Dirac gauginos in low scale supersymmetry

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Overview

Part 1:

- Motivation for studying low energy supersymmetry
- Why it requires Dirac gauginos
- Some background about Dirac gauginos
- Consequences for goldstino couplings
- Collider phenomenology

Part 2:

- Two loop Higgs mass corrections in Dirac gaugino models
- ... and beyond

Low energy SUSY breaking: what is it and why?

- We know that SUSY must be broken in a hidden sector thanks to the supertrace rule.
- Typically we imagine that this happens by dynamics at a high mass scale (e.g. gauge mediation) and possibly (very) weakly coupled to the visible sector (e.g. gravity mediation).

However, we do not know that this must be true – if we relax this assumption by lowering the scale then there are interesting consequences!

- Naturalness: alleviate the little hierarchy problem: less running from the scale Λ to the electroweak scale
- We can decouple any higher-energy physics – can study everything in field theory.
- Couplings to gravitino actually become large.
- Higher-dimensional operators contribute to Higgs potential and can raise the Higgs mass at “tree” level.

The goldstino

- In global SUSY have a massless fermion G_α corresponding to the spontaneously broken SUSY generators Q_α.
- This then corresponds to the mode eaten by the gravitino; in SUGRA

$$\mathcal{L} \supset m_{3/2}\psi_\mu \sigma^{\mu\nu}\psi_\nu - \frac{F^i}{M_P}\psi_\mu \sigma^\mu \chi_i$$

- Diagonalise these and associate F^X to the goldstino field.
- Can write as part of a supermultiplet

$$\mathbf{X} = X + \sqrt{2}\theta G + \theta^2 F_X + \ldots$$

- The sgoldstino X may be dynamical, or we may integrate it out – and find we can write in terms of a non-linear superfield

$$\mathbf{X} = \frac{\overline{G}G}{2F_X} + \sqrt{2}\theta G + \theta^2 F_X + \ldots$$

- We can use this formalism to derive all of the low-energy couplings of the goldstino by writing all the SUSY-breaking terms in terms of F_X and extending this to the superfield $\mathbf{X}$.

A light gravitino and strong couplings

- Assuming minimal supergravity couplings, $m_{3/2} = \frac{f}{\sqrt{3}M_P}$, so there is a lower bound on the mass of the gravitino.
- Couplings of the goldstino to matter look like

$$\mathcal{L} \supset -\int d^4\theta \frac{m_X^2}{f^2}\overline{X}X\overline{\Phi}\Phi + \int d^2\theta \frac{M_\lambda}{2f} X W^\alpha W_\alpha$$
$$\rightarrow \frac{m_X^2}{f}\overline{\Phi}\psi G + \frac{M}{\sqrt{2}f}\lambda^a \sigma^{\mu\nu} G F^a_{\mu\nu} + h.c. + \ldots$$

- i.e. the smaller f, the larger the coupling!
- Note that M_P does not appear so these couplings always dominate the gravitational ones.

Realising low energy SUSY

It is then important to ask:

- How could a low-scale SUSY-breaking model arise?
- What would be the relationship between f and the soft masses? These would be most interesting for $m_{\mathrm{soft}}/\sqrt{f} \sim \mathcal{O}(1)$

These can only be answered through examples:

- Gravity mediation is clearly not appropriate:

$$m_{soft} \lesssim \frac{f}{M_P} \sim m_{3/2}$$

- We could try gauge mediation:

$$m_{soft} \sim \frac{\alpha}{4\pi} \frac{f}{M_{\mathrm{Mess}}}, \qquad M^2_{\mathrm{Mess}} > f$$

- $\rightarrow$Due to the loop factor, if we take $M^2_{\mathrm{Mess}} \sim f$ then we need $\sqrt{f} \gtrsim$ 100 TeV.
- More "exotic" models are therefore called for: tree-level gauge mediation/strong dynamics.

One scenario

- [Gherghetta, Pomerol '11] considered breaking SUSY by strong dynamics, so we have

$$V \sim \frac{N\Lambda^4}{16\pi^2} \to f, D \sim \frac{\sqrt{N}\Lambda^2}{4\pi}$$

- Scalar matter fields can then couple to the strong sector linearly and obtain masses at one loop:

$$\mathcal{L} \supset g_i \Phi_i \mathcal{O}_{\text{strong}} \to m_{\text{scalar},i} \sim g_i \frac{\Lambda}{4\pi}$$

- For Λ ~TeV this is small.
- By invoking (partial) compositeness of the matter fields they can interact directly with the strong sector – we would then find

$$m_{\text{scalar},i} \sim g_i \frac{\sqrt{N}}{4\pi}\Lambda \sim \epsilon_i \Lambda, \qquad \epsilon_i \lesssim 0.3$$

- However, the Majorana gaugino masses would still have to be generated at one loop – and would thus potentially be too small.
- This problem is familiar also to "tree level gauge mediation" where scalars obtain masses through their D-term interactions at tree level, but the gauginos must still obtain masses at one loop.

Low energy SUSY and Dirac gauginos

- Instead [Gherghetta, Pomerol '11] proposed that it is necessary to give Dirac masses to the gauginos to realise this scenario.
- In this way, there are new adjoint chiral superfields Σ that couple to the strong dynamics via a superpotential term

$$W \supset \lambda_\Sigma \Sigma J_2$$

- J_2 is a current made from the fields in the strong sector and $\lambda_X \sim 4\pi/\sqrt{N}$.
- The gaugino masses generated are essentially generated by gauge mediation, and given by

$$m_{\lambda_a} = c_{m_{\lambda_a}} g_a \frac{\sqrt{N}}{4\pi}\Lambda \simeq 2\ \text{TeV} c_{m_{\lambda_a}} g_a \sqrt{\frac{N}{6}} \left(\frac{\Lambda}{10\ \text{TeV}}\right)$$

- Note in this scenario the scalars in the multiplet Σ have masses $\sim \Lambda$.

