Superconformal OPE methods for gauge-mediated supersymmetry breaking

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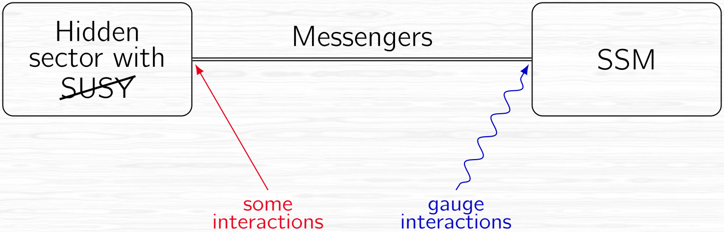
Gauge mediation of ~~SUSY~~

Typical structure of GMSB:



Good features:

- ~~SUSY~~ in SSM generated dynamically.
- Predictive.
- No problems with flavor.

Bad feature:

- μ/B_μ problem.

General Gauge Mediation (GGM)(Meade, Seiberg & Shih, 2008)

Definition

In the limit $g_i \to 0$ we recover the SSM and a separate $\not{SUSY}$ sector.

- Allows us to distinguish model-specific from universal predictions of GMSB.
- All soft masses are given by a small number of two-point correlators of hidden-sector operators.
- Gives a description of the hidden sector dynamics without referring to a Lagrangian.
- Produces mass sum rules (which could be verified experimentally).

Superconformal OPE in GGM(Fortin, Intriligator & AS, 2011)

Idea:

- Assume superconformal symmetry in the UV.
- Use constraints of the symmetry on two- and three-point functions to write down an OPE.
- Compute the Wilson coefficients (UV information).
- Express the soft terms in terms of hidden-sector operators with non-zero vacuum expectation values (IR information, incalculable).

The motivation is similar to the case of QCD, where quark and gluon condensates contribute to the hadronic part of the vacuum polarization of the photon.

The essential difference is that instead of asymptotic freedom we use superconformal symmetry in the UV.

The UV theory may be strongly coupled.

Higgs sector(Kumar, Li, Poland & AS, 2014)

The same idea can be used to study soft parameters in the Higgs sector.

These are also parametrized by two-point correlators of some hidden-sector operators.(Komargodski & Seiberg; 2008)

Using asymptotic superconformal symmetry and the OPE we developed a new language to talk about features of the Higgs sector of the SSM.

We also find a new perspective on the μ/B_μ problem.

Outline

1. Essentials of conformal and $\mathcal{N} = 1$ superconformal symmetry
2. OPE in conformal and superconformal theories
3. Superconformal OPE in GGM
4. Superconformal OPE in the Higgs sector
5. Comments on μ/B_μ

CFTs

In conformal theories the form of the two- and three-point functions of primary operators ($[K_\mu, \mathcal{O}(0)] = 0$) is determined by the symmetry:

$$\langle \mathcal{O}_i^s(x_1)\bar{\mathcal{O}}_{\bar{j}}^s(x_2)\rangle = \delta_{i\bar{j}}\frac{C}{x_{12}^{2\Delta_i}}\mathcal{I}^s(x_{12}), \quad x_{12} \equiv x_1 - x_2,$$

$$\langle \mathcal{O}_i^{s_i}(x_1)\mathcal{O}_j^{s_j}(x_2)\mathcal{O}_k^{s_k}(x_3)\rangle = \frac{c_{ijk}}{x_{12}^{\Delta_i+\Delta_j-\Delta_k}x_{13}^{\Delta_i+\Delta_k-\Delta_j}x_{23}^{\Delta_j+\Delta_k-\Delta_i}}\mathcal{I}^{s_i s_j s_k}(x_{12}, x_{13}, x_{23}).$$

Two and three-point functions of all the descendants can be obtained from the above expressions with the action of $P_\mu = i\partial_\mu$.

