# XIII. THE MSSM

In this section we will discuss an important application of supersymmetry to the physics of electroweak and strong interactions. The possibility of obtaining, through supersymmetry, better UV behaviours for quantum field theory and thus cancellation of divergences in loop diagrams, is the main ingredient which makes supersymmetry an appealing framework to go beyond the Standard Model (SM) of electroweak interactions.

In the following we will present a minimal implementation of the susy framework in the context of electroweak theories, the so called Minimal Supersymmetric Standard Model, also known as MSSM.

## · The SM and its problems

Before presenting the MSSM, it is useful to briefly recall the structure of the SM. This will also give us the possibility to better appreciate the theoretically unsatisfactory aspects of the SM which are the motivation for considering supersymmetry as a possible phenomenologically interesting option.

The SM is an extremely successful description of the properties and behaviour of elementary particles up to the energy scale of the previous generation of collider experiments (eg. LEP). It contains two sectors whose particles have been observed in high-energy collider experiments:

- Matter particles: namely quarks and leptons. They come in three families with identical properties and differing only for their masses. All matter particles are fermions with spin 1/2.

- Interaction particles: namely the three non-gravitational interactions (strong, weak and electromagnetic) which are described by a gauge theory based on an internal symmetry

$$G_{SM} = \underbrace{SU(3)_c}_{\text{Strong}} \otimes \underbrace{SU(2)_L \otimes U(1)_Y}_{\text{Electroweak}}$$

$SU(3)_c$ is the QCD part of the SM, describing strong interactions, while $SU(2)_L \otimes U(1)_Y$ is the electroweak part. The label $L$ in the $SU(2)_L$ subgroup specifies that only the L-handed fermions are charged under this subgroup, while the R-handed ones are uncharged.
The gauge particles (which have spin 1) are the gluons, for the strong interactions, and the $W^\pm$, the $Z^0$ and the photon $\gamma$ for the electroweak interactions.

The fermions can be classified according to their quantum numbers under the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ symmetry. For each family we have:

| | notation | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ |
|---|---|---|---|---|
| quarks | $Q_i = (u_L, d_L)_i$ | 3 | 2 | 1/6 |
| | $\bar{u}_i = u_{Ri}^+$ | $\bar{3}$ | 1 | -2/3 |
| | $\bar{d}_i = d_{Ri}^+$ | $\bar{3}$ | 1 | 1/3 |
| leptons | $L_i = (\nu, l_L)_i$ | 1 | 2 | -1/2 |
| | $\bar{e}_i = l_{Ri}^+$ | 1 | 1 | 1 |

As we said the SM contains three families:

- leptons:
  - electron $e, \nu_e$
  - muon $\mu, \nu_\mu$
  - tau $\tau, \nu_\tau$

- quarks:
  - up $u$, down $d$
  - charm $c$, strange $s$
  - top $t$, bottom $b$

Given that all the fermions have at least a non trivial charge we can not include Majorana mass terms in the SM. (We neglect here neutrino masses, as well as the possibility of having a R-handed neutrino.) Moreover, it is easy to see that, given that L-handed and R-handed fermions are in different representations, we can not write down Dirac mass terms as well.

To be able to give masses to the matter fields we need to break the electroweak symmetry.

- There is also another reason for which the $SU(2)_L \otimes U(1)_Y$ symmetry must be broken: an exact gauge symmetry implies that the corresponding gauge bosons are massless, so they would mediate long-range forces. But in nature we do not observe $SU(2)_L \otimes U(1)_Y$ forces at long range, we only observe the $U(1)_{em}$ electromagnetic force mediated by the photon. Hence we must have a (spontaneous) breaking

$$\underline{SU(2)_L \otimes U(1)_Y \rightarrow U(1)_{em}}.$$

(The electromagnetic group is a subgroup of $SU(2)_L \otimes U(1)_Y$, and the electromagnetic charge is given by $T_{3L} + Y$, where $T_{3L}$ is the charge under the $SU(2)_L$ subgroup generated by $\sigma_3$.)