More motivation for Dirac gauginos

The study of Dirac gaugino masses is of course an ongoing large research project with many additional motivations:

- Dirac gauginos allow to relax LHC search bounds as production of squarks is suppressed since no chirality flip is possible. Gluino production is enhanced a little relative to MSSM, but this is greatly suppressed when $m_{\tilde{q}_{1,2}} \gg m_{\tilde{g}}$.
- They typically suppress processes such as $B \to s\gamma$ and $\Delta F = 2$.
- They allow for increased naturalness: supersoft masses do not lead to large corrections to stop mass.
- They allow new Higgs couplings, permitting increased Higgs mass $\to$ compatibility with e.g. light stops.
- There would have been/could still be clear signals from accompanying adjoint scalars if light (this would have been a surprise)
- If gauginos are found at the LHC, we will have to determine whether they are Majorana or Dirac in nature, and this is very difficult to do directly: maybe only possible at ILC

Dirac gauginos: top down

Some attractive theoretical motivations:

- Nelson-Seiberg Theorem: existence of R symmetry (chiral symmetry under which bosons are also charged: $\Phi \to e^{i\alpha R_\Phi}\Phi$, $\theta \to e^{i\alpha}\theta$, $W \to e^{2i\alpha}W$) required for F-term SUSY breaking
- Many ~~SUSY~~ models preserve R symmetry (e.g. original O'Raifertaigh model)
- Dirac gaugino mass may preserve R , Majorana does not: [Fayet, 78] suggested this as the original way to obtain gaugino masses!
- Alternatively Majorana gaugino mass may be too small (e.g. from many O'Raifeartaigh models [Komargodsky and Shih, 2008], [Abel, Jaeckel, Khoze 09])

Chiral Adjoint Fields

- MSSM chiral superfields are in singlet, fundamental and antifundametal reps; vector superfields are in adjoint reps.
- To allow Dirac masses for the gauginos, must add chiral ▸ adjoint field:

$$\mathbf{\Sigma} = \Sigma + \sqrt{2}\theta^{\alpha}(\chi)_{\alpha} + (\theta\theta)F_{\Sigma} + \ldots \rightarrow \mathcal{L} \supset -m_D \chi\lambda$$

- $\rightarrow$ Adjoint superfields will contain fermions to partner gauginos, but scalars too.
- $N \geqslant 2$ supersymmetry - chiral adjoint is superpartner of vectors
- Higher dimensional models:

$$\partial^5_{[\mu} A^5_{\nu]} - [A^5_{\mu}, A^5_{\nu}] = \begin{cases} \partial^4_{[\mu} A^4_{\nu]} - [A^4_{\mu}, A^4_{\nu}] \\ \partial^4_{\mu} A_5 - [A^4_{\mu}, A_5] \end{cases}$$

- $\rightarrow$Current in warped models
- Seiberg dualities (e.g. ISS):

$$Q_i \tilde{Q}_j = \mu X_{ij} = \mu \delta_{ij} \mathrm{tr} X_{ii} + \mu (X_{ij} - \delta_{ij} \mathrm{tr} X_{ii})$$

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MSSM with Adjoints

Names		Spin 0	Spin 1/2	Spin 1	$SU(3), SU(2), U(1)_Y$
Quarks (×3 families)	$\mathbf{Q}$ $\mathbf{u^c}$ $\mathbf{d^c}$	$\tilde{Q} = (\tilde{u}_L, \tilde{d}_L)$ $\tilde{u}_L^c$ $\tilde{d}_L^c$	(u_L, d_L) u_L^c u_L^c		(**3**, **2**, 1/6) ($\bar{\mathbf{3}}$, **1**, -2/3) ($\bar{\mathbf{3}}$, **1**, 1/3)
Leptons (×3 families)	$\mathbf{L}$ $\mathbf{e^c}$	$(\tilde{\nu}_{eL}, \tilde{e}_L)$ $\tilde{e}_L^c$	(ν_{eL}, e_L) e_L^c		(**1**, **2**, -1/2) (**1**, **1**, 1)
Higgs	$\mathbf{H_u}$ $\mathbf{H_d}$	(H_u^+, H_u^0) (H_d^0, H_d^-)	$(\tilde{H}_u^+, \tilde{H}_u^0)$ $(\tilde{H}_d^0, \tilde{H}_d^-)$		(**1**, **2**, 1/2) (**1**, **2**, -1/2)
Gluons	$\mathbf{W}_{3\alpha}$		$\lambda_{3\alpha}$ $[\equiv \tilde{g}_\alpha]$	g	(**8**, **1**, 0)
W	$\mathbf{W}_{2\alpha}$		$\lambda_{2\alpha}$ $[\equiv \tilde{W}^\pm, \tilde{W}^0]$	$W^\pm, W^0$	(**1**, **3**, 0)
B	$\mathbf{W}_{1\alpha}$		$\lambda_{1\alpha}$ $[\equiv \tilde{B}]$	B	(**1**, **1**, 0)
DG-octet	$\mathbf{O_g}$	O_g $[\equiv \Sigma_g]$	χ_g $[\equiv \tilde{g}']$		(**8**, **1**, 0)
DG-triplet	$\mathbf{T}$	$\{T^0, T^\pm\}$ $[\equiv \{\Sigma_0^W, \Sigma_W^\pm\}]$	$\{\chi_T^0, \chi_T^\pm\}$ $[\equiv \{\tilde{W}'^\pm, \tilde{W}'^0\}]$		(**1**,**3**, 0)
DG-singlet	$\mathbf{S}$	S $[\equiv \Sigma_B]$	χ_S $[\equiv \tilde{B}']$		(**1**, **1**, 0)

Supersymmetric Couplings

Here are the most general renormalisable superpotential couplings:

- SUSY couplings contained in superpotential:
  $W = W_{\mathrm{Yukawa}} + W_{\mathrm{Higgs}} + W_{\mathrm{Adjoint}}$
- No new Yukawas:

$$W_{\mathrm{Yukawa}} = Y_U^{ij} \mathbf{Q_i} \cdot \mathbf{H_u u_j^c} + Y_D^{ij} \mathbf{Q_i} \cdot \mathbf{H_d d_j^c} + Y_E^{ij} \mathbf{L_i} \cdot \mathbf{H_d e_j^c}$$

