5 The minimal supersymmetric SM (MSSM)

5.1 Supersymmetrizing the SM

- 3 (sets of) real SFs V_i ($i = 1,2,3$ for U_1, SU_2, SU_3)
- chiral SFs $\underbrace{Q}_{\text{l.h.}}\,,\ \underbrace{U\,,\ D}_{\text{r.h.}}\,,\ \underbrace{L}_{\text{l.h.}}\,,\ \underbrace{E}_{\text{r.h.}}\,,\ \underbrace{H_u\,,\ H_d}_{\text{two Higgs doublets}}$

  quarks leptons

U_1-charges: $1/6$ $-2/3$ $1/3$ $-1/2$ 1 $1/2$ $-1/2$

- <u>lagrangian:</u> gauge: $\sum_{i=1}^{3}\left(\frac{1}{2g_i^2}\,\mathrm{tr}\,W^2\Big|_{\theta^2} + \text{h.c.}\right)$

  Kähler: $Q^\dagger e^{2V} Q + U^\dagger e^{2V} U + \cdots$

  superpot.: $W = \mu H_u H_d$ ("μ-term")

  (Yukawa couplings) → $+\lambda_u Q H_u U + \lambda_d Q H_d D + \lambda_e L H_d E$

  (extra terms) → $+ a L H_u + b QLD + c UDD + d LLE$

  [only dim. 3 & 4 operators]

- <u>Note:</u> The second Higgs doublet, H_d, is needed since, in contrast to the SM, we can not write terms like $Q\bar{H}_u D$; $L\bar{H}_u E$ because of holomorphicity.

- To forbid the "extra" operators (which induce fast proton decay) we introduce <u>R-symmetry</u>.

$U_{1,R}$: $\phi(y,\theta) \longrightarrow e^{in\alpha}\phi(y, e^{-i\alpha}\theta)$ [$n = R(\phi)$ is the "R-charge" of the SF ϕ.]

in components: $A \to e^{in\alpha} A$

$\psi \to e^{i(n-1)\alpha}\psi$

$F \to e^{i(n-2)\alpha} F$

- Invariance of terms in the Lagrangian:
  - $\phi^+\phi|_{\theta^2\bar\theta^2}$ is invariant for any n
  - $\phi^2|_{\theta^2} \longrightarrow e^{2in\alpha}e^{-2i\alpha}\phi^2|_{\theta^2}$ is invariant only for $n=1$
  - more generally: $W(\phi)|_{\theta^2}$ is invariant if the scalar component of $W(\phi)$ has R-charge 2
  - since e^{2V} appears in the Lagrangian, $R(V)=0$.

    $\Rightarrow R(W_\alpha) = R(\bar D^2 D_\alpha V) = 2-1+0 = 1$

    $\Rightarrow W^\alpha W_\alpha|_{\theta^2}$ is invariant.

- With the assignment

	Q	U	D	L	E	H_u	H_d
R:	1	1	1	1	1	0	0

, the only terms of our $W(Q,U,\cdots)$ which are allowed are the Yukawa couplings.

- This is inacceptable since we need a non-zero μ-term to give mass to the Higgsinos

- Way out: Together with SUSY-breaking, the U_1 R-symm. will be broken to its Z_2 subgroup (of the transformations $e^{i\alpha}$ ($\alpha \in [0,2\pi)$), only those with $\alpha=0$ & $\alpha=\pi$, i.e. 1 & -1, survive as symmetries).

- An independent reason for breaking $U_{1,R}$ to $Z_{2,R}$ is the need for a gaugino mass term $\sim (\lambda^a)^\alpha(\lambda^a)_\alpha$.
  (λ_α is the lowest component of W_α and hence has R-charge 1 $\Rightarrow$ $\lambda^\alpha\lambda_\alpha$ is not inv. under $U_{1,R}$ but <u>is</u> invariant under $Z_{2,R} \subset U_{1,R}$.)