OPE in CFTs

The OPE between conformal primaries is

$$\mathcal{O}_i^{s_i}(x_1)\mathcal{O}_j^{s_j}(x_2) = \frac{c_{ij}\mathcal{I}^{s_i s_j}(x_{12})}{x_{12}^{2\Delta_i}}\mathbb{1} + \sum_{k'} \frac{c_{ij}^{k'}}{x_{12}^{\Delta_i+\Delta_j-\Delta_{k'}}}[F_{ij}^{k'}(x_{12},P),\mathcal{O}_{k'}]^{(s_{k'})}(x_2).$$

The two-point function is reproduced by the first term.

The three-point function is reproduced by the second term, something that determines F.(Ferrara, Gatto & Grillo, 1973)

The contributions of all descendants in the OPE are thus determined by the corresponding contributions of the primaries.

$\mathcal{N} = 1$ SCFTs

In $\mathcal{N} = 1$ SCFTs the story is a little different.

Similarly to CFTs, the two-point function of superconformal primaries ($[S^\alpha, \mathcal{O}(0)\} = [\bar{S}^{\dot{\alpha}}, \mathcal{O}(0)\} = 0$) is fixed by the symmetry:

$$\langle \mathcal{O}^s(z_1) \bar{\mathcal{O}}^{\bar{s}}(z_2) \rangle = C_{\mathcal{O}} \frac{\mathcal{I}^{s\bar{s}}(x_{1\bar{2}}, x_{\bar{1}2})}{x_{1\bar{2}}{}^{2q} x_{\bar{1}2}{}^{2\bar{q}}},$$

where

$$x_{\bar{1}2} = -x_{2\bar{1}} = x_{12} - i\theta_1\sigma\bar{\theta}_1 - i\theta_2\sigma\bar{\theta}_2 + 2i\theta_2\sigma\bar{\theta}_1,$$

and

$$\Delta = q + \bar{q}, \qquad R = \tfrac{2}{3}(q - \bar{q}).$$

$\mathcal{N} = 1$ SCFTs

Unlike CFTs, the three-point function of superconformal primaries is not fixed by the symmetry:(Park, 1997; Osborn, 1998)

$$\langle \mathcal{O}_1^{s_1}(z_1)\mathcal{O}_2^{s_2}(z_2)\bar{\mathcal{O}}_3^{\bar{s}_3}(z_3)\rangle = \frac{\mathcal{I}_1^{s_1\bar{s}_1}(x_{1\bar{3}}, x_{\bar{1}3})\mathcal{I}_2^{s_2\bar{s}_2}(x_{2\bar{3}}, x_{\bar{2}3})}{{x_{1\bar{3}}}^{2q_1}{x_{\bar{1}3}}^{2\bar{q}_1}{x_{2\bar{3}}}^{2q_2}{x_{\bar{2}3}}^{2\bar{q}_2}}\bar{t}^{\bar{s}_3}_{\bar{s}_1\bar{s}_2}(X_3, \Theta_3, \bar{\Theta}_3).$$

X, Θ and $\bar{\Theta}$ are convenient "coordinates", with $\Theta \sim \theta$ and thus nilpotent.

The function $\bar{t}$ is arbitrary (up to some conditions). Its appearance here is due to the fact that one can construct superconformal invariants out of three points in superspace.

In conformal field theories one needs four points in space to construct the cross ratios u and v.

OPE in SCFTs

Since the three-point function has a degree of arbitrariness, so does the OPE.

Different classes of operators contribute to the OPE, depending on the independent contributions to the function $\bar{t}(X, \Theta, \bar{\Theta})$.

Perhaps surprisingly, if we take two superconformal primaries and we know the contribution of another primary to their OPE, then we do not necessarily know the contributions of the superdescendants of that primary.

Sometimes we do, though, in particular in some cases where short multiplets are considered.

We will be interested in $\mathcal{N} = 1$ linear, chiral, and antichiral superfields.