By a spontaneous breaking the $W^\pm$ and the $Z^0$ bosons acquire a mass, so they mediate only short-range forces, while the photon remains massless.

The breaking of the electroweak symmetry in the SM is induced by a scalar field, the Higgs boson H. It has the following quantum numbers

| | | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ |
|---|---|---|---|---|
| Higgs | $H$ | $1$ | $2$ | $1/2$ |

that is it is a complex $SU(2)_L$ doublet. The most general renormalizable Lagrangian for the Higgs field includes a potential

$$V = m_H^2 |H|^2 + \lambda |H|^4 .$$

Electroweak symmetry breaking is induced by a non-vanishing VEV for $H$, this is realized if $m_H^2 < 0$ and $\lambda > 0$, in which case

$$\langle H \rangle = \sqrt{-\frac{m_H^2}{2\lambda}} .$$

Through the gauge interactions, the Higgs VEV generates a mass term for the gauge fields

$$(D_\mu H)^\dagger (D^\mu H) \longrightarrow \frac{g^2}{4} \langle H \rangle^2 W_\mu^+ W^{-\mu} + \frac{g^2}{8\cos^2\theta_W} \langle H \rangle^2 Z_\mu Z^\mu .$$

From the experimental values of $m_W$ and $m_Z$ we get

$$\underline{\langle H \rangle \simeq 246\ \mathrm{GeV}.}$$

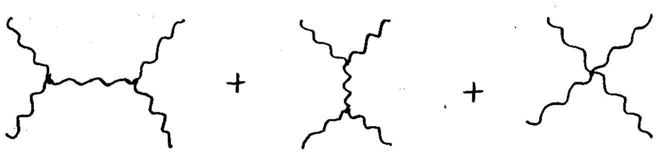
The Higgs induces also mass terms for the fermions through the Yukawa interactions

$$y_{ij}^u\, Q_i H \bar{u}_j + y_{ij}^d\, Q_i H^c \bar{d}_j + h.c.$$

· <u>The W-W scattering</u>

An important consequence of the presence of an Higgs doublet in the SM is the unitarization of the W-W scattering. If we consider a model with spontaneous breaking of the electroweak symmetry, but without the physical Higgs mode, we get that the scattering of the longitudinal components of the W's has an amplitude which grows with the energy.

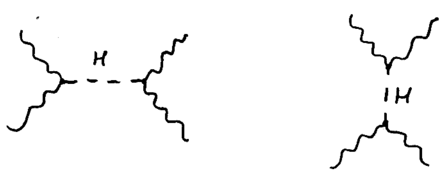
The sum of the diagrams



gives a result

$$\sigma(W_L W_L \to W_L W_L) \simeq \frac{g^2}{4 m_W^2}(s+t) \propto E^2$$

and we get a violation of unitarity at energies $E \sim 1\ \mathrm{TeV}$. This behaviour is modified by the presence of the physical Higgs mode, through the diagrams of the type



Adding the contribution of these diagrams we find that, for energies above the Higgs mass,

$$\sigma(W_L W_L \to W_L W_L) \sim \text{const.}$$

Hence, if the SM includes an Higgs with a relatively low mass ($m_H \lesssim 1\,\text{TeV}$) the W–W scattering is unitarized and the theory remains perturbative at energies much larger than the scale of electroweak breaking ($E \sim 1\,\text{TeV}$).

On the other hand, if the physical Higgs is absent or is too heavy, the theory becomes non perturbative at energies $E \gtrsim 1\,\text{TeV}$.

NOTE. The physical Higgs mode is not necessary to write an effective theory with electroweak symmetry breaking. What we need are the Goldstone modes which correspond to the breaking of a global $SU(2)_L \otimes U(1)_Y$ to the $U(1)_{EM}$ subgroup. These modes are necessary because they play the role of the longitudinal components of the $W^\pm$ and the $Z^0$ bosons, which are present when the gauge bosons acquire a mass. (Recall that a massive gauge boson has 3 degrees of freedom, while a massless gauge boson has only 2 degrees of freedom.)