- Transformations under $Z_{2,R}$: Higgs scalars $\to$ + Higgs scalars
  Higgsinos $\to$ − Higgsinos
  fermions $\to$ + fermions
  sfermions $\to$ − sfermions
  gauge bosons $\to$ + gauge bosons
  gauginos $\to$ − gauginos

- This Z_2-symm. forbids all terms in $W(Q, U, \cdots)$ except for the Yukawa couplings & the μ-term

- In addition, the lightest superpartner (LSP) will be absolutely stable since its decay to SM-particles would violate $Z_{2,R}$ ("R-parity")

- Next, we need to break SUSY to give large masses to the superpartners (except the Higgsinos, which already have a mass because of the μ-term).

5.2 SUSY Breaking

- Spont. Breaking : S (or $\hat{H}$) − SUSY-inv.
  φ_0 (or $|0\rangle$) − not SUSY-inv.
  $\uparrow$
  (classical ground state)

F-term breaking

- recall:

$$\delta_\xi A = \sqrt{2}\, \xi \psi$$
$$\delta_\xi \psi = i\sqrt{2}\, \sigma^\mu \bar{\xi} \partial_\mu A + \sqrt{2}\, \xi F$$
$$\delta_\xi F = i\sqrt{2}\, \bar{\xi} \bar{\sigma}^\mu \partial_\mu \psi$$

- Poinc.-invariance $\Rightarrow$ $\psi_0 = 0$; $\partial_\mu A_0 = 0$

$\Rightarrow$ r.h. side can only be non-zero if $F_0 = \text{const.} \neq 0$.

⇒ $(A_0, \psi_0, F_0) = (0, 0, \text{const.})$ is the only way in which a chiral multiplet can break SUSY in a Lorentz-inv. way. SUSY is broken in this vacuum since $\delta_\xi \psi_0 = \sqrt{2}\,\xi F_0 = \sqrt{2}\,\xi\,\text{const} \neq 0$

- The simplest renormalizable model of this type is the

O'Raifeartaigh model:

$$\mathcal{L} = \sum_{i=1}^{3} \bar{\Phi}_i \Phi_i \Big|_{\theta^2\bar{\theta}^2} + \Big[\Phi_1 \left(M^2 + \frac{\lambda}{2}\Phi_3^2\right) + \mu \Phi_2 \Phi_3\Big]\Big|_{\theta^2} + \text{h.c.}$$

scalar potential:

$$V(\phi, \bar{\phi}) = \sum_i |F_i|^2 \underset{\uparrow}{=} \sum_i \left|\frac{\partial W}{\partial \phi_i}\right|^2$$

Problem: check that this follows from the EOMs for F_i

$$\Rightarrow V = \left|M^2 + \frac{\lambda}{2}\phi_3^2\right|^2 + |\mu\phi_2 + \lambda\phi_1\phi_3|^2 + |\mu\phi_3|^2$$

This is minimized by $\phi_2 = \phi_3 = 0$ and any value of ϕ_1 (which is a "flat direction"). Hence, $F_1 = M^2$ in the vacuum.

A simpler model with F-term breaking

(more generally, chiral SF models of this type are sometimes also referred to as O'Raifeartaigh models)

- Consider $\mathcal{L} = \bar{\Phi}\Phi\big|_{\theta^2\bar{\theta}^2} + c\,\Phi\big|_{\theta^2} + \text{h.c.}$
- Component form: $\mathcal{L} = F\bar{F} + cF + \text{h.c.} + \cdots$

$$\Rightarrow F \neq 0$$

- However, SUSY is not really broken since this is a free theory with degenerate boson & fermion.

- Introducing interactions in the form $\phi^2|_{\theta^2}$ or $\phi^3|_{\theta^2}$ does not work since since then the linear term can be absorbed in a shift of ϕ and SUSY will again be unbroken.