Soft masses in GGM

Starting with a linear multiplet,

$$\mathcal{J}(z) = J(x) + i\theta j(x) - i\bar{\theta}\bar{j}(x) - \theta\sigma^{\mu}\bar{\theta}j_{\mu}(x) + \text{derivatives},$$

where the derivatives are determined by

$$D^2\mathcal{J} = \bar{D}^2\mathcal{J} = 0,$$

we get gaugino and sfermion masses by (Buican, Meade, Seiberg & Shih, 2008)

$$M_{\text{gaugino}} = \pi i\alpha \int d^4x\, \langle Q^2(J(x)J(0))\rangle,$$
$$m^2_{\text{sfermion}} = \frac{i\alpha^2 c_2}{8} \int d^4x \ln(x^2 M^2)\langle \bar{Q}^2 Q^2(J(x)J(0))\rangle.$$

sOPE in GGM

Since momentum-space current-current two-point functions $\int d^4x\, e^{-ip\cdot x}\langle J(x)J(0)\rangle$ determine the soft masses, we will use the OPE $J(x) \times J(0)$.

Assumption: Superconformal symmetry in the UV, spontaneously broken in the IR.

In an SCFT we know the form of the two- and three-point functions involving two $\mathcal{J}$'s, and so we can write down the OPE $J(x) \times J(0)$.

Then, we can act with Q^2 and $\bar{Q}^2 Q^2$ to get expressions for gaugino and sfermion masses.

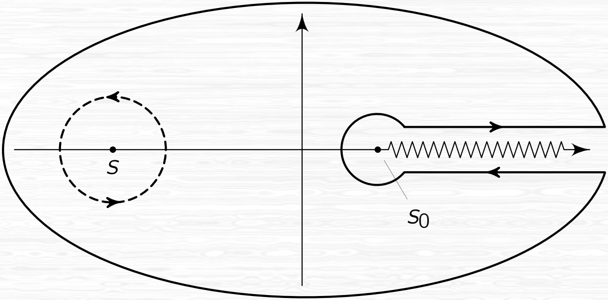
Soft masses from the sOPE

We need

$$i \int d^4x\, e^{-ip\cdot x} J(x)J(0) = \sum_k \tilde{c}^{\,k}_{JJ}(s)\mathcal{O}_k(0), \qquad s = p^2.$$

Compute the OPE coefficients and use a dispersion relation:

$$A(s) = \frac{1}{2\pi i}\int_{s_0}^{\infty} ds' \, \frac{\mathrm{Disc}\, A(s')}{s'-s} = \frac{1}{\pi}\int_{s_0}^{\infty} ds' \, \frac{\mathrm{Im}\, A(s')}{s'-s}.$$



Soft masses from the sOPE

We find

$$M_{\text{gaugino}} \approx \sum_k \frac{\alpha\, \text{Im}[s^{d_k/2}\, \tilde{c}^k_{JJ}(s)]}{2^{d_k-1} d_k M^{d_k}} \langle Q^2(\mathcal{O}_k(0))\rangle,$$

$$m^2_{\text{sfermion}} \approx -\sum_k \frac{\alpha^2 c_2\, \text{Im}[s^{d_k/2}\, \tilde{c}^k_{JJ}(s)]}{2^{d_k+1}\pi d_k^2 M^{d_k}} \langle \bar{Q}^2 Q^2(\mathcal{O}_k(0))\rangle.$$

In minimal GMSB,

$$W = \lambda X \Phi \widetilde{\Phi}, \qquad \langle X \rangle = X + \theta^2 F,$$

$$\mathcal{J} = \Phi^\dagger \Phi - \widetilde{\Phi}\widetilde{\Phi}^\dagger,$$

these methods reproduce the full gaugino and sfermion masses at one loop.

Higgs sector

Our starting point is the simple superpotential

$$W = \lambda_u \mathcal{H}_u \mathcal{O}_u + \lambda_d \mathcal{H}_d \mathcal{O}_d.$$

The couplings $\lambda_{u,d}$ are assumed perturbative.

The operators $\mathcal{O}_{u,d}$ are $SU(2)$ doublets with hypercharge opposite to that of $\mathcal{H}_{u,d}$ respectively, and belong to a sector that breaks SUSY.

$\mathcal{O}_{u,d}$ can be fundamental or composite operators in the hidden sector.

We assume that there is no bare μ-term in the superpotential.

