Supersymmetry Breaking

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Abstract

Supersymmetry has to be a broken symmetry in order to have any relevance for the real world. Like any other symmetry in field theory it could either be broken explicitly by non invariant terms in the Lagrangian or by a spontaneos breaking mechanism. In this report I will recall supersymmetric theories of chiral supermultiplets interacting via gauge supermultiplets, show that supersymmetry cannot be an exact symmetry of the known particle spectrum and discuss spontaneous breaking mechanisms that could lead to the effective symmetry breaking Lagrangians appearing in the MSSM.

Contents

1 Introduction

The standard model of elementary particles and interactions employs three different kinds of fields on space time: bosonic gauge fields (observed) , bosonic scalar Higgs fields (not observed) as well as fermionic (Weyl) spinor fields (observed).

Assuming that besides Poincaré symmetry and the internal $SU(3)\otimes SU(2)\otimes U(1)$ gauge symmetry group there exists an additional symmetry of Nature transforming bosons to fermions and vice versa, forces us to extend the particle spectrum by superpartners for every particle in the standard model. Since particles and their superpartners are equal in mass and every gauge quantum number, they should have been observed long time ago. Therefore supersymmetry cannot be a perfect symmetry of the particle spectrum, but has to be a broken symmetry if it has any phenomenological relevance at all .

This report aims at giving a hopefully understandable introduction to the issue of supersymmetry breaking presumably occuring in the supersymmetric extension of non-abelian gauge theories. I will mainly follow the notation and arguments given in reference [1,2,3].

2 Supersymmetric extended gauge theories

To study supersymmetry breaking in a comfortable environment I will first introduce the supersymmetric extension of gauge theories. Besides Lorentz invariance such a theory should obey the following requirements:

1. It should be invariant under supersymmetric field transformations pairing the massive fermions or the gauge bosons with a superpartner having the same number of degrees of freedom.
2. The theory should remain gauge invariant.
3. The interactions of the extended theory should stay renormalizable.

2.1 Chiral Multiplets

In supersymmetry the matter fields of the SM are replaced by so called chiral supermultiplets. They consist of the fermionic Weyl spinor fields χ^i_α, their complex scalar superpartners ϕ^i and complex auxiliary fields F^i required to close the superalgebra off shell in perturbation theory.

- Chiral Supermultiplets $\left\{\phi^i, \chi^i_\alpha, F^i\right\}, \alpha = 1, 2$:

$$\begin{aligned} \delta_\xi \phi^i &= \xi \cdot \chi^i, & (1) \\ \delta_\xi \chi_{i\,\alpha} &= -i(\sigma^\mu \xi^\dagger)_\alpha \partial_\mu \phi^i + \xi_\alpha F^i, & (2) \\ \delta_\xi F^i &= -i(\xi^\dagger \overline{\sigma}^\mu) \cdot (\partial_\mu \chi^i), & (3) \end{aligned}$$

where the global infinitessinmal parameter ξ is itself an anticommuting spinor.

- Chiral Lagrangian:

$$L_{chiral} = L_{free} + L_{int}, \tag{4}$$
$$L_{free} = -\partial_\mu \phi^{\dagger i} \partial^\mu \phi_i + \chi^{\dagger i} i \bar{\sigma}^\mu \partial_\mu \chi_i + F^{\dagger i} F_i, \tag{5}$$
$$L_{int} = (-\frac{1}{2} \frac{\partial^2 W}{\partial \phi_i \partial \phi_j} \chi_i \chi_j + \frac{\partial W}{\partial \phi_i} F_i) + h.c., \tag{6}$$

where the so called superpotential W is given by

$$W = \frac{1}{2} M_{ij} \phi_i \phi_j + \frac{1}{6} y_{ijk} \phi_i \phi_j \phi_k. \tag{7}$$