- This problem can be overcome by introducing interactions via $(\bar\phi\phi)^2|_{\theta^2\bar\theta^2}$ which, however makes the model non-renormalizable. If accept this, we have a nice model with F-term breaking:

$$\mathcal{L} = [\bar\phi\phi + (\phi\bar\phi)^2]|_{\theta^2\bar\theta^2} + \phi|_{\theta^2} + h.c.$$

(all prefactors are set to 1 for simplicity)

- Ignoring derivatives and fermions, we find

$$\mathcal{L} = F\bar F + 4F\bar F\phi\bar\phi + F + \bar F + \cdots$$

$$\frac{\delta}{\delta\bar\phi}\cdot\cdot \Rightarrow \phi = 0$$

$$\frac{\delta}{\delta\bar F}\cdot\cdot \Rightarrow F + 1 = 0 \Rightarrow F = -1 \neq 0$$

<u>A brief excursion</u>

The full lagrangian of such more general models of the type $\mathcal{L} = K(\phi^i, \bar\phi^{\bar j})|_{\theta^2\bar\theta^2} + W(\phi^i)|_{\theta^2} + h.c.$

can be given in an elegant form using the definitions

$$g_{i\bar j} = \frac{\partial}{\partial\phi^i}\frac{\partial}{\partial\bar\phi^{\bar j}} K(\phi,\bar\phi)$$ (Here we treat $\phi^i, \bar\phi^{\bar j}$ as complex variables, <u>not</u> SFs)

$$\Gamma_{ik}{}^{m} = g^{m\bar{j}}\, g_{i\bar{j},k} \qquad \left((\cdots)_{,k} = \partial_k(\cdots) = \frac{\partial}{\partial\phi^k}(\cdots)\right)$$

$$R_{i\bar{j}k\bar{l}} = g_{m\bar{l}}\, \partial_{\bar{j}}\, \Gamma_{ik}{}^{m}$$

(Kähler metric, Kähler geometry, hence: Kähler potential)

$$\Rightarrow \mathcal{L} = -g_{i\bar{j}}\,\partial_\mu\phi^i\partial^\mu\bar{\phi}^{\bar{j}} - i g_{i\bar{j}}\,\bar{\psi}^{\bar{j}}\bar{\sigma}^\mu D_\mu\psi^i$$
$$+\frac{1}{4}R_{i\bar{j}k\bar{l}}(\psi^i\psi^k)(\bar{\psi}^{\bar{j}}\bar{\psi}^{\bar{l}}) - \frac{1}{2}(D_iD_jW)\psi^i\psi^j + h.c.$$
$$-g^{i\bar{j}}(D_iW)(D_{\bar{j}}\bar{W})$$

where $D_\mu\psi^i = \partial_\mu\psi^i + \Gamma^i_{jk}(\partial_\mu\phi^j)\psi^k$

$D_iW = \partial_iW$; $D_i\partial_jW = \partial_i\partial_jW - \Gamma_{ij}^k\partial_kW$

and $g^{i\bar{j}}$ is defined by $g^{i\bar{j}}g_{i\bar{k}} = \delta^{\bar{j}}_{\bar{k}}$.

Applying this to our simple 1-field-model we find the scalar potential

$$V(\phi,\bar{\phi}) = g^{\phi\bar{\phi}}|\partial_\phi W|^2 = g^{\phi\bar{\phi}} = (g_{\phi\bar{\phi}})^{-1} = \frac{1}{1-4\phi\bar{\phi}},$$

which indeed has a minimum at $\phi = 0$.

SUSY breaking is visible since $F^i = -g^{i\bar{j}}\partial_{\bar{j}}\bar{W} \neq 0$.