Higgs Lagrangian

$$\begin{aligned}
\mathscr{L} = \mathcal{Z}_u F^\dagger_{H_u} F_{H_u} + \mathcal{Z}_d F^\dagger_{H_d} F_{H_d} + \left(\mu \int d^2\theta \, \mathcal{H}_u \mathcal{H}_d + \text{c.c.} \right) \\
- V^{(\text{soft})}_{\text{Higgs}} - V^{(\text{other})}_{\text{Higgs}},
\end{aligned}$$

$$\begin{aligned}
V^{(\text{soft})}_{\text{Higgs}} &= m^2_{H_u} H^\dagger_u H_u + m^2_{H_d} H^\dagger_d H_d + (B_\mu H_u H_d + \text{c.c.}) \\
&\quad + (A_u H_u F^\dagger_{\mathcal{H}_u} + A_d H_d F^\dagger_{\mathcal{H}_d} + \text{c.c.}), \\
V^{(\text{other})}_{\text{Higgs}} &= (a'_u H_u F_{\mathcal{H}_d} + \text{c.c.}) + (a'_d F_{\mathcal{H}_u} H_d + \text{c.c.}) \\
&\quad + (\gamma F_{\mathcal{H}_u} F_{\mathcal{H}_d} + \text{c.c.}).
\end{aligned}$$

Soft terms

We can integrate out the dynamics of the hidden sector and generate the quadratic terms in the Higgs Lagrangian, with

$$\mu = \frac{i}{8}\lambda_u \lambda_d \int d^4x \, \langle Q^\alpha(O_u(x)) Q_\alpha(O_d(0))\rangle,$$
$$B_\mu = \frac{i}{2^5}\lambda_u \lambda_d \int d^4x \, \langle Q^2(O_u(x)) Q^2(O_d(0))\rangle,$$
$$A_{u,d} = -\frac{i}{8}|\lambda_{u,d}|^2 \int d^4x \, \langle Q^2(O_{u,d}(x)\bar{O}_{u,d}(0))\rangle,$$
$$m^2_{H_{u,d}} = -\frac{i}{2^5}|\lambda_{u,d}|^2 \int d^4x \, \langle Q^2\bar{Q}^2(O_{u,d}(x)\bar{O}_{u,d}(0))\rangle,$$

and similar formulas for the remaining parameters.

Chiral-chiral three-point function (Vichi, 2011)

Analysis of the three-point function

$$\langle \Phi_1(z_{1+})\Phi_2(z_{2+})\bar{\mathcal{O}}^{\bar{I}}(z_3)\rangle = \frac{\lambda_{12\mathcal{O}}}{x_{\bar{3}1}{}^{2\Delta_1} x_{\bar{3}2}{}^{2\Delta_2}} \bar{t}^{\bar{I}}(\bar{X}_3, \Theta_3, \bar{\Theta}_3),$$

reveals three possible contributions to $\bar{t}$:

1. $\bar{t}_1(\bar{X}_3, \bar{\Theta}_3) = 1, \qquad (\Delta_{\mathcal{O}} = \Delta_1 + \Delta_2,\ R_{\mathcal{O}} = R_1 + R_2)$
2. $\bar{t}_{2\,\alpha_2\ldots\alpha_\ell}^{\dot{\alpha}_1\ldots\dot{\alpha}_\ell}(\bar{X}_3, \bar{\Theta}_3) = \bar{\Theta}_3^{(\dot{\alpha}_1}\bar{\mathrm{X}}_{3\alpha_2}^{\dot{\alpha}_2}\cdots\bar{\mathrm{X}}_{3\alpha_\ell}^{\dot{\alpha}_\ell)},$
   $(\Delta_{\mathcal{O}} = \Delta_1 + \Delta_2 + \ell - \tfrac{1}{2},\ R_{\mathcal{O}} = R_1 + R_2 - 1),$
3. $\bar{t}_3^{\mu_1\ldots\mu_\ell}(\bar{X}_3, \bar{\Theta}_3) = \bar{\Theta}_3^2 \dfrac{\bar{X}_3^{\mu_1}\cdots\bar{X}_3^{\mu_\ell}}{\bar{X}_3^{\Delta_1+\Delta_2-\Delta_{\mathcal{O}}+\ell+1}},$
   $(\Delta_{\mathcal{O}} \geq |\Delta_1 + \Delta_2 - 3| + \ell + 2,\ R_{\mathcal{O}} = R_1 + R_2 - 2).$

Chiral-chiral sOPE

Given the three classes of contributions to the three-point function, there are three classes of operators that contribute to the chiral-chiral OPE.