- Two new Higgs couplings (c.f. NMSSM):

$$W_{\mathrm{Higgs}} = \mu \mathbf{H_u} \cdot \mathbf{H_d} + \lambda_S \mathbf{S H_d} \cdot \mathbf{H_u} + 2\lambda_T \mathbf{H_d} \cdot \mathbf{T H_u}$$

- Several possible new Adjoint couplings which violate R:

$$\begin{aligned} W_{\mathrm{Adjoint}} =& L\mathbf{S} + \frac{M_S}{2}\mathbf{S}^2 + \frac{\kappa_S}{3}\mathbf{S}^3 + M_T \mathrm{tr}(\mathbf{TT}) + \lambda_{ST}\mathbf{S}\mathrm{tr}(\mathbf{TT}) \\ &+ M_O \mathrm{tr}(\mathbf{OO}) + \lambda_{SO}\mathbf{S}\mathrm{tr}(\mathbf{OO}) + \frac{\kappa_O}{3}\mathrm{tr}(\mathbf{OOO}). \end{aligned}$$

D-term masses

- Write the Dirac mass as a Holomorphic term; gives new D-term interactions:

$$\int d^2\theta 2\sqrt{2} m_D \theta^\alpha \mathrm{tr}(W_\alpha \Sigma) \supset -m_D(\lambda_a \chi_a) + \sqrt{2} m_D \Sigma_a D_a$$

The New D-term couplings have two main effects:

- Adjoint scalar masses and B-type masses are modified:

$$\mathcal{L} \supset \frac{1}{2} D_a^2 + \sqrt{2}(m_D \Sigma_a + \bar{m}_D \overline{\Sigma}_a) D_a$$
$$\xrightarrow{m_D \text{ real}} -\frac{1}{2} m_D^2 (\Sigma_a + \overline{\Sigma}_a)^2$$

- Trilinear terms modify Higgs mass matrix

$$\frac{1}{\sqrt{2}} g_Y m_D (S + \overline{S})(H_u^* H_u - H_d^* H_d) \supset -g_Y m_D c_{2\beta} v \,(s_R h)$$

Naturalness

Scalar masses thus enter into naturalness bounds:

$$\delta m^2_{H_{u,d}} \supset -\frac{1}{16\pi^2}(2\lambda_S^2 m_S^2 + 6\lambda_T^2 m_T^2)\log\left(\frac{\Lambda}{\mathrm{TeV}}\right)$$

where Λ is the UV cutoff of the theory. Using $\Delta \equiv \frac{\delta m_h^2}{m_h^2}$ then we have

$$m_S \lesssim \mathrm{TeV}\left(\frac{1}{\lambda_S}\right)\left(\frac{\log \Lambda/\mathrm{TeV}}{3}\right)^{-1/2}\left(\frac{\Delta^{-1}}{20\%}\right)^{-1/2}$$

$$m_T \lesssim 5\ \mathrm{TeV}\left(\frac{0.1}{\lambda_T}\right)\left(\frac{\log \Lambda/\mathrm{TeV}}{3}\right)^{-1/2}\left(\frac{\Delta^{-1}}{20\%}\right)^{-1/2}$$

(in the absence of a large tree-level contribution). This is good, because

$$\Delta\rho \simeq \frac{v^2(g_2 m_{D2} c_{2\beta} + \sqrt{2}\tilde{\mu}\lambda_T)^2}{(m_T^2 + |M_T|^2 + B_T + 4|m_{D2}|^2)^2} \lesssim 8\times 10^{-4} \rightarrow m_T \gtrsim 1.4\ \mathrm{TeV}$$

NB for light stops

$$\Delta\rho^{\mathrm{stops}} \simeq 4\times 10^{-4}\left(\frac{500\mathrm{GeV}}{m_{\tilde{t}_1}}\right)^2$$

Getting 126 GeV

- In limit of large m_S, m_T, can integrate out adjoint scalars to obtain

$$m_h^2 \simeq M_Z^2 c_{2\beta}^2 + \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) s_{2\beta}^2 + \frac{3}{2\pi^2}\frac{m_t^4}{v^2}\left[\log\frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2} + \frac{\mu^2 \cot^2\beta}{m_{\tilde{t}_1} m_{\tilde{t}_2}}\left(1 - \frac{\mu^2\cot^2\beta}{12 m_{\tilde{t}_1} m_{\tilde{t}_2}}\right)\right]$$
$$+ v^2\left[\lambda_1 c_\beta^4 + \lambda_2 s_\beta^4 + 2(\lambda_3+\lambda_4+\lambda_5) c_\beta^2 s_\beta^2 + 4(\lambda_6 c_\beta^2 + \lambda_7 s_\beta^2) s_\beta c_\beta\right]$$
$$\overset{\tan\beta\to\infty}{\longrightarrow} M_Z^2 + \lambda_2 v^2 + \frac{3}{2\pi^2}\frac{m_t^4}{v^2}\log\frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2}$$

Can enhance the Higgs mass naturally!

- At small $\tan\beta$, do not need heavy stops or large stop mixing etc: for large λ_S or λ_T we can take just the tree-level part: $m_h^2 \simeq M_Z^2 c_{2\beta}^2 + \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) s_{2\beta}^2$ $\rightarrow \lambda_S \sim 0.7$ to obtain correct Higgs mass as at small $\tan\beta$ as in NMSSM/λSUSY.
- Also the origin of the potential may be a maximum rather than saddlepoint as in MSSM
- For large $\tan\beta$, scalar and triplet scalars can do the same job if they are heavy (e.g. for $\lambda_S = 1.8$ or $\lambda_T = 1.2$ with no stop contribution)

$$32\pi^2\lambda_2 \supset 2\lambda_S^4 \log\frac{m_S^2}{v^2} + (g_2^4 - 4g_2^2\lambda_T^2 + 10\lambda_T^4)\log\frac{m_T^2}{v^2}$$
$$+ \frac{4\lambda_S^2\lambda_T^2}{m_S^2 - m_T^2}\left[m_S^2\log\frac{m_S^2}{v^2} - m_T^2\log\frac{m_T^2}{v^2} - (m_S^2 - m_T^2)\right]$$

Dirac gauginos and a low SUSY scale

Let us see how a low SUSY scale affects this story.