· The Hierarchy problem

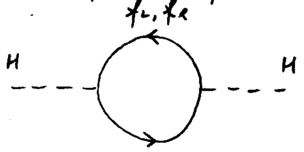
The Higgs mass term is the only mass term allowed in the SM Lagrangian (before electroweak symmetry breaking) and is also the only superrenormalizable term.

As we already discussed in the introduction of the course the Higgs mass is sensitive to the UV physics through loop effects. This sensitivity comes from the fact that there is no symmetry which protects the Higgs mass.

The situation for the gauge boson masses and for the fermion masses is radically different, because in these cases there are some symmetries which protect them from large radiative corrections.

- Gauge invariance does not allow to write a mass term for the gauge bosons. The masses arise only when the symmetry is spontaneously broken, so they are determined by the scale of symmetry breaking.
- The chiral structure of the theory does not allow mass terms for the fermions. Moreover if the Yukawa couplings are absent, the fermionic sector has a larger invariance under chiral transformations (transformations which rotate independently the L-handed and the R-handed fermions). Chiral symmetry protects the Yukawa couplings from large radiative corrections because loop corrections to the couplings are proportional to the tree level couplings, so that they are small if the tree-level couplings are small.

Let's now consider the Higgs field. We can compute the radiative corrections to the Higgs mass induced by loops of fermions (eg. the top) by using a cut-off regularization. We find, for a fermion with a Yukawa $\mathcal{L}_{Yuk} = -\frac{y}{\sqrt{2}} H \bar{f}_L f_R + h.c.$,




$$\Rightarrow \quad \delta m_H^2\big|_{f_L,f_R} = -\frac{|y|^2}{8\pi^2}\left[\Lambda^2 - 3 m_f^2 \ln\left(\frac{\Lambda^2+m_f^2}{m_f^2}\right) + \dots\right].$$

This shows that there are quadratic and logarithmically divergent corrections to the Higgs mass. The physical Higgs mass is given by

$$m_H^2\big|_{phys} = m_H^2\big|_{bare} + \delta m_H^2\big|_{f_L,f_R}.$$

This implies that, if we assume that the SM has an high cut-off, in order to get a light Higgs we need to fine-tune the bare Higgs mass against the radiative corrections (and we need to tune the bare mass at each loop level). For example if we assume that the SM is valid up to the Planck mass ($M_{Pl} \sim 10^{19}\,GeV$), in order to have $m_H^2\big|_{phys} \lesssim 1\,TeV$, we need a tuning on $m_H^2\big|_{bare}$ of the order $10^{-30}$. This kind of tuning seems very unnatural, although it is not an inconsistency of the theory.

To obtain a theory without this unwanted tuning, we expect that some new physics should appear at an energy scale $E \lesssim 1\,TeV$, which could stabilize the Higgs mass.

Supersymmetry provides a symmetry which protects the Higgs potential from large corrections. To understand the mechanism, we consider the contribution to the Higgs mass coming from loops of scalar particles. We take a pair of scalars $\phi_L$ and $\phi_R$ with the interactions

$$\mathcal{L}_{scalar} = -\frac{\lambda}{2} H^2\left(|\phi_L|^2 + |\phi_R|^2\right) - H\left(\mu_L |\phi_L|^2 + \mu_R |\phi_R|^2\right) - m_L^2 |\phi_L|^2 - m_R^2 |\phi_R|^2.$$

We get


$$\Rightarrow \quad \delta m_H^2\big|_1 = \frac{\lambda}{16\pi^2}\left[2\Lambda^2 - m_L^2 \ln\left(\frac{\Lambda^2+m_L^2}{m_L^2}\right) - m_R^2 \ln\left(\frac{\Lambda^2+m_R^2}{m_R^2}\right) + \dots\right]$$

and


$$\Rightarrow \quad \delta m_H^2\big|_2 = -\frac{1}{16\pi^2}\left[\mu_L^2 \ln\left(\frac{\Lambda^2+m_L^2}{m_L^2}\right) + \mu_R^2 \ln\left(\frac{\Lambda^2+m_R^2}{m_R^2}\right) + \dots\right].$$

We notice that

- if $\lambda = |y|^2$: the quadratic divergences are cancelled.
- if also $m_f = m_L = m_R$ and $\mu_L^2 = \mu_R^2 = 2\lambda m_f^2$: also the logarithmic divergences are cancelled.