D-term Breaking

For a U(1) SUSY gauge theory with $V \hat{=} (A_\mu, \lambda, D)$ we have

$$\delta_\xi A_\mu = i\xi\sigma_\mu\bar{\lambda} + h.c.$$
$$\delta_\xi \lambda = i\xi D + \xi\sigma^{\mu\nu}F_{\mu\nu}$$
$$\delta_\xi D = \bar{\xi}\bar{\sigma}^\mu\partial_\mu\lambda + h.c.$$

- ~~SUSY~~ is again signalled by a non-zero auxiliary field: $D = \text{const.} \neq 0$.

- Consider $\mathcal{L} = \frac{1}{2g^2} W^2\big|_{\theta^2} + \underbrace{2\kappa V\big|_{\theta^2\bar{\theta}^2}}$

Fayet-Iliopoulos or "FI" term

(Note: such a term is only allowed in the U_1-case since it would not be gauge-inv. for non-abelian models)

- In Components: $\mathcal{L} = \frac{1}{2g^2} D^2 + \kappa D + \cdots$

$$\Rightarrow D = -\kappa g^2 \neq 0 \Rightarrow \text{~~SUSY~~}$$

- However, the actual dynamical model is still supersymmetric since, because of missing interactions, the spectrum is not affected by $D \neq 0$.

- ~~SUSY~~ becomes apparent in SUSY-QED:

$$\mathcal{L} = \frac{1}{2g^2} W^2\big|_{\theta^2} + \bar{\Phi}_1 e^{2V} \phi\big|_{\theta^2\bar{\theta}^2} + \bar{\Phi}_2 e^{-2V} \phi_2\big|_{\theta^2\bar{\theta}^2}$$

$$+ m\,\phi_1\phi_2\big|_{\theta^2} + \text{h.c.} + 2\kappa V\big|_{\theta^2\bar{\theta}^2}$$

!

$\Rightarrow D \neq 0$; fermionic masses unaffected by $D \neq 0$,

$m_{1,2}^2 = m^2 \pm \kappa g^2$ for charged bosons.

- This feature of a "symmetric mass splitting" induced by SUSY breaking has an important generalization to all renormalizable models with spontaneous ~~SUSY~~:

Mass sum rule: $\mathrm{str}\, M^2 \equiv \mathrm{tr}(M_0^2) - 2\,\mathrm{tr}(M_{1/2}^+ M_{1/2}) + 3\,\mathrm{tr}(M_1^2) \overset{!}{=} 0$

"supertrace" of the mass matrix

scalar mass matrix (for real scalars in the sense of $\mathrm{Re}\,\phi$, $\mathrm{Im}\,\phi$ or ϕ, $\bar\phi$ as indep. degrees of freedom

fermionic mass matrix (for Weyl spinors ψ_i; note that mass terms are holomorphic (or anti-holom.) in ψ

vector boson mass matrix

In words: The sum of all squared particle masses, taking spin multiplicities into account and giving fermionic contributions an opposite sign, vanishes.

- This sum rule is one of the reasons for the fact that:

  It is impossible to break SUSY spontaneously (in a phenomenologically acceptable way) within the supersymmetric SM as a renormalizable theory. (There are very many light fermions and far fewer light bosons.)

Other reasons are the limited availability of F-term VEVs since chiral SFs are charged) and the limited availability of D-terms (just $U_{1,Y}$).

5.3 ~~SUSY~~ in the supersymmetric SM

The simplest solution is to break SUSY spontaneously in a "hidden sector" (a sector coupled to the SM just via non-renormalizable operators). The transmission of $\not{\text{SUSY}}$ to the SM is realized by these operators, which involve the non-zero F-term-VEVs of the hidden sector & SUSY-SM fields.

Simplest realization:

- Let the hidden sector contain a chiral SF S with $\langle S \rangle = \theta^2 M_s^2$ ("S" stands for "spurion"; M_s is the SUSY breaking scale).

[How SUSY breaking is in detail realized in the hidden sector will not be relevant for what follows. One may imagine some O'Raifeartaigh-like model. All masses in the hidden sector are assumed to be large, so that the dynamics of this sector is irrelevant at low energies. The F-term-VEV of S is all that we need to know.]