- Bounds on squark masses are at about 1.38 TeV for the first two generations, and 700 GeV for the third. Now, naturalness of the Higgs mass demands

$$\sqrt{m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2} < 670 \text{ GeV} \frac{\sin\beta}{\sqrt{1+x_t^2}} \left(\frac{\log\frac{\Lambda}{\text{TeV}}}{\log 10}\right)^{-1/2} \frac{m_h}{125 \text{ GeV}} \left(\frac{\Delta}{20\%}\right)^{-1/2}$$

- The bounds on natural Dirac gluino masses are derived in the leading log approximation from the correction to the stop mass:

$$\delta m_H^2 = -\frac{3y_t^2}{8\pi^2}(m_Q^2 + m_U^2 + A_t^2) \log\frac{\Lambda}{\text{TeV}}$$
$$\delta m_{\tilde{t}}^2 = \frac{2g_s^2 m_{\tilde{g}}^2}{3\pi^2} \log\frac{m_O}{m_{\tilde{g}}}$$

- Hence

$$m_{\tilde{g}} < \frac{\pi^2 m_h}{y_t g_s} \frac{1}{\sqrt{\Delta}} \frac{1}{\sqrt{\log \Lambda/\text{TeV}}} \frac{1}{\sqrt{\log m_O/m_{\tilde{g}}}}$$
$$< 1800 \text{ GeV} \sin\beta \left(\frac{m_h}{125 \text{ GeV}}\right) \left(\frac{\Delta}{20\%}\right)^{-1/2} \left(\frac{\log\frac{\Lambda}{\text{TeV}}}{\log 10}\right)^{-1/2} \left(\log\frac{m_O}{m_{\tilde{g}}}\right)^{-1/2}$$

Goldstino couplings with Dirac gauginos

Recall that the operator for the Dirac gaugino mass is a holomorphic operator:

$$\int d^2\theta 2\sqrt{2} m_D \theta^\alpha \mathrm{tr}(W_\alpha \Sigma) \rightarrow -\frac{m_D}{4\sqrt{2} f^2}\int d^2\theta\, \overline{D}^2 D^\alpha (X^\dagger X)\, W^a_\alpha \Sigma^a$$

By extending X from a spurion to include (now a proto-)goldstino ψ_X, we can write down all of the couplings. Up to second order in ψ_X we have:

$$\begin{aligned}
\mathcal{L}_{Dirac} = & -m_D\lambda\chi + 2|m_D| D S_R + \frac{|m_D|}{f^2}\Big[-iDS_R(\partial_\mu\psi_X\sigma^\mu\overline{\psi}_X - \psi_X\sigma^\mu\partial_\mu\overline{\psi}_X) \\
& + \frac{1}{2} S_I F_{\mu\nu}\epsilon^{\mu\nu\rho\lambda}(\psi_X\sigma_\lambda\partial_\rho\overline{\psi}_X - \overline{\psi}_X\overline{\sigma}_\lambda\partial_\rho\psi_X) + S_R F_{\mu\nu}(\psi_X\sigma^\mu\partial^\nu\overline{\psi}_X - \overline{\psi}_X\overline{\sigma}^\mu\partial^\nu\psi_X)\Big] \\
& + \Big[\frac{m_D}{f}\Big(-\psi_X\lambda F_S - Si\lambda^\alpha\sigma^\mu_{\alpha\dot{\beta}}\partial_\mu\overline{\psi}^{\dot{\beta}}_X \\
& - iS\psi^\alpha_X\sigma^\mu_{\alpha\dot{\alpha}} D_\mu\overline{\lambda}^{\dot{\alpha}} + \frac{1}{\sqrt{2}}\psi_X\chi D - \frac{i}{2\sqrt{2}}\psi^\alpha_X(\sigma^\mu\overline{\sigma}^\nu)_\alpha^{\ \beta} F_{\mu\nu}\chi_\beta\Big) \\
& + \frac{im_D}{f^2}\Big((\chi\sigma^\mu\partial_\mu\overline{\psi}_X)\psi_X\lambda - (\chi\partial_\mu\psi_X)\lambda\sigma^\mu\overline{\psi}_X\Big) + \mathrm{h.c.}\Big].
\end{aligned}$$

From these generic couplings, we have – in principle – everything we need for phenomenological studies.

Observations

(a) The sgoldstino does not enter at second order.

(b) A vacuum expectation value for the adjoint scalar $v_S \equiv \langle S_R \rangle$ induces kinetic mixing between the goldstino and the gaugino

$$\mathcal{L} \supset -\frac{\sqrt{2} m_D v_S}{f} \left[i\lambda\sigma^\mu \partial_\mu \overline{\psi}_X + i\psi_X \sigma^\mu D_\mu \overline{\lambda} \right].$$

(c) The Dirac operator contains a coupling of two goldstinos to a gauge boson and the corresponding adjoint scalar. For a light enough scalar, this brings phenomenological signatures that are absent in the Majorana case – and therefore potentially interesting phenomenology.

Conventional signals

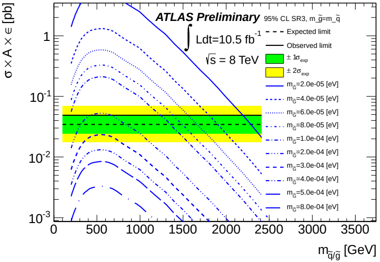
Gravitinos have been the subject of much experimental interest, involving one of two approaches:

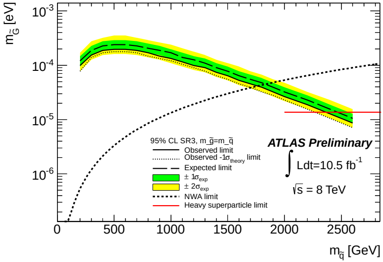
1. Consider all supersymmetric particles to be heavy and integrate them out, leaving only the SM fields and the gravitino. Place limits on $\sqrt{f}$ from the resulting higher-dimensional operators.
2. Make assumptions about the spectrum of superpartners and place limits on $\sqrt{f}$ as a function of their masses.