We have already seen that supersymmetry ensures the above relations, so if susy is unbroken all the divergences are cancelled.

Of course, supersymmetry must be broken for phenomenological reasons. However, if we want to solve the fine-tuning problem we need to cancel the quadratic sensitivity to the cut-off $\Lambda$ (the logarithmic divergences only introduce a very mild dependence on $\Lambda$). To ensure this cancellation we need to break susy in a "controlled" way, so that the relation $\lambda = |y|^2$ is preserved. As we will see, this is obtained by assuming that susy is broken at low energy only by soft-breaking terms (that is by superrenormalizable terms, such as supersymmetry-breaking mass terms, which introduce a mass split between the SM particles and their superpartners).

· The particle content of the MSSM.

The field content of the MSSM consists of the SM fields and the corresponding superpartners. Only the Higgs sector requires to be enlarged as we will discuss in the following.

Before discussing the Higgs sector we summarize the MSSM field content (and our notation) in a table

| | | bosons | fermions | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ |
|---|---|---|---|---|---|---|
| Quarks | $Q_i$ | $(\tilde{u}_L, \tilde{d}_L)_i$ | $(u_L, d_L)_i$ | 3 | 2 | 1/6 |
| | $\bar{u}_i$ | $\tilde{u}^*_{Ri}$ | $\bar{u}_i = u^+_{Ri}$ | $\bar{3}$ | 1 | -2/3 |
| | $\bar{d}_i$ | $\tilde{d}^*_{Ri}$ | $\bar{d}_i = d^+_{Ri}$ | $\bar{3}$ | 1 | 1/3 |
| Leptons | $L_i$ | $(\tilde{\nu}, \tilde{e}_L)_i$ | $(\nu, e_L)_i$ | 1 | 2 | -1/2 |
| | $\bar{e}_i$ | $\tilde{e}^*_{Ri}$ | $\bar{e}_i = e^+_{Ri}$ | 1 | 1 | 1 |
| Gauge bosons | $G$ | $G^a_\mu$ | $\tilde{g}^a$ | 8 (adj) | 1 | 0 |
| | $W$ | $W^3_\mu, W^\pm_\mu$ | $\tilde{W}^3, \tilde{W}^\pm$ | 1 | 3 (adj) | 0 |
| | $B$ | $B_\mu$ | $\tilde{B}$ | 1 | 1 | 0 |

NOTE: Bosonic superpartners get the name of the SM particles with a prefix s- (eg. electron → selectron). Fermionic superpartners get the name from the SM particles with a suffix -ino (eg. gluon → gluino, Higgs → Higgsino)

We denote by a ~ the superpartners of the usual SM particles; notice that the bar (‾) is part of the names of the conjugates of the right-handed fields. The particle generations are the same as in the SM:

$$u_i = (u, c, t) \qquad d_i = (d, s, b)$$

$$\nu_i = (\nu_e, \nu_\mu, \nu_\tau) \qquad e_i = (e, \mu, \tau)$$

· The Higgs sector needs to be modified with respect to the SM. We need at least twice as many Higgs doublets as in the SM. Let's discuss why this happens.

- If we introduce a supersymmetric Higgs doublet, its fermionic superpartner (the Higgsino) generates $U(1)_Y^3$ and $U(1)_Y\, SU(2)_L^2$ anomalies. To cancel them we need to introduce an extra doublet with opposite hypercharge. We can also include several Higgs doublets, but anomaly cancellation forces us to put them in pairs with opposite hypercharges.

