 Spontaneous susy breaking is not easy

The best option would be to break supersymmetry spontaneously, as we do for example in gauge theories. We will see that this is not easy to do in two examples.

In supersymmetric theories the hamiltonian is proportional to the sum of the squares of the generators of supersymmetry algebra. So the energy is zero iff susy is preserved, and positive if susy is spontaneously broken. We saw in the previous lecture that the potential (energy) is made from two terms, the F-term (1) and the D-term (3). We will consider here only the so called F-term breaking:

$$\frac{\partial W}{\partial \phi_i} \neq 0 \quad \text{for at least some } i \ . \tag{23}$$

Notice that such a requirement automatically implies a massless fermion. In fact, from the minimization of (1), the second derivative of the superpotential in in the field space direction defined by (23) must vanish. This is expected from the Goldstone theorem. In the case of spontaneously broken internal symmetry a massless Goldstone boson appears. Here the spontaneously broken supersymmetry generator is a Grassman object, thus a fermionic zero mode follows. This object is called the goldstino.

Although W is a gauge singlet, in the SM ϕ_i are not. Thus the only field that could break supersymmetry without breaking some unwanted extra gauge symmetry is one of the two Higgses H_u or H_d. Imagine we do it with H_d. Then suppose that we are able to properly change the superpotential so that

$$F_{H_d} \equiv \left\langle \frac{\partial W}{\partial H_d} \right\rangle \neq 0 \ . \tag{24}$$

To see why this cannot work, let us concentrate on the mass of the selectron:

$$W = y_e e h_d^0 e^c + ... \tag{25}$$

From (1) and (24) the mass terms for the selectron in the potential are

$$V = \left(\tilde{e}^* \quad \tilde{e}^c \right) \begin{pmatrix} y_e^2 v_d^2 & y_e F_{H_d}^* \\ y_e F_{H_d} & y_e^2 v_d^2 \end{pmatrix} \begin{pmatrix} \tilde{e} \\ \tilde{e}^{c*} \end{pmatrix} \ , \tag{26}$$

where v_d is the vev of H_d. The eigenvalues satisfy the mass sum rule

$$m_{\tilde{e}}^2 + m_{\tilde{e}^c}^2 - 2m_e^2 = 0 \ . \tag{27}$$

This is in contradiction with what we know from low energy physics: there is no scalar with quantum numbers of the electron and a smaller (or equal) mass.

It is possible to show that any spontaneously broken global supersymmetric model has similar unacceptable mass sum rules at tree order. There are two possible ways out: either one goes local, i.e. to supergravity, or one transmits the information of susy breaking not at tree order, but at one loop. In both cases the mass sum rules change and unwanted constraints get relaxed.

Either way, the mechanism should roughly look as follows: a sector with no interaction with the SM fields (and thus called the hidden sector) is responsible for the spontaneous breaking of supersymmetry. The information that susy is broken in this hidden sector gets thus transmitted to our SM sector either by $1/M_{Pl}$ suppressed higher dimensional terms (as in supergravity) or through an intermediate (messenger) field, that couples to both SM and hidden sector fields. In this scenario loops with external SM fields and internal messenger fields (with susy breaking couplings and/or masses) transmit the information on susy breaking to our sector.

Let us see now in more details the above mechanisms of susy breaking mediation.

2.3 Gravity mediation

In the case of local supersymmetry (supergravity), one needs to introduce the **gravity multiplet**, which essentially means the spin 2 graviton (2 d.o.f.o.s.) plus the spin 3/2 gravitino (also 2 d.o.f.o.s.). Analogous to the case of spontaneously breaking of a local symmetry, where the would-be goldstone boson gets eaten by the longitudinal component of the vector boson, here the gravitino eats the goldstino, acquiring a nonzero mass, $m_{3/2}$. This mass turns out to be the typical scale, and soft masses will be proportional to it.

The potential in the supergravity case becomes (φ_i denote both SM fields ϕ_i and hidden sector fields X)

$$V_F = e^{K/M_*^2}\left[\left(\frac{\partial W}{\partial \varphi^i} + \frac{\partial K}{\partial \varphi^i}\frac{W}{M_*^2}\right)\left(K^{-1}\right)^i_j\left(\frac{\partial W^*}{\partial \varphi^*_j} + \frac{\partial K}{\partial \varphi^*_j}\frac{W^*}{M_*^2}\right) - 3\frac{|W|^2}{M_*^2}\right] , \tag{28}$$

where $M_* = M_{Pl}/\sqrt{8\pi} \approx 2 \times 10^{18}$ GeV is the so called reduced Planck mass and $(K^{-1})^i_j$ is the inverse matrix of $\partial^2 K/\partial\varphi^i\partial\varphi^*_j$. The Kähler function $K(\varphi, \varphi^*)$ is a real function of the fields, which we call canonical, iff $K_{can} = \sum_i \varphi^*_i \varphi_i$, but in a general nonrenormalizable model can be actually modified with higher dimensional polinomials.