Approach (1) is dubious: current limit has $\sqrt{F} > 240$ GeV from monophoton searches at LEP, makes no sense given latest superparticle bounds.

Limits for Majorana gauginos

Approach (2): current best limits (from ATLAS-CONF 2012-147):





- Corresponds to $\sqrt{F} \gtrsim 700$ GeV for lighter squarks/gluinos (LEP limit is red line)
- One problem is that squark-gluino parameter space is potentially very large.
- Most importantly: assume a Majorana gaugino, which we claim should be revisited!

Sgluons

- In models with Dirac gauginos, we also have the scalar octet superpartners, the sgluons O.
- In typical explicit models (e.g. gauge mediation) these are potentially the heaviest particles in the theory
- However, if they are light then they can have very interesting phenomenology
- Current bounds are surprisingly weak: only about 800 GeV, and for its conventional decays.

This lead us to propose a new approach to low-energy SUSY searches:

1.5. Consider the effective theory of sgluons, gluinos and gravitinos, integrating out the squarks (except for perhaps the third generation).
This provides a much lower-dimensional parameter space with interesting phenomenology!

Conventional interactions of sgluons

The mass terms of the scalar octet are given by

$$\begin{aligned}\mathcal{L}_{M_O} &= -\, m_O^2 |O^a|^2 - \frac{1}{2}(B_O O^a O^a + h.c.) - (m_D O^a + m_D^* O^{a*})^2 \\ &= -\, \tilde{m}_O^2 |O^a|^2 - \frac{1}{2}(\tilde{B}_O O^a O^a + h.c.)\end{aligned}$$

where $\tilde{m}_O^2 = m_O^2 + 2|m_D|^2$ and $\tilde{B}_O = B_O + 2m_D^2$. The two mass eigenstates are then

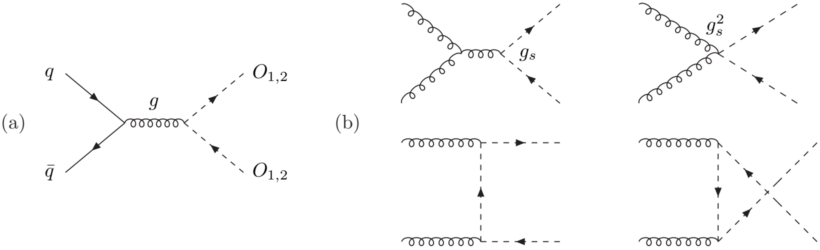
$$O_1^a = \frac{1}{\sqrt{2}}\left(e^{\frac{i}{2}\phi_O} O^a + e^{-\frac{i}{2}\phi_O} O^{a*}\right), \quad O_2^a = -\frac{i}{\sqrt{2}}\left(e^{\frac{i}{2}\phi_O} O^a - e^{-\frac{i}{2}\phi_O} O^{a*}\right),$$

with $M_{O_{1,2}}^2 = \tilde{m}_O^2 \pm |\tilde{B}_O|$.

- Note that if m_D dominates over m_O, B_O or $m_O \sim B_O$ then (e.g.) the pseudoscalar component can be much lighter than the scalar.

Octet tree couplings

The octet scalars have the usual gauge couplings and so can be produced in pairs at tree level:



Tree level decays

They have trilinear couplings with the squarks and gauginos

$$\mathcal{L}_{Dirac} = -\int d^2\theta \frac{m_D}{4\sqrt{2}f^2}\overline{D}^2 D^\alpha (X^\dagger X) W^a_\alpha \mathbf{O}^a \supset \sqrt{2} m_D (O^a + O^{a*}) D^a_c ,$$

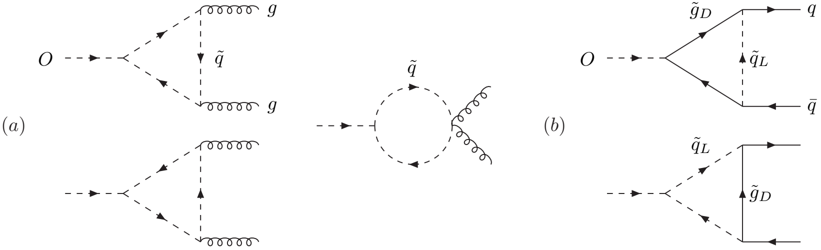
$$\rightarrow -2 g_s m_D T^a_{xy} \sum_{\tilde{q}_L, \tilde{q}_R} (\tilde{q}^*_{Lxi}\tilde{q}_{Lyi} - \tilde{u}^*_{Rxi}\tilde{u}_{Ryi} - \tilde{d}^*_{Rxi}\tilde{d}_{Ryi}) \left(\cos(\frac{\phi_O}{2}) O^a_1 + \sin(\frac{\phi_O}{2}) O^a_2\right)$$

$$\mathcal{L}_{Gauge} \supset i f^{abc} \overline{O}^b \lambda^a \chi^c + h.c.$$

These lead to rapid decays if the squarks or gluinos are ligher than half the octet mass $\rightarrow$ but this would mean rather heavy octets anyway.

Octet loop couplings

More interestingly, the above generate couplings at one loop with the quarks and gluons, which provide the conventional decay modes:



Loop couplings

- The widths to quarks are parametrised by

$$\mathcal{L} \supset c_{1\bar{t}t}\bar{t}O_1 t + c_{2\bar{t}t} i\bar{t}O_2\gamma_5 t,$$

- i.e. they split into scalar and pseudoscalar.
- The widths to gluons are given by

$$\Gamma(O_1 \to gg) = \frac{5\alpha_s^3}{192\pi^2}\frac{m_{D3}^2}{M_{O_1}}\cos^2(\frac{\phi_O}{2})|\lambda_{g_1}|^2, \quad \Gamma(O_2 \to gg) = \frac{5\alpha_s^3}{192\pi^2}\frac{m_{D3}^2}{M_{O_2}}\sin^2(\frac{\phi_O}{2})|\lambda_{g_2}|^2.$$

- In both cases the loop functions vanish for degenerate left and right-handed squarks – which we expect to be approximately true for flavour-blind low-energy SUSY breaking!
- If the B_O mass conserves CP, $\phi_O = 0$ and the pseudoscalar octet does not decay to gluons! From flavour constraint grounds we also expect this to be true ...
- Hence the conventional decays of the octet scalars should actually be much more suppressed than we would naively think.