- With two doublets ($H_u$ and $H_d$), the superpotential contains the terms

$$W_{Higgs} = \bar{u}\, Y_u\, Q\, H_u - \bar{d}\, Y_d\, Q\, H_d - \bar{e}\, Y_e\, L\, H_d + \mu\, H_u H_d .$$

The dimensionless Yukawa coupling parameters $Y_{u,d,e}$ are 3x3 matrices in family space. The "$\mu$ term" can be written out as

$$\mu (H_u)_\alpha (H_d)_\beta\, \varepsilon^{\alpha\beta} ,$$

where $\varepsilon^{\alpha\beta}$ is used to tie together $SU(2)_L$ indices in a gauge invariant way. Analogously the term $\bar{u}\, Y_u\, Q H_u$ can be written as

$$\bar{u}^{ia} (Y_u)_i{}^j\, Q_{j\alpha a}\, (H_u)_\beta\, \varepsilon^{\alpha\beta}$$

where $i,j = 1,2,3$ are family indices and $a = 1,2,3$ is a color index.

Notice that in the SM we can write a mass term for the top and the bottom with one Higgs multiplet because we can use $H$ and $H^*$ to write a Yukawa coupling. In a supersymmetric theory this is not possible, because $\bar{H}_u$ and $\bar{H}_d$ can not appear in the superpotential (since it must be holomorphic). Hence we are forced to introduce at least two doublets to give mass to all the fermions.

The third generation of fermions are the heaviest ones, so we can approximate the Yukawa matrices as

$$Y_u \approx \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & y_t \end{pmatrix} \qquad Y_d \approx \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & y_d \end{pmatrix} \qquad Y_e \approx \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & y_\tau \end{pmatrix} .$$

In this limit only the third family contributes to the MSSM superpotential, so:

$$W_{MSSM} \approx y_t (\bar{t}\, t\, H_u^0 - \bar{t}\, b\, H_u^+) - y_b (\bar{b}\, t\, H_d^- - \bar{b}\, b\, H_d^0) - y_\tau (\bar{\tau}\, \nu_\tau H_d^- - \bar{\tau}\,\tau\, H_d^0)$$
$$+ \mu (H_u^+ H_d^- - H_u^0 H_d^0) .$$

When $H_u^0$ and $H_d^0$ get a VEV, the Yukawa couplings generate the masses for the fermions.

The minimal Higgs sector can be summarized in the table

| | bosons | fermions | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ |
|---|---|---|---|---|---|
| $H_u$ | $(H_u^+, H_u^0)$ | $(\tilde{H}_u^+, \tilde{H}_u^0)$ | 1 | 2 | 1/2 |
| $H_d$ | $(H_d^0, H_d^-)$ | $(\tilde{H}_d^0, \tilde{H}_d^-)$ | 1 | 2 | -1/2 |

In a supersymmetric theory the Yukawa-like terms in the superpotential generate not only the usual Yukawas couplings as in the SM (Higgs-quark-quark, Higgs-lepton-lepton), but also interactions involving the superpartners (eg. squark-Higgsino-quark and slepton-Higgsino-lepton). By integrating out the auxiliary $F$ components we also get quartic couplings (eg. (squark)$^4$, (slepton)$^4$, (squark)$^2$(Higgs)$^2$).

(exercise. Derive the interaction vertices.)

The dimensionful coupling in the supersymmetric part of the MSSM Lagrangian are all depending on $\mu$. Integrating out the auxiliary fields we get the Higgs mass terms

$$\mathcal{L}_{\text{higgsino}} = -\mu\left(\tilde{H}_u^+ \tilde{H}_d^- - \tilde{H}_u^0 \tilde{H}_d^0\right) + h.c.$$

and

$$\mathcal{L}_{\text{Higgs}} = -|\mu|^2 \left(|H_u^0|^2 + |H_u^+|^2 + |H_d^0|^2 + |H_d^-|^2\right).$$

The Higgs potential has a quadratic term with a positive squared mass, so there is a stable minimum at the origin with $\langle H_u\rangle = \langle H_d\rangle = 0$. To have electroweak symmetry breaking we have to add <u>soft susy breaking terms</u> in the Lagrangian. A discussion of this aspect will be presented later.