Assuming that it is a singlet field X that breaks susy, the order parameter in supergravity is defined as

$$F_X = \frac{\partial W}{\partial X} + \frac{\partial K}{\partial X}\frac{W}{M_*^2} . \tag{29}$$

It must be nonzero, so to break supersymmetry. In flat spacetime one can parametrize this breaking by the gravitino mass:

$$F_X = \sqrt{3} m_{3/2} M_* . \tag{30}$$

As we said before, this field X should have small enough couplings with the SM fields ϕ_i. This is most easily obtained, assuming

$$W(X,\phi) = W(X) + W_{SM}(\phi)\ . \tag{31}$$

Then all the couplings between the two sectors are through $1/M_{Pl}$ suppressed operators.

One then needs the following requirements at $X = \langle X \rangle$:

$$W = m_{3/2} M_*^2 \qquad \text{(zero cosmological constant)}\ , \tag{32}$$

$$\frac{\partial V}{\partial X} = 0 \qquad \text{(extremum of the potential)}\ , \tag{33}$$

$$\left|\frac{\partial^2 V}{\partial X^2}\right| \le \frac{\partial^2 V}{\partial X \partial X^*} \qquad \text{(no tachyons)}\ . \tag{34}$$

As an example take the so-called Polonyi superpotential:

$$W(X) = aX + b \tag{35}$$

and a canonical Kähler. With properly chosen constants a and b, so to satisfy (29), (32)-(34), it is possible to determine all the soft parameters in (22) in terms of few parameters at the Planck scale, which are at the moment still compatible with any experimental constraint.

We were thus able to reduce the 100 and so parameters to very few of them. It has to be kept in mind however, that this is no more than just a reparamentrization of the original soft terms. In fact, there is absolutely no real reason to believe that the superpotential should look like (35) and even less that the Kaehler potential is canonical. In fact, taking for example

$$K = X^* X + \phi_i^* \phi_i + c_{ij} \frac{X^* X \phi_i^* \phi_j}{M_{Pl}^2}, \tag{36}$$

gives the sfermion soft mass squares

$$\left(m_{\tilde f}^2\right)_{ij} = (\delta_{ij} + c_{ij}) \frac{|F_X|^2}{M_{Pl}^2}\ , \tag{37}$$

which not only introduces new couplings (c's), but may very easily be in contradiction with experiment due to possible large contribution to the flavour changing neutral currents. Although there are some ways of making these c's small in the infrared via running, this requires extra physics, and is certainly not a minimalist's approach. It is safe to conclude that supergravity cannot explain the structure, less the particular values of the soft parameters.

2.4 Gauge mediation

The problem with the previous example was, that gravity is not actually flavour blind (the masses of fermions of different genrations are very different!), which can be parametrized by a nontrivial matrix c above. This led people to use as mediators of susy breaking gauge interactions, which are known to be flavour blind

As before, we have to forbid that supersymmetry breaking is transmitted to the SM particles at tree order, so no SM field can satisfy (23). Thus imagine that we have a gauge singlet field X with a non zero F-term

$$F_X = \left\langle \frac{\partial W}{\partial X} \right\rangle \neq 0 \ . \tag{38}$$

Can it be coupled to any of the standard model fields? Allowing only renormalizable couplings the only possibility is to transform the dimension 2 term (6) to a dimension 3 term [14]

$$W_\mu = \lambda_X X H_u H_d \ . \tag{39}$$

The higgsino mass appears from the vev of X,

$$\mu = \lambda_X \langle X \rangle \ . \tag{40}$$

The sum rule (27) applied to the Higgs supermultiplets is not dangerous in this case, since no mass of this multiplet has been measured yet. On the other side, the sfermions do not couple directly to X, so they are not influenced by it at tree order and no mass sum rule thus applies. Unfortunately the one loop correction to the masses of the sfermions gets negative and proportional to the relevant Yukawas [15]. This would destabilize the stop, breaking SU(3).

In other words, one cannot couple only the Higgs to the X field, but must use another pair of multiplets, call them Φ and $\bar{\Phi}$:

$$W = W(X) + \lambda_X X \bar{\Phi} \Phi + W_{SM} \ . \tag{41}$$

These new multiplets are not gauge singlets, so they can interchange gauge bosons (and gauginos) with the SM fields. Thus at leading order in a small F_X/M_Φ^2 ratio, where $M_\Phi = \lambda_X \langle X \rangle$ is the susy preserving Φ mass, the gaugino mass gets a 1-loop contribution

$$m_\lambda = c_i \frac{\alpha_i}{4\pi} \frac{\lambda_X F_X}{M_\Phi} \ , \tag{42}$$

while the sfermions get a two-loop contribution

$$m_{\tilde{f}}^2 = d_i \left(\frac{\alpha_i}{4\pi}\right)^2 \frac{|\lambda_X F_X|^2}{M_\Phi^2} \,, \tag{43}$$

with c_i, d_i $(i = 1, 2, 3)$ depending on the representation under the SM gauge group SU(3)×SU(2)×U(1). Similar relations are possible also for the other soft parameters. Although this seems to solve the flavour problem mentioned above, one must be careful. In fact, there is no gaurantee that the Φ and $\bar{\Phi}$ do not couple or mix with the SM fields. Typically these mediators have the quantum numbers of the sfermions. This mixing can in principle spoil the flavour blindness of the mediation [15].

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