- Let ϕ be a generic chiral SM SF and V (with corresponding W_α) be a generic SM real SF.

- Assume that the following non-renormalizable couplings exist:

① $\frac{1}{M^2}\bar{\phi}\phi\bar{S}S\Big|_{\theta^2\bar\theta^2}$ or $\frac{1}{M^2}\phi^2\bar{S}S\Big|_{\theta^2\bar\theta^2}$

② $\frac{1}{M}\phi^3 S\Big|_{\theta^2}$ + h.c.

③ $\frac{1}{M}W^2 S\Big|_{\theta^2}$ + h.c.

- Effects: ① $\Rightarrow$ $\frac{M_s^4}{M^2}|A_\phi|^2 \equiv \underset{\uparrow}{M_0^2|A_\phi|^2}$ or $\underset{\nearrow}{M_0^2 A_\phi^2}$

"squark, slepton and soft Higgs masses"

(the holomorphic version appears only in the Higgs sector)

② $\Rightarrow$ $$\frac{M_s^2}{M} A_\phi^3 \equiv \underset{\substack{\uparrow \\ \text{"trilinear couplings"}}}{A} \cdot A_\phi^3$$

③ $\Rightarrow$ $$\frac{M_s^2}{M} \lambda^\alpha\lambda_\alpha \equiv \underset{\substack{\uparrow \\ \text{"gaugino masses"}}}{M_{1/2}} \lambda^\alpha\lambda_\alpha$$

- For $M \sim M_P \sim 10^{19}\,\text{GeV}$ ("gravity mediation") and $M_s \sim 10^{11}\,\text{GeV}$ we find
$$M_{1/2} \sim A \sim M_0 \sim 1\,\text{TeV}\,,$$ as required.

- Clearly, a large set of new parameters are introduced in this way (3 gaugino masses, as many trilinear terms as there are entries in the Yukawa matrices; 3x3 matrices (in generation space) of squark & slepton masses for every matter multiplet, …).

  – FCNCs & other effects rule out most generic structures –

- Need to avoid (forbid by symmetries) renormalizable terms like $\phi^2 S|_{\theta^2}$, since this would give large $\cancel{\text{SUSY}}$ masses.

- Note that the above set of higher-dim. operators (①⋯③) is not complete. (For example, terms like $\bar{\phi}\phi S|_{\theta^2\bar\theta^2}$ may also be present.) However, they are sufficient to generate all so-called "soft $\cancel{\text{SUSY}}$ parameters" or "soft terms" in the low-energy eff. theory.

- The adjective "soft" means that these terms do not destabilize the hierarchy between the electroweak scale (= SUSY breaking scale) and some high scale.

- In the absence of SUSY breaking, the only mass parameter (which appears in the Higgs potential and hence has to be of the order of the electroweak scale) is the μ-term. It is automatically stable (i.e. receives no loop corrections $\sim \Lambda$ (= cutoff)) due to the

  <u>non-renormalization theorem:</u>

  The superpotential receives no loop corrections. (proof see later)

- The terms $\sim W^2|_{\theta^2}$; $\sim \Phi\phi|_{\theta^2\bar\theta^2}$ receive only logarithmically divergent corrections since their coefficients are dimensionless (as in renormalizable non-SUSY QFT).

- The dimensionfull SUSY-breaking terms all come (within the present approach) from higher-dim. operators. Such operators are generically less divergent than renormalizable operators (they can not receive loop corrections $\sim \Lambda$ for dimensional reasons).

- <u>In summary:</u> If we consider a SUSY-SM with SUSY-breaking terms, with $\mu \sim M_s^2/M \sim m_{EW}$, then the electroweak scale (defined in this way) is not destabilized by loop corrections.