We are only interested in conformal primary operators that can get a vev consistently with Poincaré invariance and SM gauge invariance.

Such operators are not necessarily superconformal primaries.

In order to contain a scalar component an $\mathcal{N} = 1$ superfield must have $j + \bar{\jmath} = 0, \frac{1}{2}, 1$.

So we will consider the superconformal primary operators

$$\mathcal{O}_{1,0}, \quad \mathcal{O}^{\alpha}_{2,1}, \quad \mathcal{O}_{3,0}, \quad \mathcal{O}^{\mu}_{3,1}.$$

Chiral-antichiral three-point function (Poland & Simmons-Duffin, 2010)

Here we have a simple three-point function:

$$\langle \Phi(z_{1+})\bar{\Phi}(z_{2-})\bar{\mathcal{O}}^{\mu_1\ldots\mu_\ell}(z_3)\rangle = \frac{\lambda_{1\bar{2}\mathcal{O}^{\mu_1\ldots\mu_\ell}}}{x_{\bar{2}3}{}^{2\Delta_\Phi}x_{\bar{3}1}{}^{2\Delta_\Phi}}\bar{X}_3^{\Delta_{\mathcal{O}}-2\Delta_\Phi-\ell}\bar{X}_3^{\mu_1}\ldots\bar{X}_3^{\mu_\ell}.$$

All three-point functions of, say, the zero components of the first two operators with various components of $\bar{\mathcal{O}}$ are determined by the coefficient of the three-point function $\langle\phi(x_1)\bar{\phi}(x_2)\bar{O}^{\mu_1\ldots\mu_\ell}(x_3)\rangle$.

Here the unitarity bound is

$$\Delta_{\mathcal{O}} \geq \ell + 2.$$

Chiral-antichiral sOPE

The chiral-antichiral sOPE in this case is also particularly simple.

The only operators that can contribute have to have zero R-charge.

Since the parameters we are interested in ($A_{u,d}$ and $m^2_{H_{u,d}}$) are given by Q^2 and $Q^2\bar{Q}^2$ on the OPE, $O_{u,d}(x) \times \bar{O}_{u,d}(0)$, we only need to consider superconformal primary scalar operators

$$\mathcal{O}_0.$$

Only the zero component of $\mathcal{O}_0$ is of interest.

μ

We need $Q^\alpha O_u(x) \times Q_\alpha O_d(0)$.

We start from the known expression for the three-point function $\langle \mathcal{O}_u(z_{1+}) \mathcal{O}_d(z_{2+}) \bar{\mathcal{O}}^{\bar{I}}(z_3) \rangle$, and we pick the $\theta_1\theta_2$ component.

We then further expand in $\theta_3, \bar{\theta}_3$ and we can finally find

$$
\begin{aligned}
Q^\alpha O_u \times Q_\alpha O_d = {} & c_{\mu;1} Q^2 O_{1,0} + c_{\mu;2} Q^\alpha O_{2,1\,\alpha} \\
& + c_{\mu;3} O_{3,0} + c_{\mu;4} [Q^2 \bar{Q}^2 O_{3,0}]_p \\
& + c_{\mu;5} [Q \sigma_\mu \bar{Q} O^\mu_{3,1}]_p + \cdots ,
\end{aligned}
$$

Determining the coefficients c_μ is rather complicated, especially for the primaries in higher θ-components of $\mathcal{O}_{3,0}$ and $\mathcal{O}^\mu_{3,1}$.

The complication arises from the presence of conformal descendants.

Aside: Conformal primary norms

In order to complete this computation we need the projection of the $\mathcal{N}=1$ superconformal two-point function to the conformal subgroup.