Let's now comment on the size of the "$\mu$ term". To obtain a reasonable scale for the electroweak breaking we need $\mu \sim O(m_{soft}) \sim O(\text{TeV})$ otherwise we should have unnatural cancellations. In principle there is no reason for which $\mu$ should be of TeV order, it could be as well of the order of the Planck mass, so that a hierarchy of scales is reintroduced in the theory. This is known as the <u>$\mu$-problem.</u> Notice that this problem is different from the hierarchy problem in the SM. In the latter the Higgs mass is <u>sensitive</u> to the UV physics, so it is driven to large values by radiative corrections. In supersymmetry the divergences are cancelled, so $\mu$ is only slightly sensitive to the UV physics and receives small radiative corrections. This means that it is "technically natural", or, in other words, once we take it to be of a certain order, its value is not completely changed by radiative corrections, so we do <u>not</u> need a fine tuning of the bare mass against loop corrections. The "$\mu$ problem" is thus only the lack of an explanation for the size of $\mu$.

<u>NOTE.</u> A somewhat similar fact happens for the SM Yukawa couplings. The Yukawas for the light families (eg. the electron) are much smaller than the others (eg. the top). This difference is not explained by the theory. Nevertheless small Yukawas are stable against radiative corrections due to the chiral symmetry in the limit of vanishing Yukawas, so there is no need for a tuning to cancel loop corrections.

· Integrating out the auxiliary fields we also get cubic scalar interactions proportional to $\mu$:

$$\mathcal{L}_{\mu,\text{cubic}} = \mu^*\Big(\tilde{u}_R^* Y_u \tilde{u}_L H_d^{0*} + \tilde{d}_R^* Y_d \tilde{d}_L H_u^{0*} + \tilde{e}_R^* Y_e \tilde{e}_L H_u^{0*} + \tilde{u}_R^* Y_u \tilde{d}_L H_d^{-*} + \tilde{d}_R^* Y_d \tilde{u}_L H_u^{+*} + \tilde{e}_R^* Y_e \tilde{\nu}_L H_u^{+*}\Big) + h.c.$$

<u>NOTE.</u> Various solutions have been proposed for the $\mu$ problem. They work by assuming that the $\mu$ term is absent at tree-level before electroweak symmetry breaking, and then it arises from the VEV of some field. This VEV is generated by a potential which depends on the soft susy-breaking terms, so the value of $\mu$ is related to the value of soft breaking and is naturally of the same order.

· **R-parity and its consequences.**

The superpotential we wrote contains the minimal set of terms needed to produce a phenomenologically viable model. However there are other terms that one can write that are gauge-invariant, holomorphic and renormalizable and could be included in the superpotential. These terms are

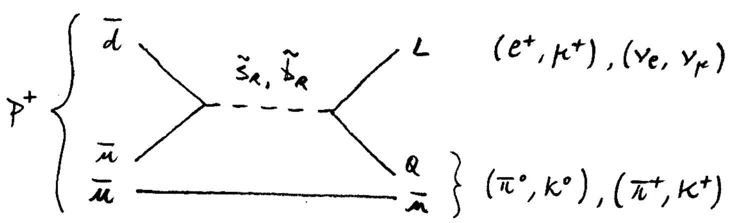
$$W_{\Delta L=1} = \frac{1}{2} \lambda^{ijk} L_i L_j \bar{e}_k + \lambda'^{ijk} L_i Q_j \bar{d}_k + \mu'^{i} L_i H_u$$

$$W_{\Delta B=1} = \frac{1}{2} \lambda''^{ijk} \bar{u}_i \bar{d}_j \bar{d}_k$$

where $i, j, k$ are family indices. Here we consider the usual assignment of baryon and lepton numbers, that is equal numbers for all the components in a superfield with:

- $B = +1/3$ for $Q_i$; $B = -1/3$ for $\bar{u}_i, \bar{d}_i$; $B = 0$ for all the others,
- $L = +1$ for $L_i$; $L = -1$ for $\bar{e}_i$; $L = 0$ for all the others.