- <u>Comment:</u> We argued that in renormalizable SUSY models no divergences $\sim \Lambda^n$ $(n > 0)$ arise. The argument is roughly that the Kähler potential has dimension-ful parameters (only terms $\sim \Phi\phi$ are allowed) while the superpotential is not renormalized for fundamental reasons. There is, however, one exception to this rule:

Terms of the type $\kappa V|_{\theta^2\bar{\theta}^2}$ have the same structure as Kähler terms ($\theta^2\bar{\theta}^2$-projection) and can indeed receive loop corrections $\sim \Lambda^2$. Hower, in the MSSM only one such term (related to $U_{1,Y}$) can be written down. For this term, the Λ^2-divergence has a vanishing coefficient becaus of a cancellation between the various contributions (from $U_{1,Y}$ – charged matter). The deeper reason for this cancellation is the fact that $U_{1,Y}$ is non-anomalous.

- **Comment:** We can now be more specific concerning the idea that $U_{1,R}$ should be broken to $\mathbb{Z}_{2,R}$: let $R(S) = 0$. At the component level this means $R(A_s) = 0$; $R(F_s) = -2$. Thus,

$$U_{1,R}: \quad F_s \to e^{-2i\alpha} F_s \,,$$

and $\langle F_s \rangle \neq 0$ breaks $U_{1,R}$ to $\mathbb{Z}_{2,R}$ (only $U_{1,R}$-trfs. with $\alpha = 0$ & $\alpha = \pi$ leave the vacuum invariant). This consistent with our earlier claim that gaugino masses require the breaking of $U_{1,R}$ to $\mathbb{Z}_{2,R}$: For $R(S) = 0$ both

$W^2|_{\theta^2}$ & $\frac{1}{M} S W^2|_{\theta^2}$ are allowed. The second term induces an operator

$\frac{F_s}{M}\lambda^2$, which indeed breaks $U_{1,R}$ to $\mathbb{Z}_{2,R}$.

- **Comment:** As we will see latter, we will need $\mu \sim \frac{M_s^2}{M} \sim m_{EW}$

phenomenologically. This seems unnatural since, a priori, the two scales μ & M_s^2/M are completely unrelated. However, if $U_{1,R}$ is indeed a symmetry of the underlying theory, $\mu = 0$. If we allow for an operators

$$\frac{1}{M}\bar{S} H_u H_d\Big|_{\theta^2\bar{\theta}^2} \qquad \& \qquad \frac{1}{M^2}\bar{S} S H_u H_d\Big|_{\theta^2\bar{\theta}^2},$$

we find at a (partial) component level

$$\frac{\bar{F}_S}{M}\underbrace{H_u H_d}_{\text{superfields}}\Big|_{\theta^2} \qquad \& \qquad \frac{|F_S|^2}{M^2}\underbrace{H_u H_d}_{\text{A-terms of Higgs-superfields}}$$

$\Rightarrow$ $\mu H_u H_d|_{\theta^2}$ with $\mu \sim M_s^2/M$ ("μ-term")

& $(B\mu) H_u H_d$ with $(B\mu) \sim (M_s^2/M)^2$ ("holomorphic

↑ Higgs mass term or "$B\mu$-term")

(This is just a name - there is no deep reason for writing the coefficient of this operator as a product of μ with some new parameter B.)

This way of generating the supersymmetric μ-term and the SUSY-breaking $B\mu$-term with the same mass scale is known as the "Giudice-Masiero mechanism".

Note: $\bar{S} H_u H_d|_{\theta^2\bar{\theta}^2}$ is $U_{1,R}$-invariant since $R(S) = R(H_{u,d}) = 0$ and the projection on $\theta^2\bar{\theta}^2$ does not change the R-charge.

5.4 Soft terms

SUSY breaking in the MSSM can also be viewed from a different perspective:

- Work out the component lagrangian of the SUSY version of the SM (The μ-term is the only dim.-ful parameter).
- Add all SUSY-breaking terms which do not introduce power-like divergent loop-corrections ("soft terms").
- Without proof, these terms are given in $\mathcal{L}_{soft}$ below.