The classical equation of motion for the auxiliary field gives $F^{\dagger i} = -(\frac{\partial W}{\partial \phi_i})$ and $F_i = -(\frac{\partial W}{\partial \phi_i})^\dagger$. The Lagrangian desity L_{chiral} transforms under supersymmetry transformations as

$$\delta L_{chiral} = \frac{1}{2} i \partial_\mu (...) \tag{8}$$

and therefore by the usual argument of boundary conditions the action S remains invariant. Since by Lorentz invariance only scalar fields are able to have a non zero vacuum expectation value (VeV) a useful quantity to consider is the scalar potential $V(\phi, \phi^\dagger)$ given by

$$V_{chiral} = F^{\dagger i} F_i \tag{9}$$
$$= (\frac{\partial W}{\partial \phi_i})(\frac{\partial W}{\partial \phi_i})^\dagger \tag{10}$$

Additionally one encounters fermion mass terms $(-\frac{1}{2} M^{ij} \psi_i \psi_j + h.c.)$ and Yukawa type interactions (-$\frac{1}{2} y^{ijk} \phi_i \psi_j \psi_k + h.c$).

2.2 Gauge Interactions

In analogy to standard QED the theory becomes interacting when we promote the global internal symmetry of the free Lagrangian to a space-time dependent gauge symmetry. In the supersymmetric invariant extension this leads to a gauge supermultiplet $G = \{A^{\mu a}, \lambda^a, D^a\}$, where the index a runs over the adjoint representation of the gauge group in consideration. The massless spin 1 gauge boson $A^{\mu a}$ is paired with a massless Weyl fermion λ^a and a real bosonic auxilliary field D^a. Under supersymmetry a gauge supermultiplet transforms as:

$$\delta_\xi A^{\mu a} = -\frac{1}{\sqrt{2}} \left(\xi^\dagger \overline{\sigma}^\mu \lambda^a + \lambda^{a\dagger} \overline{\sigma}^\mu \xi \right) \tag{11}$$
$$\delta_\xi \lambda^a = \frac{i}{2\sqrt{2}} \sigma^\mu \overline{\sigma}^\nu \xi F^a_{\mu\nu} + \frac{1}{\sqrt{2}} \xi D^a \tag{12}$$
$$\delta_\xi D^a = \frac{i}{\sqrt{2}} \left(\xi^\dagger \overline{\sigma}^\mu D_\mu \lambda^a - D_\mu \lambda^{a\dagger} \overline{\sigma}^\mu \xi \right) \tag{13}$$

Under the gauge group the supermultiplet behaves as:

$$\begin{aligned} \delta_{gauge} A^{\mu a} &= \partial_\mu \Lambda^a + g f^{abc} A^b_\mu \Lambda^c && (14)\\ \delta_{gauge} \lambda^a &= g f^{abc} \lambda^b \Lambda^c, && (15)\\ \delta_{gauge} D^a &= \partial_\mu \Lambda^a + g f^{abc} A^b_\mu \Lambda^c && (16) \end{aligned}$$

where the structure constants f^{abc} of the gauge Lie algebra are defined by $\left[T^a, T^b\right] = i f^{abc} T^c$. This is exactly what is ment by "gauge fields transform in the adjoint representation of the gauge group"; the transformation matrix for the gauge supermultiplet is given by the structure constants of the Lie algebra. The free Lagrangian for the gauge supermultiplet is

$$L_{gauge} = -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu a} + i \lambda^{\dagger a} \overline{\sigma}^\mu D_\mu \lambda^a + \frac{1}{2} D^a D^a \tag{17}$$

where the field strength tensor $F^a_{\mu\nu}$ is given by

$$F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + g f^{abc} A^b_\mu A^c_\mu, \tag{18}$$

and the covariant derivative for the Weyl fermion is

$$D_\mu \lambda^a = \partial \lambda^a + g f^{abc} A^b_\mu \lambda^c. \tag{19}$$

Combining chiral and gauge supermultiplets in a common interacting theory is achieved by gauge transformations of the chiral supermultiplets:

$$\begin{aligned} \delta_{gauge} X_i &= i g \Lambda^a ((T^a)^j_i X_j) && (20)\\ D_\mu X_i &= \partial_\mu X_i - i g A^a_\mu (T^a)^j_i X_j && (21) \end{aligned}$$

where $X_i \epsilon \left\{\phi^i, \chi^i_\alpha, F^i\right\}$. This leads to the desired interaction terms of the form $i g A^a_\mu (T^a)^j_i X_j$. Furthermore one can construct additional terms that couple the gaugino λ^a or the auxiliary field D^a to the fields of the chiral supermultiplet. Renormalizabilty again restricts the form of these tems i.e.:

$$(\phi^{\dagger i} T^{aj}_i \chi_j) \lambda^a, \; h.c. \; (\phi^{\dagger i} T^{aj}_i \phi_j) D^a \tag{22}$$