We worked this out for scalar and spin-one operators.

For example,

$$\langle [Q^2\bar{Q}^2 O]_p(x)[Q^2\bar{Q}^2 O]_p^\dagger(0)\rangle = \\ 2^{12} C_{\mathcal{O}} \frac{q\bar{q}(q-1)(\bar{q}-1)(q+\bar{q})(q+\bar{q}+1)}{(q+\bar{q}-1)(q+\bar{q}-2)} \frac{1}{x^{2(q+\bar{q}+2)}},$$

where

$$[Q^2\bar{Q}^2 O]_p = Q^2\bar{Q}^2 O - 2^4 \frac{\bar{q}(\bar{q}-1)}{(q+\bar{q}-1)(q+\bar{q}-2)} P^2 O \\ - 8\frac{\bar{q}-1}{q+\bar{q}-2} P_\mu Q\sigma^\mu \bar{Q} O.$$

μ

$$
\begin{aligned}
Q^\alpha O_u \times Q_\alpha O_d = {} & c_{\mu;1} Q^2 O_{1,0} + c_{\mu;2} Q^\alpha O_{2,1\,\alpha} \\
& + c_{\mu;3} O_{3,0} + c_{\mu;4} [Q^2 \bar{Q}^2 O_{3,0}]_p \\
& + c_{\mu;5} [Q \sigma_\mu \bar{Q} O^\mu_{3,1}]_p + \cdots ,
\end{aligned}
$$

To find μ we need to Fourier-transform this and go to the zero momentum limit.

The coefficients $c_{\mu;1}$ and $c_{\mu;2}$ have no x-dependence.

Assuming the simple analytic structure we saw before, it turns out that $Q^2 O_{1,0}$ and $Q^\alpha O_{2,1\,\alpha}$ do not actually contribute to the μ-term.

One way to see this:

$$
i \int d^4x \, e^{-ip\cdot x} \frac{1}{(x^2)^\epsilon} = \pi^2 \frac{\Gamma(2-\epsilon)}{2^{2\epsilon-4}\Gamma(\epsilon)} \frac{1}{(p^2)^{2-\epsilon}} .
$$

μ

We can finally write

$$\mu = \lambda_u \lambda_d \left(\hat{c}_{\mu;3} \langle O_{3,0} \rangle + \hat{c}_{\mu;4} \langle Q^2 \bar{Q}^2 O_{3,0} \rangle + \hat{c}_{\mu;5} \langle Q \sigma_\mu \bar{Q} O^\mu_{3,1} \rangle \right),$$

with

$$\hat{c}_{\mu;3} = \frac{i}{8} \check{c}_{\mu;3} \int d^4x \, e^{-ip\cdot x} x^{\Delta_{\mathcal{O}_{3,0}} - \Delta_{\mathcal{O}_u} - \Delta_{\mathcal{O}_d} - 1} \Big|_{p\to 0},$$

$$\hat{c}_{\mu;4} = \frac{i}{8} \check{c}_{\mu;4} \int d^4x \, e^{-ip\cdot x} x^{\Delta_{\mathcal{O}_{3,0}} - \Delta_{\mathcal{O}_u} - \Delta_{\mathcal{O}_d} + 1} \Big|_{p\to 0},$$

$$\hat{c}_{\mu;5} = \frac{i}{8} \check{c}_{\mu;5} \int d^4x \, e^{-ip\cdot x} x^{\Delta_{\mathcal{O}^\mu_{3,1}} - \Delta_{\mathcal{O}_u} - \Delta_{\mathcal{O}_d}} \Big|_{p\to 0}.$$

B_μ

For B_μ we need the OPE $Q^2 O_u(x) \times Q^2 O_d(0)$, i.e. we need to look at the $\theta_1^2 \theta_2^2$ order of the three-point function $\langle \mathcal{O}_u(z_{1+}) \mathcal{O}_d(z_{2+}) \bar{\mathcal{O}}^{\bar{I}}(z_3) \rangle$.