The terms in $W_{\Delta L=1}$ violate lepton number by 1, while those in $W_{\Delta B=1}$ violate baryon number by 1. This violation can give rise to proton decays. The decay to a lepton and a meson is given by the Feynman diagram




We can estimate the approximate decay width since the decay probability amplitude is proportional to $\lambda'' \cdot \lambda' / m_{\tilde{q}}^2$, where $m_{\tilde{q}}$ is the mass of the exchanged squark. By dimensional analysis we get

$$\Gamma_p \sim m_p^5 \cdot \frac{|\lambda' \cdot \lambda''|^2}{m_{\tilde{q}}^4} .$$

The proton lifetime can be estimated as

$$\tau_p = \frac{1}{\Gamma_p} \sim \frac{1}{|\lambda' \lambda''|^2} \cdot \left( \frac{m_{\tilde{q}}}{1\,\text{TeV}} \right)^4 \cdot 10^{-12}\,\text{s}.$$

Experimentally, the bound on the proton lifetime is $\tau_p > 10^{32}$ years $\simeq 3 \cdot 10^{39}$ s. So we would need $|\lambda' \lambda''| < 10^{-25}$ (if $m_{\tilde{q}} \sim 1$ TeV), which is an unnaturally small coupling.

To solve this problem one could assume B and L conservation in the MSSM. However this is a step back with respect to the SM, where conservation of these quantum numbers is <u>not</u> assumed; but rather it is an "accidental" symmetry coming from the fact that there are no possible renormalizable terms that violate B or L. Moreover it is known that B and L are violated by non-perturbative effects, so it seem unnatural to impose B and L as symmetries by hand in the MSSM.

A more elegant solution is to introduce a new discrete symmetry in the MSSM. This symmetry is called "matter parity":

Matter parity is a multiplicatively conserved quantum number defined as

$$P_M = (-1)^{3(B-L)}$$

for each particle in the theory. It is easy to check that the quark and lepton supermultiplets have $P_M = -1$, while the Higgs supermultiplets have $P_M = +1$. Gauge fields and gauginos have $B = L = 0$, so they have $P_M = +1$. A term is allowed in the Lagrangian only if it is even under matter parity. It is easy to see that all the terms in $W_{\Delta L=1}$ and $W_{\Delta B=1}$ are forbidden, while all the terms we previously included in the superpotential are allowed.

<u>NOTE</u>. Even if matter parity is an exact symmetry, baryon and lepton number conservation could be violated in the MSSM. However the MSSM does not have renormalizable interactions that violate B or L, if matter parity conservation is assumed.

It is often useful to recast matter parity in terms of R-parity, defined as

$$P_R = (-1)^{3(B-L)+2s},$$

where s is the spin of the particle. Matter parity and R-parity are equivalent, since the product of $(-1)^{2s}$ for the particles involved in an interaction vertex in a theory that conserves angular momentum is always equal to $+1$ (equivalently one can notice that all the terms in the Lagrangian have an even number of fermions).
Particles in the same multiplet do not have the same R-parity, so this symmetry does not commute with supersymmetry, it is an R-symmetry.
R-parity transforms the particles as

$$(\text{SM particle}) \to (\text{SM particle}),$$
$$(\text{superpartner}) \to -(\text{superpartner}).$$

R-parity conservation has a series of extremely important phenomenological consequences:

- The lightest sparticle with $P_R = -1$, called the "lightest supersymmetric particle" or LSP, must be absolutely stable. If the LSP is electrically neutral, it interacts only weakly with ordinary matter and could be a dark matter candidate.
- Each sparticle other than the LSP must eventually decay into a state that contains an odd number of LSP's.
- In collider experiments, sparticles can only be produced in even numbers (usually in pairs).

<u>NOTE</u>. R-parity or matter parity could originate from a gauged U(1) symmetry which is spontaneously broken at high energy to a discrete subgroup.