Adding them to the Lagrangian and demanding supersymmetry invariance i.e. $\delta_\xi S = 0$ forces us to modify the transformations of the chiral supermultiplet:

$$\begin{aligned} \delta_\xi \phi^i &= \xi \cdot \chi^i, && (23)\\ \delta_\xi \chi_{i\alpha} &= -i(\sigma^\mu \xi^\dagger)_\alpha D_\mu \phi^i + \xi_\alpha F^i, && (24)\\ \delta_\xi F^i &= -i(\xi^\dagger \overline{\sigma}^\mu) \cdot (D_\mu \chi^i) + \sqrt{2} g (T^{ai}_j \phi^j) \xi^\dagger \lambda^{\dagger a} && (25) \end{aligned}$$

Finally the total Lagrangian of the interacting supersymmetric theory takes the form:

$$\begin{aligned} L &= L_{chiral} + L_{gauge} && (26)\\ & \quad - \sqrt{2} g (\phi^{\dagger i} T^{aj}_i \chi_j) \lambda^a + h.c. + g (\phi^{\dagger i} T^{aj}_i \phi_j) D^a && (27) \end{aligned}$$

where all ordinary derivatives have been replaced by the covariant one i.e.

$$L_{chiral} = -D_\mu \phi^{\dagger i} D^\mu \phi_i + \chi^{\dagger i} i\bar{\sigma}^\mu D_\mu \chi_i + F^{\dagger i} F_i \quad (28)$$

$$+(-\frac{1}{2}\frac{\partial^2 W}{\partial\phi_i \partial\phi_j}\chi_i\chi_j + \frac{\partial W}{\partial\phi_i}F_i) + h.c. \quad (29)$$

The scalar potential V is now

$$V = F^{\dagger i} F_i + \frac{1}{2}\sum_a D^a D^a \quad (30)$$

$$= \sum_i \frac{\partial W}{\partial\phi_i}\left(\frac{\partial W}{\partial\phi_i}\right)^\dagger + \frac{1}{2}\sum_a g_a^2 (\phi^{\dagger i} T_i^{aj} \phi_j)^2 \quad (31)$$

where $\frac{\partial W}{\partial\phi_i} = M^{ij}\phi_j + \frac{1}{2}y^{ijk}\phi_j\phi_k$. Before coming to the main part of this report I briefly collect all possible gauge and non gauge interaction of the supersymmetric theory.

- Non-gauge dimensionless interactions:
  - Yukawa interaction:
  $$-\frac{1}{2}y^{ijk}\chi_i\chi_j\phi_k + h.c. \quad (32)$$
  - quartic scalar interaction:
  $$+\frac{1}{4}y^{ijn}y^*_{kln}\phi_i\phi_j\phi^{\dagger k}\phi^{\dagger l} \quad (33)$$
- Non-gauge dimensionful interactions:
  - scalar^3 interaction:
  $$\frac{1}{2}M^{in}y^*_{jkn}\phi_i\phi^{\dagger j}\phi^{\dagger k} + h.c. \quad (34)$$
  - scalar squared mass terms:
  $$M^*_{ik}M^{kj}\phi^{\dagger i}\phi_j \quad (35)$$
  - fermion mass terms:
  $$-\frac{1}{2}M^{ij}\chi_i\chi_j + h.c. \quad (36)$$
- Gauge interactions:
  - non abelian boson-boson interactions coming from
  $$-\frac{1}{4}F^a_{\mu\nu}F^{\mu\nu a} \quad (37)$$