The answer is

$$B_\mu = \lambda_u \lambda_d \hat{c}_{B_\mu} \langle Q^2 O_{3,0} \rangle,$$

where

$$\hat{c}_{B_\mu} = -\tfrac{1}{4} \hat{c}_{\mu;3}.$$

This relation follows from

$$Q^2(Q^\alpha(O_u(x)) Q_\alpha(O_d(0))) = -Q^2(O_u(x)) Q^2(O_d(0)).$$

Go to soft terms

$A_{u,d}$ and $m^2_{H_{u,d}}$

From

$$\langle \mathcal{O}_{u,d}(z_{1+})\bar{\mathcal{O}}_{u,d}(z_{2-})\bar{\mathcal{O}}_0(z_3)\rangle = \lambda_{\mathcal{O}_{u,d}\bar{\mathcal{O}}_{u,d}\mathcal{O}_0} \frac{x_{\bar{2}1}{}^{\Delta_{\mathcal{O}_0}-2\Delta_{\mathcal{O}_{u,d}}}}{x_{\bar{2}3}{}^{\Delta_{\mathcal{O}_0}} x_{\bar{3}1}{}^{\Delta_{\mathcal{O}_0}}},$$

we can extract the OPE

$$O_{u,d} \times \bar{O}_{u,d} = c^i_{u,d} O^{u,d}_{0;i} + \cdots \,,$$

and obtain

$$A_{u,d} = |\lambda_{u,d}|^2 \hat{c}_{A_{u,d}} \langle Q^2 O^{u,d}_0 \rangle,$$
$$m^2_{H_{u,d}} = |\lambda_{u,d}|^2 \hat{c}_{m^2_{H_{u,d}}} \langle Q^2 \bar{Q}^2 O^{u,d}_0 \rangle,$$

with

$$\hat{c}_{m^2_{H_{u,d}}} = \frac{1}{4} \hat{c}_{A_{u,d}}.$$

Summary of results

We have found expressions that allow us to determine the Higgs parameters by identifying hidden-sector operators that can get a vev.

$$\begin{aligned} \mu &: \quad O_{3,0}, \;\; Q^2\bar{Q}^2 O_{3,0}, \;\; Q\sigma_\mu \bar{Q} O^\mu_{3,1}, \\ B_\mu &: \quad Q^2 O_{3,0}, \\ A_{u,d} &: \quad Q^2 O_0^{u,d}, \\ m^2_{H_{u,d}} &: \quad Q^2\bar{Q}^2 O_0^{u,d}. \end{aligned}$$

Example: in a spurion model $O_{3,0} = X^\dagger F^\dagger$, $O_0 = X^\dagger X$.

The soft parameters are determined at the messenger scale M.

To study their phenomenology we have to RG evolve them down to the weak scale.

μ/B_μ problem

Electroweak symmetry breaking in the MSSM requires

$$B_\mu \approx \mu^2.$$

In gravity mediation this is achived by the Giudice–Masiero mechanism:

$$\frac{1}{M_P}\int d^4\theta\, X^\dagger \mathcal{H}_u \mathcal{H}_d, \qquad \frac{1}{M_P^2}\int d^4\theta\, X^\dagger X \mathcal{H}_u \mathcal{H}_d.$$

In gauge mediation, however,

$$B_\mu \gg \mu^2.$$

In the weakly-coupled example

$$W = \lambda_u \mathcal{H}_u \Phi_1 \Phi_2 + \lambda_d \mathcal{H}_d \tilde{\Phi}_1 \tilde{\Phi}_2 + \lambda X(\Phi_1 \tilde{\Phi}_1 + \Phi_2 \tilde{\Phi}_2),$$

we can compute

$$B_\mu \approx 16\pi^2 \mu^2.$$

μ/B_μ problem

From our OPE analysis the problem in a weakly-coupled spurion model is also clear:

$$\mu \sim \frac{c_\mu}{M^2}\langle O_{3,0}\rangle \sim c_\mu \frac{F}{M},$$

$$B_\mu \sim \frac{c_\mu}{M^2}\langle Q^2 O_{3,0}\rangle \sim c_\mu \frac{F^2}{M^2},$$

and so

$$\frac{B_\mu}{\mu^2} \sim c_\mu^{-1} \sim 16\pi^2 \gg 1.$$

One way out is to find models where $c_\mu \sim 1$, but that requires strong coupling.