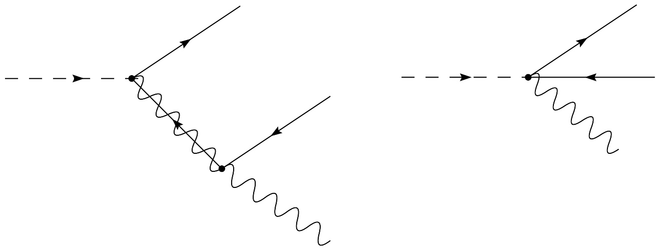
Octet goldstino decays

Now let us consider the new, goldstino, decay channels:

- If the octet is heavier than the gluino, then it will decay rapidly to that:

$$\Gamma(O_i \to \tilde{g} G) = \frac{(M_{O_i}^2 - m_D^2)^4}{32\pi f^2 M_{O_i}^3}$$

- If instead the gluino and squarks are heavy, then it can decay to a gluon and two goldstini via $O \to GGg$ and $O \to G\tilde{g} \to GGg$:



Octet goldstino couplings

- These processes involve highly non-renormalisable operators:

$$\mathcal{L}_{OGGg} = \frac{m_D}{f}\partial^\mu(G\sigma^\nu \overline{G})G^a_{\mu\nu}O^a_1 + \frac{m_D}{2f}\epsilon^{\mu\nu\rho\lambda}\partial_\rho(G\sigma_\lambda\overline{G})G^a_{\mu\nu}O^a_2\,,$$

,

- Naively this looks like the decay rate should increase as m_D increases! However, we also have the couplings for the other process:

$$\mathcal{L} \supset \Bigg(\frac{m_D}{\sqrt{2}f}G\sigma^{\mu\nu}\chi^a G^a_{\mu\nu} + \frac{i}{\sqrt{2}f}M^2_{O_2}O^a_2 G\chi^a - \frac{1}{\sqrt{2}f}(M^2_{O_1} - 2m^2_D)O^a_1 G\chi^a + h.c.$$

$$+ \frac{im_D}{\sqrt{2}f}\partial_\mu O_1(G\sigma^\mu\overline{\lambda} - \lambda\sigma^\mu\overline{G}) + \frac{m_D}{\sqrt{2}f}\partial_\mu O_2(\lambda\sigma^\mu\overline{G} + G\sigma^\mu\overline{\lambda})\,.$$

- They actually interfere so that m_D can decouple; we find ($y \equiv \frac{M^2_{O_i}}{m^2_D}$)

$$\Gamma(O_i \to \tilde{g}G) = \frac{(M^2_{O_i} - m^2_D)^4}{32\pi f^2 M^3_{O_i}},$$

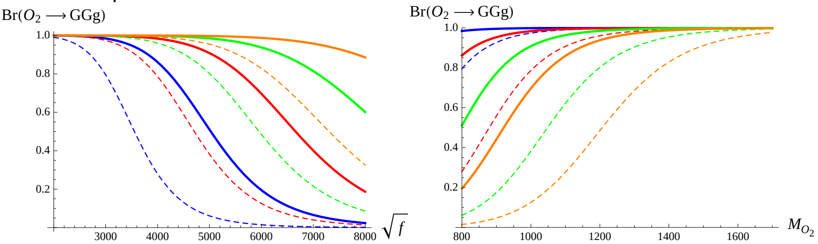
$$g(y) \equiv \frac{60\,(3-y)\,(1-y)^3\log(1-y)}{y^5} + \frac{6y^4 - 155y^3 + 480y^2 - 510y + 180}{y^4}$$

$$= \frac{2}{7}y^2 + \frac{3}{14}y^3 + \frac{1}{7}y^4 + \dots$$

$$g(1) = 1\,.$$

Branching ratios

So then we may have regions of parameter space where the goldstino decays of the octet are important:



Full lines are drawn for $\delta\tilde{m} = M_Z = 90$ *GeV and dotted lines for* $\delta\tilde{m} = 180$ *GeV.*
Also, we have taken $m_{\tilde{g}} = 1.7$ *TeV and* $m_{\tilde{q}} = 1.5$ *TeV.*
Left plot: (Blue, Red, Green, Orange): $M_{O_2} = (0.8$ *TeV,* 1 *TeV,* 1.2 *TeV,* 1.4 *TeV).*
Right plot: (Blue, Red, Green, Orange): $\sqrt{f} = (3$ *TeV,* 4 *TeV,* 5 *TeV,* 6 *TeV).*

Conventional searches for low-scale SUSY

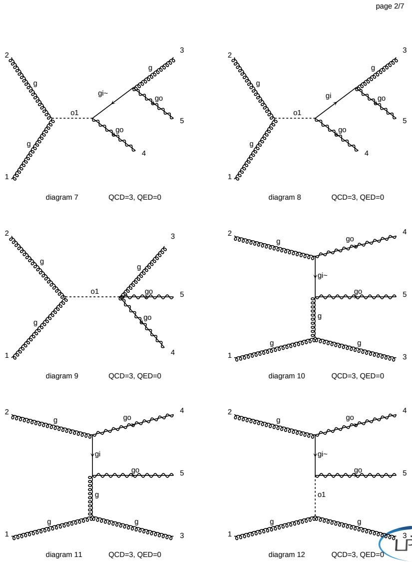
- The standard search channel for low-scale SUSY-breaking is monojet/monophoton events:

$$\begin{aligned} g + g &\to G + G + g \\ g + q &\to G + G + q \\ q + \overline{q} &\to G + G + \gamma/g \end{aligned} \tag{1}$$

- As I mentioned earlier, either we integrate out all of the SUSY particles and look at the model-independent effective operators à la Brignole, Feruglio, Mangano, Zwirner $\to$ does not matter whether gluinos are Dirac or Majorana, but unfortunately the bounds are too weak to be consistent.
- ... or we need to revist the standard Majorana case.
- Have argued that the theory with light octets is an interesting alternative.