– boson-fermion and boson-scalar interactions coming from

$$-D_\mu \phi^{\dagger i} D^\mu \phi_i + \chi^{\dagger i} i \bar{\sigma}^\mu D_\mu \chi_i + i\lambda^{\dagger a} \overline{\sigma}^\mu D_\mu \lambda^a \tag{38}$$

where the covariant derivative is either given by

$$D_\mu \lambda^a = \partial \lambda^a + g f^{abc} A^b_\mu \lambda^c. \tag{39}$$

or

$$\begin{aligned} D_\mu \phi_i &= \partial_\mu \phi_i - ig A^a_\mu (T^a)^j_i \phi_j & (40) \\ D_\mu \chi_i &= \partial_\mu \chi_i - ig A^a_\mu (T^a)^j_i \chi_j & (41) \end{aligned}$$

– gaugino-scalar-fermion interactions:

$$-\sqrt{2} g (\phi^{\dagger i} T^{aj}_i \chi_j) \lambda^a + h.c \tag{42}$$

– scalar quartic interaction:

$$g^2_a (\phi^{\dagger i} T^{aj}_i \phi_j)^2 \tag{43}$$

3 Supersymmetry breaking

As mentioned in the introduction supersymmetry has to be a broken symmetry of the present universe. This can be seen from a simple argument. The fermionic generators of supersymmetry transformations, transforming the bosonic states of a supermultiplet to fermionic ones and vice versa

$$Q_\alpha \left| fermion \right\rangle = \left| boson \right\rangle, \quad Q_\alpha \left| boson \right\rangle = \left| fermion \right\rangle, \tag{44}$$

fulfill the (anti-) commutation relations

$$\begin{aligned} \left\{ Q_\alpha, Q^\dagger_{\dot\beta} \right\} &= -2(\sigma^\mu)_{\alpha\dot\beta} P_\mu & (45) \\ \left\{ Q_\alpha, Q_\beta \right\} &= 0 = \left\{ Q^\dagger_{\dot\alpha}, Q^\dagger_{\dot\beta} \right\} & (46) \\ \left[Q_\alpha, P_\mu \right] &= 0 = \left[Q^\dagger_{\dot\alpha}, P_\mu \right] & (47) \end{aligned}$$

This justifies the calculation:

$$\begin{aligned} m^2_b \left| boson \right\rangle &= P^2 \left| boson \right\rangle = Q_\alpha P^2 \left| fermion \right\rangle = m^2_f \left| boson \right\rangle & (48) \\ \Rightarrow m^2_b &= m^2_f & (49) \end{aligned}$$

So we conclude that superpartners have equal mass in unbroken supersymmetry. Furthermore consider that the interactions of a supersymmetric extended gauge theory, summarized in the last part of the report couple particles to superparticles. Therefore superpartners should have inevitably been detected so far in collider experiments. Another important feature of unbroken supersymmetry is the zero vacuum energy. To see this consider

$$\begin{aligned}\langle 0\left|P^0\right|0\rangle &= \frac{1}{4}\left\langle 0\left|\left\{Q_1,Q_1^\dagger\right\}+\left\{Q_2Q_2^\dagger\right\}\right|0\right\rangle &(50)\\ &= \frac{1}{4}\left(\left|\left|Q_1^\dagger\left|0\right\rangle\right|\right|^2+\left|\left|Q_1\left|0\right\rangle\right|\right|^2+\left|\left|Q_2^\dagger\left|0\right\rangle\right|\right|^2+\left|\left|Q_2^\dagger\left|0\right\rangle\right|\right|^2\right) &(51)\end{aligned}$$

which is clearly equal to zero in unbroken supersymmetry where the generators Q_α annihilate the vacuum i.e. $Q_\alpha\left|0\right\rangle=0$. Conversely if supersymmetry is spontaneously broken $Q_\alpha\left|0\right\rangle\neq 0$ by the same argument the vacuum energy is strictly larger than zero.