Another way out is to assume some dynamics between M and $\sqrt{F}$ that suppress B_μ and give $B_\mu \approx \mu^2$ at the weak scale. (Conformal sequestering(Murayama, Nomura & Poland, 2007; Roy & Schmaltz, 2007))

μ/B_μ problem

An interesting direction to explore is if there are models where $O_{3,0}$ does not appear.

In that case μ is still generated by $Q\sigma_\mu \bar{Q} O^\mu_{3,1}$:

$$\mu \sim \frac{F^2}{M^3}.$$

$A_{u,d}$ and $m^2_{H_{u,d}}$ are generated at the messenger scale M.

B_μ is zero at the messenger scale, but it is generated by RG runnning down to the weak scale:

$$B_\mu \approx \mu(M_{\text{gaugino}} + A).$$

Then, at the weak scale,

$$\frac{B_\mu}{\mu^2} \sim \frac{M_{\text{gaugino}} + A}{\mu} \sim \left(\frac{F}{M^2}\right)^{-1} > 1.$$

Summary

We used an OPE analysis to classify hidden-sector operators that can generate soft parameters in the MSSM.

This approach offer a new perspective on the quadratic Higgs Lagrangian, and elegantly explains known features of models found in the literature.

We relied on superconformal symmetry in the UV, and also on various assumptions on the analytic structure of the relevant two-point functions.

Future work:

- Study effects of poles in the relevant correlation functions.
- Explore further the μ/B_μ problem.
- Study the quartic terms in the Higgs Lagrangian using similar methods.

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- Explore further the μ/B_μ problem.
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Thank you!

BACKUP

$a'_{u,d}$

$$a'_{u,d} = \frac{i}{8}\lambda_u\lambda_d \int d^4x \, \langle Q^\alpha(O_{u,d}(x) Q_\alpha(O_{d,u}(0)))\rangle .$$

We have the OPE

$$\begin{aligned} Q^\alpha(O_{u,d} \times Q_\alpha O_{d,u}) = {} & c_{a'_{u,d;1}} Q^2 O_{1,0} + c_{a'_{u,d;2}} Q^\alpha O_{2,1\,\alpha} \\ & + c_{a'_{u,d;3}} [Q^2 \bar{Q}^2 O_{3,0}]_p \\ & + c_{a'_{u,d;4}} [Q \sigma_\mu \bar{Q} O^\mu_{3,1}]_p + \cdots . \end{aligned}$$

We have exactly the same contributions as for μ, except, of course, for the $O_{3,0}$ contribution.

We can obtain

$$a'_{u,d} = \lambda_u\lambda_d(\hat{c}_{a'_{u,d;3}} \langle Q^2 \bar{Q}^2 O_{3,0}\rangle + \hat{c}_{a'_{u,d;4}} \langle Q\sigma_\mu \bar{Q} O^\mu_{3,1}\rangle).$$

Smallness of $a'_{u,d}$

In the presence of singlets $a'_{u,d}$ is not soft.

Typically a' is neglected because it comes out suppressed in most cases.

This is easy to understand with our methods:

$$\mu \sim \frac{c_\mu}{M^2}\langle O_{3,0}\rangle \sim c_\mu \frac{F}{M},$$
$$a'_{u,d} \sim \frac{c_{a'_{u,d}}}{M^4}\langle Q^2\bar{Q}^2 O_{3,0}\rangle \sim c_{a'_{u,d}}\frac{F^2}{M^3}.$$

If, however, we assume that the operator that generates μ and $a'_{u,d}$ is not $O_{3,0}$ but rather $Q\sigma_\mu \bar{Q} O^\mu_{3,1}$, then μ and $a'_{u,d}$ are of the same order and both suppressed,

$$\mu,\, a'_{u,d} \sim \frac{F^2}{M^3}.$$