Octet scalar monojet events

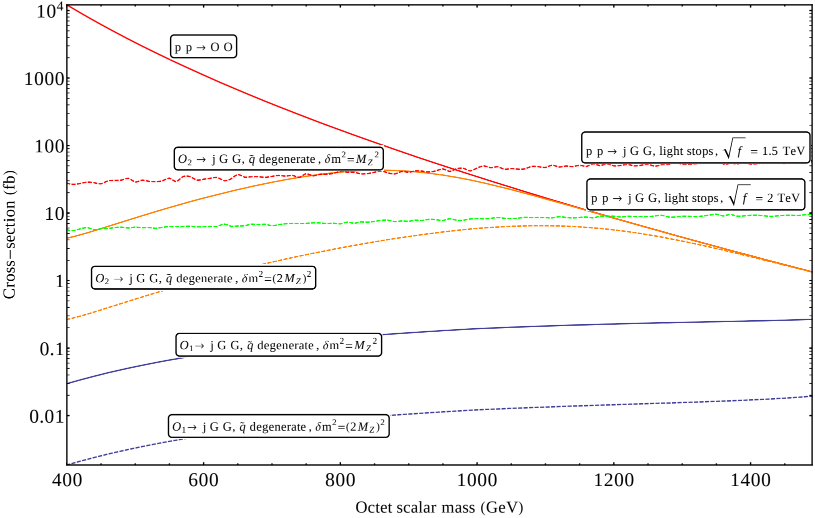
- There are many diagrams which contribute to these processes, but they are dominated at LHC by gluon fusion.
- Must include the full effective theory of scalar octet, gluino and goldstino.
- Crucially depends on the single octet production process via the loop-induced coupling.



Implementation

- As an illustration and precursor to a full study, we have implemented the model in `Feynrules` and then `MadGraph` and `CalcHEP`.
- Calculated cross-sections at LHC13 for goldstino events when one octet scalar is light as a function of the octet scalar mass, with the total double octet production cross-section given as reference.
- Events where two sgluons are produced and at least one decays to goldstinos (as opposed to two jets) are labelled $O_1 \to j\,G\,G$ and $O_2 \to j\,G\,G$.
- $\sqrt{f} = 7.5$ TeV was chosen since then $m_{\tilde{q}} \sim m_{\tilde{g}} \sim 0.2\sqrt{f} \sim 1.5$ TeV, with the squark masses varying from a common SUSY-breaking mass as $\sqrt{m_{\tilde{q}}^2 \pm \frac{1}{2}M_Z^2}, \sqrt{m_{\tilde{q}}^2 \pm 2M_Z^2}$.
- Monojet events are labelled $p\ p \to j\ G\ G$; for these, two different, lower, values of $\sqrt{f}$ are shown, and the spectrum of other sparticles has the first two generations of squarks and the right-handed squarks of the third generation at 2 TeV, with left-handed third-generation squarks at 755 GeV.
- In all cases the gluino mass was fixed at 1500 GeV.

Projected cross-sections



Conclusions

- Future searches to place limits on a low SUSY-breaking scale should be revisited in terms of a Dirac gluino.
- The phenomenology of octet scalars can then be important and potentially very interesting.
- The decays of octet scalars to goldstinos could dominate over their conventional decays – at the same time weakening already weak existing limits and allowing interesting monojet events.
- Hence we proposed a new effective scenario with few parameters with which to interpret these searches.

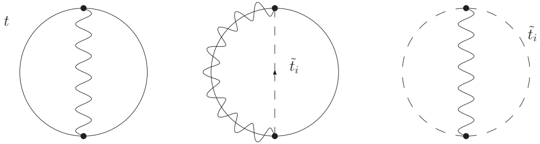
Two loop Higgs masses

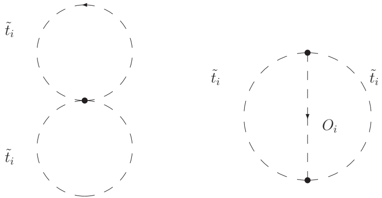
Work with P. Slavich

- In general, the tools are now available to study the phenomenology of Dirac gaugino masses to almost the same precision as the MSSM.
- Spectra at one loop, UFO files, etc can all be produced by `SARAH`.
- One of the major deficiencies is the lack of a two loop Higgs mass calculation.
- Have performed the analytic calculation of the leading (in $\alpha_s \alpha_t$) two loop corrections
- Have written a code (available on request) for numerical results.
- Have useful analytic approximations.

2 loop potential

We compute the two loop potential for the following diagrams:





There is one extra diagram compared to MSSM/NMSSM, and we now allow for two Majorana gaugino eigenstates/one Dirac.

Analytic results

Now have a 4×4 mass matrix, so the corrections are

$$\left(\Delta\mathcal{M}_S^2\right)_{11} = \frac{1}{2}\tilde{\mu}^2 s_{2\theta_t}^2 y_t^2 F_3 + \frac{A_t \tilde{\mu} y_t^2 \tan\beta}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{12} = -m_t \tilde{\mu} s_{2\theta_t} y_t^2 F_2 - \frac{1}{2} A_t \tilde{\mu} s_{2\theta_t}^2 y_t^2 F_3 + \frac{-A_t \tilde{\mu} y_t^2}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{22} = 2m_t^2 y_t^2 F_1 + 2A_t m_t s_{2\theta_t} y_t^2 F_2 + \frac{1}{2} A_t^2 s_{2\theta_t}^2 y_t^2 F_3 + \frac{A_t \tilde{\mu} y_t^2 \cot\beta}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{13} = \frac{1}{2} m_t \tilde{\mu} s_{2\theta_t}^2 y_t \lambda_S \cot\beta F_3 + \frac{-m_t y_t \lambda_S \left(-2\tilde{\mu}\cot\beta + A_t\right)}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{23} = -m_t^2 s_{2\theta_t} y_t \lambda_S \cot\beta F_2 - \frac{1}{2} A_t m_t s_{2\theta_t}^2 y_t \lambda_S \cot\beta F_3 + \frac{-A_t m_t y_t \lambda_S \cot\beta}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{33} = \frac{1}{2} m_t^2 s_{2\theta_t}^2 \lambda_S^2 \cot^2\beta F_3 + \frac{m_t^2 v_3^{-1} \lambda_S \cot\beta \left(\left(-\tilde{\mu} + v_3 \lambda_S\right)\cot\beta + A_t\right)}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{14} = \frac{1}{2} m_t \tilde{\mu} s_{2\theta_t}^2 y_t \lambda_T \cot\beta F_3 + \frac{-m_t y_t \lambda_T \left(-2\tilde{\mu}\cot\beta + A_t\right)}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\left(\Delta\mathcal{M}_S^2\right)_{24} = -m_t^2 s_{2\theta_t} y_t \lambda_T \cot\beta F_2 - \frac{1}{2} A_t m_t s_{2\theta_t}^2 y_t \lambda_T \cot\beta F_3 + \frac{-A_t m_t y_t \lambda_T \cot\beta}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} F$$