So we are left with three possibilities, either supersymmetry is broken explicitly by adding non supersymmetric invariant terms to the gauge theory Lagrangians, a spontaneous breaking mechanism can be identified or it is just an analaptic mind game. Dismissing the last possibility we will see that a possible solution to this problem could lie in a combination of the first two alternatives.

3.1 Explicit soft supersymmetry breaking

The idea behind explicit supersymmetry breaking is that supersymmetry is in fact broken spontaneously at a very high energy scale that up to now has not been experimentally accessed. Therefore our ignorance of the breaking mechanism is parametrized by effective Lagrangians. This is done by taking the supersymmetric invariant Lagrangian density of the extended gauge theory designed to describe elementary particle interactions and include additional terms designated with L_{soft} that manifestly are not supersymmetric invariant. Since we still want a renormalizable theory and in fact one that will continue to cure the hirarchy problem the form of these terms is restricted and one calls them "soft". To summarize:

$$\begin{aligned}L &= L_{SuSy}+L_{soft} &(52)\\ \delta_\xi\int dx^4 L_{soft} &\neq 0 &(53)\end{aligned}$$

where the almost most general form of L_{soft} is given by:

$$L_{soft}=-\left(\frac{1}{2}M_a\lambda^a\lambda^a+\frac{1}{6}a^{ijk}\phi_i\phi_j\phi_k+\frac{1}{2}b^{ij}\phi_i\phi_j\right)+h.c.-(m^2)^i_j\phi^{j\dagger}\phi_i \quad (54)$$

If this is done for the MSSM we are left with more than hundred adjustable parameters as coupling constants, particle masses etc. . Therefore the identification of a consistent supersymmetry breaking mechanism that could be used to determine these parameters is one of the more important questions (besides observation of the Superpartners of course) in particle physics.

3.2 Spontaneously broken symmetries in field theory

In order to break an internal symmetry (in our case supersymmerty) of a quantum field theory it is necessary that a field, which is not invariant under the symmetry i.e. is member of a symmetry multiplet (in our case supermultiplet) aquires a non-vanishing vacuum expectation value (VeV). Consider a continuous field transformation with group generators Q_a.

$$\phi(x) \to \phi'(x) = \phi(x) + i\theta^a[(Q_a), \phi(x)],$$

Spontaneous symmetry breaking occurs if

$$\int d^4x \frac{\delta L[\phi']}{\delta\phi(x)}(Q_a)\phi(x) = 0 \, and \ \langle\Omega| \, \phi'(x) \, |\Omega\rangle \neq 0,$$

This implies that the vacuum is not invariant under the symmetry $Q_a \, |\Omega\rangle \neq 0$. In the case of SuSy we will have something like

$$[Q, \phi(x)] = i\psi(x), \quad \langle\Omega| \, \psi(x) \, |\Omega\rangle \neq 0,$$

where $\phi(x)$ and $\psi(x)$ are fields of different spin type belonging to the same supermultiplet. From $\langle\Omega| \, \psi(x) \, |\Omega\rangle = \langle\Omega| \, [Q, \phi(x)] \, |\Omega\rangle \neq 0$ thereby follows:

$$Q \, |\Omega\rangle \neq 0 \tag{55}$$

Since the VeV has to be Lorenz invariant the field $\psi(x)$ that is responsible for the symmetry breaking has to be a scalar member of the multiplet. It has already been mentioned that the vacuum energy now is strictly positive: $Q_a \, |0\rangle \neq 0 \Rightarrow \langle 0 \left| P^0 \right| 0 \rangle > 0$. Spacetime independence and Lorentz invariance of the vacuum energy implies that its value is given by the scalar potential:

$$\begin{aligned} \langle 0 \left| P^0 \right| 0 \rangle &= \langle 0 \left| V \right| 0 \rangle \\ V &= F^{*i}F_i + \frac{1}{2}\sum_k D^{a_k} D^{a_k} \end{aligned}$$

where the sum over k indicates the possibility of more than one gauge group.