[$\tilde{g}$ is the superpartner of the gluon g (i.e. a Weyl fermion); $\tilde{e}$ is the superpartner of the r.h. electron (i.e. a scalar); etc.]

$$\begin{aligned}\mathcal{L}_{soft} = &- \frac{1}{2}\left(M_1 \tilde{B}\tilde{B} + M_2 \tilde{A}\tilde{A} + M_3 \tilde{g}\tilde{g}\right) + h.c.\\ &- \left(\tilde{u}^T A_u \tilde{q} H_u + \tilde{d}^T A_d \tilde{q} H_d + \tilde{e}^T A_e \tilde{l} H_d\right) + h.c.\\ &- \left(\tilde{q}^{+} m_Q^2 \tilde{q} + \tilde{l}^{+} m_L^2 \tilde{l} + \tilde{u}^{+} m_u^2 \tilde{u} + \tilde{d}^{+} m_D^2 \tilde{d} + \tilde{e}^{+} m_E^2 \tilde{e}\right) + h.c.\\ &- m_{H_u}^2 \bar{H}_u H_u - m_{H_d}^2 \bar{H}_d H_d - \left(\underset{\uparrow}{b} H_u H_d + h.c.\right)\end{aligned}$$

$b \equiv B\mu$ for simplicity

(M_1, M_2, M_3 are the bino, $\underset{\uparrow}{\text{wino}}$ & gaugino mass)

because of its relation to the W-bosons.

- As we have shown above, all these terms can originate from higher-dim. operators coupling the SUSY-SM to a hidden sector with F-term breaking. This is the conceptual reason why these terms do not introduce power-like divergences. (recall that higher-dim. operators are less divergent than renorm. oper.)

<u>Comment:</u> This argument may appear to be not very convincing since, in the eff. theory below M_S, one simply has given certain scalars a $\cancel{\text{SUSY}}$ mass term, and it is not clear why this mass should not receive corrections $\sim \Lambda$.

- To understand this more deeply, one may perform the following (advanced) exercise:

- Consider

$$\mathcal{L} = \underbrace{(S\bar S - c_1 S\bar S))\Big|_{\theta^2\bar\theta^2} + c_2 S\Big|_{\theta^2} + \text{h.c.}}_{\text{hidden sector}} + \underbrace{\phi\bar\phi\Big|_{\theta^2\bar\theta^2}}_{\text{"Higgs"}} + \underbrace{\frac{1}{M^2}\phi\bar\phi S\bar S\Big|_{\theta^2\bar\theta^2}}_{\cancel{\text{SUSY}}\text{-mediating operator}} .$$

- Work out the general form of loop corrections (SUSY will <u>not</u> be broken by loop corrections $\Rightarrow$ the general superfield structure will be respected; no superpotential for ϕ will be induced; all Kähler terms may receive small loop corrections; new terms like $\phi\bar\phi S\big|_{\theta^2\bar\theta^2}$ etc. may be induced (but they will be suppressed by $\frac{1}{M}$))

- Work out the low-energy spectrum (i.e. the mass of ϕ) before & after these loop corrections.

- The $\cancel{\text{SUSY}}$ mass of ϕ will be $\sim \frac{1}{M}$ in both cases (i.e. the hierarchy will <u>not</u> be destabilized).

<u>Note:</u> For this to work it is important that no linear terms in ϕ (like $\bar S\phi\big|_{\theta^2\bar\theta^2}$) are induced.

This can be achieved by requiring a symmetry like $\phi \to e^{i\alpha}\phi$, which is in fact present in the above lagrangian. In our simplistic model this symmetry is, unfortunately, anomalous. However, in the SM all chiral SFs are protected by non-anomalous gauge symmetries. Thus, no terms linear in the light fields ("tadpoles") are induced. The above argument goes through!

(See the review of Bagger referenced at my SUSY/SUGRA lecture Web page for more details.)