$$\vdots$$

Analytic approximations

We find analogues of the MSSM approximation formulae:

$$\begin{aligned}
\Delta m_h^2 =& \frac{3m_t^4}{2\pi^2 v^2}\left[\ln\frac{m_{\tilde t}^2}{m_t^2} + \frac{X_t^2}{m_{\tilde t}^2} - \frac{X_t^4}{12 m_{\tilde t}^4}\right] \\
&+ \frac{\alpha_s m_t^4}{\pi^3 v^2}\left\{\ln^2\frac{m_{\tilde t}^2}{m_t^2} - 2\ln^2\frac{m_t^2}{Q^2} + 2\ln^2\frac{m_{\tilde t}^2}{Q^2} + \ln\frac{m_{\tilde t}^2}{m_t^2} - 1 + \frac{X_t^2}{m_{\tilde t}^2}\left[1 - 2\ln\frac{m_{\tilde t}^2}{Q^2}\right] - \frac{X_t^4}{12 m_{\tilde t}^4}\right\} \\
&- \frac{2\alpha_s m_t^4}{\pi^3 v^2}\frac{c_{\phi_O}^2}{(1-4x_1)}\left\{1 - 4x_1 + 4x_1\ln\frac{m_{\tilde t}^2}{Q^2} - \ln\frac{m_{O_1}^2}{Q^2} + x_1\phi(x_1)\right\} \\
&+ \frac{2\alpha_s m_t^4}{\pi^3 v^2}\frac{c_{\phi_O}^2}{(1-4x_1)}\frac{X_t^2}{m_{\tilde t}^2}\left\{2x_1\ln x_1 - (1-4x_1)(\ln\frac{m_{O_1}^2}{Q^2} - 1) + 2x_1^2\phi(x_1)\right\} \\
&- \frac{\alpha_s m_t^4}{6\pi^3 v^2}\frac{c_{\phi_O}^2}{(1-4x_1)}\frac{X_t^4}{m_{\tilde t}^4}\left\{3(1-2x_1)x_1\ln x_1 + (1-4x_1)(1+3x_1 - \ln\frac{m_{O_1}^2}{Q^2}) + 6x_1^3\phi(x_1)\right\} \\
&+ (s_{\phi_O} \leftrightarrow c_{\phi_O}, m_{O_1} \leftrightarrow m_{O_2}, x_1 \leftrightarrow x_2)
\end{aligned}$$

where

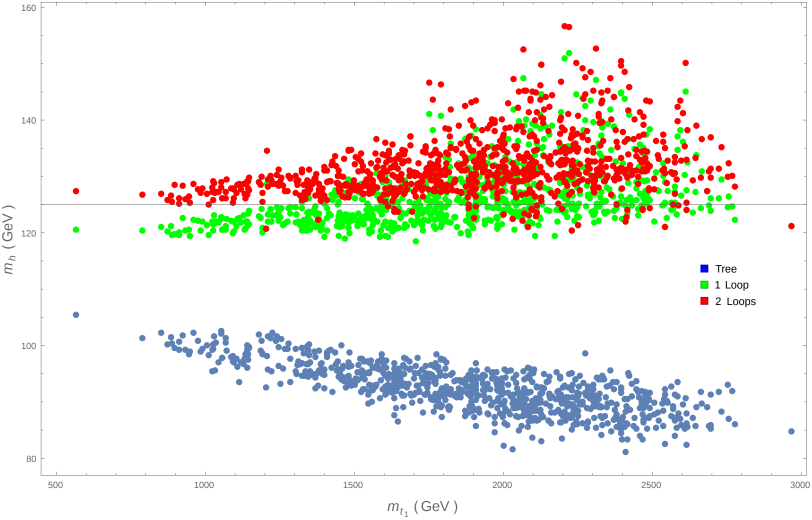
$$x_i \equiv \frac{m_{\tilde t}^2}{m_{O_i}^2}$$

$$\phi(x_i) \equiv \Phi(m_{\tilde t}^2, m_{\tilde t}^2, m_{O_i}^2).$$

and also a useful formula when the gaugino is heavy.

Example

From reprocessing some GUT models:



Conclusions

- Bottom line: the octet scalars do not induce a large correction; it is still mostly the stop contribution that matters.
- We do, however, find that the Higgs mass can be raised by around 5 GeV in various CMDGSSM benchmark points.

Two loop Higgs masses in general models

Work with F. Staub and K. Nickel

- In principle, the expressions for the two-loop effective potential for general field theories have been calculated by S. Martin in 2001.

There are then three approaches to using this to calculate the two loop higgs mass for <u>any</u> general model:

1. Numerically evaluate the effective potential and, by varying it, take the first and second derivatives.
2. Use analytic expressions for the derivatives of the loop functions, but numerically find the derivatives of the masses and couplings, composing them together via the chain rule.
3. Analytically compute the expression for the second derivatives in terms of couplings of the theory and new loop functions.

We have implemented all of these! They should appear in a forthcoming version of `SARAH` ...