Therefore only two kind of scalar fields could aquire the non-vanishing VeV in order to break supersymmetry. These are called F-type or D-type breaking:

$$\langle 0 \left| F \right| 0 \rangle \neq 0, \ or \ \langle 0 \left| D^a \right| 0 \rangle \neq 0,$$

This sets the procedure for the construction of spontaneous supersymmetry breaking mechanisms: Look for models where $F_i = 0$ *and* $D^a = 0$ simultaneously is impossible for any field configuration.

3.3 Fayet-Iliopoulos D-term breaking

For theories with a abelian $U(1)$ gauge (sub)group Pierre Fayet and John Iliopoulos proposed a breaking mechanism involving the auxiliary filed D of the gauge supermultiplet. Let us introduce $L_{F-I} = -\kappa D$, with $[\kappa] = M^2$ and $\delta_\xi D = \partial_\mu(...)$ to the Lagrangian. The the D-term involving part of the scalar potential then becomes:

$$V|_D = \kappa D - \frac{1}{2}D^2 - gD\sum_i q_i \left|\phi_i\right|^2,\ D = \kappa - g\sum_i q_i \left|\phi_i\right|^2,$$

$$\Rightarrow V = \sum_i \left|m_i\right|^2 \left|\phi_i\right|^2 + \frac{1}{2}\left(\kappa - g\sum_i q_i \left|\phi_i\right|^2\right)^2.$$

We see that this is clearly larger than zero $V > 0$ for every field configuration. This implies that $Q_a \left|0\right\rangle \neq 0$ and we find that the mass degeneracy is removed: The scalar masses become $\left|m_i\right|^2 - gq_i\kappa$,whereas the fermion masses stay $\left|m_i\right|^2$.

3.4 F-type supersymmetry breaking

We allready mentioned that due to Lorentz invariance only the auxilliary fields appearing in the scalar potential of a supersymmertic extended gauge theory could eventually aquire a non-vanishing VeV. So it remains to consider symmerty breaking caused by the F_i fields of the chiral supermultiplets. The procedure for constructing such models is the following:

1. Take set of chiral supermultiplets $\left\{\phi^i, \chi^i_\alpha, F^i\right\}$, $i = 1, ..., n$
2. Construct superpotential W such that $F_i = -(\frac{\partial W}{\partial \phi_i})^\dagger = 0\ \forall i$ is impossible.
3. Then $V = \sum_i \left|F_i\right|^2 > 0$ for every field configuration.
4. Then $Q_a \left|0\right\rangle \neq 0$.

To see how this works let us consider a simple model conceived by the Irish physicist Lochlainn O'Raifeartaigh.

3.4.1 The O'Raifeartaigh model

Consider a theory with three chiral supermultiplets$\left\{\phi^i, \chi^i_\alpha, F^i\right\}$, $i = 1, ..., 3$. The superpotential W is constructed to be

$$W = -k\phi_1 + m\phi_2\phi_3 + \frac{g}{2}\phi_1\phi_3^2,\ k, m, g > 0, \tag{56}$$

where the term linear in ϕ_1 is crucial for the spontaneous breaking to occur. The scalar potential is now given by:

$$V = \left(k - \frac{g}{2}\left|\phi_3\right|^2\right)^2 + m^2\left|\phi_3\right|^2 + \left|m\phi_2 + g\phi_1\phi_3\right|^2 > 0,$$

We imediatly see that this cannot be equal to zero. For $m^2 > gk$ it is clearly minimized by: $\phi_3 = 0 = \phi_2$, where the scalar ϕ_1 remains undetermined . This is called a flat direction in field space. So without loss of generality ϕ_1 can also be taken to zero. Since $V = \sum_i |F_i|^2$ we see that the non vanishing VeV for the auxiliary field is $\langle F_1 \rangle = k$. If we expand the scalar potential of the three chiral supermultiplets around the minimum field configuration i.e. $\phi_i = 0$ we get

$$V = V(0) + m^{ij}\phi_i\phi_j + O(\phi^3), \tag{57}$$

with mass matrix

$$m^{ij} = \frac{\partial^2 V}{\partial\phi_i\partial\phi_j}|_{\phi_0} = m^{ji} \tag{58}$$

Doing this one gets 6 real scalar field with masses 0, 0, m^2, m^2, $m^2 - gk$, $m^2 + gk$, and also 3 Weyl fermions with masses (coming from $\frac{1}{2}\frac{\partial^2 W}{\partial\phi_i\partial\phi_j}\chi_i\chi_j$): 0, m^2, m^2, where the zero mass corresponds to the field χ_1 of the multiplet $\{\phi_1, \chi_1, F_1\}$, which can be interpreted as the massless Goldstone spinor corresponding to the broken generator Q_α.

3.5 Is a seperate sector responsible for supersymmetry breaking?

So far we have seen that for theories with a non-abelian gauge symmetry group like QCD or the Standard Model, supersymmetry breaking is supposed to be of F-type. Demanding that gauge invariance is preseverd by the breaking mechanism and noting that it relied on a term linear in a scalar field we conclude that supersymmetry breaking is achieved by a gauge singlet. This forces us to extend the MSSM by a hidden sector of new particles.

Besides the nature of this new hypothetical sector the main question then is: How is the breaking in the new sector communicated to the MSSM particles to abolish mass degeneracy? Considerations in this direction come under the name of symmetry breaking mediation. The idea is the following: The terms in L_{soft} should emerge by very weak couplings of the MSSM fields to a non vanishing F-term VeV in the hidden sector.

Since this is a highly speculative topic theorist came up with many different, almost arbitrarly sophisticated models of breaking mediation. In this report I only mention two of them, while illustrating one of them with an example. The historically most prominent models for breaking mediation are

1. Gauge mediated supersymmetry breaking:

   The terms in L_{soft} are generated by loop diagrams involving new chiral supermultiplets that couple to the F-term VeV of a gauge singlet in the hidden sector. These so called masseger particles should also have the gauge interactions of the MSSM to mediate the breaking to the visible part of the universe.

2. Gravity or Planck scale medited supersymmetry breaking:

   Symmetry breaking is associated with the "new physics" or "quantum gravity" that is supposed to enter near the Planck scale where the Compton wavelength $\frac{h}{mc}$ of a particle becomes compareable to its Schwarzschild radius $\frac{2Gm}{c^2}$.

3.5.1 Example: Gravity mediated supersymmetry breaking

As an example for breaking mediation consider the non-renormalizable Lagrangian designed to describe interactions of the auxiliary field F of the "new physics" sector with the gauginos and sparticles of the MSSM:

$$\begin{aligned} L_{NR} &= -\frac{1}{M_p} F \left(\frac{1}{2} d_a \lambda^a \lambda^a + \frac{1}{6} y'^{ijk} \phi_i \phi_j \phi_k + \frac{1}{2} \mu'^{ij} \phi_i \phi_j \right) \\ &\quad + h.c. - \frac{1}{M_p^2} F F^* k^i_j \phi_i \phi^{\dagger j} \end{aligned}$$

If now F aquires a VeV $\langle 0 | F | 0 \rangle \neq 0$, L_{soft} emerges:

$$L_{soft} = -\frac{1}{2} M_a \lambda^a \lambda^a - \frac{1}{6} a^{ijk} \phi_i \phi_j \phi_k - \frac{1}{2} b^{ij} \phi_i \phi_j - (m)^{2i}_j \phi^{j*} \phi_i$$

where $M_a = \frac{\langle 0 | F | 0 \rangle}{M_p} d_a$ for example. Although this example is somewhat artificial it clearly illustrates the spirit of breaking mediation.

4 Summary

To summarize the main assertions of this report:

- SuSy breaking is inevitable.
- so far effective soft Lagrangians parametrize supersymmetry breaking in the MSSM.
- Spontaneous F-type breaking in a hidden sector of particles is proposed.
- Mediation of supersymmetry breaking to the MSSM sector is a wide field.

5 References

[1] John Iliopoulos: Beyound the standard model. Lecture given at ENS Paris

[2] Ian Aitchison: Supersymmetry in particle physics

[3] Stephen P. Martin: A Supersymmerty primer (hep-ph/9709356)