Exploring Dirac Gaugino GUTs

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1206.6697;

...

Overview

- Motivation for Dirac gauginos
- Dirac gaugino models and the Higgs
- A new constrained GUT scenario

Dirac gauginos

- In the MSSM have Majorana gauginos described by one Weyl fermion λ in adjoint rep of each gauge group, mass term $\mathcal{L} \supset -\frac{1}{2}M_\lambda \lambda\lambda + \mathrm{h.c.}$
- To make give a Dirac mass, add an extra adjoint fermion χ to give mass term

$$\mathcal{L} \supset -m_D \chi\lambda + \mathrm{h.c.}$$

- This also requires a scalar Σ by supersymmetry, fit in an adjoint chiral multiplet (Σ, χ).

Motivation: bottom up

- 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
- Challenge is to study the possible spectra and Higgs properties

Motivation: 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])
- Adjoint multiplets appear in many UV models - would be nice to use them rather than throwing them away!

Status

Studying non-(N)MSSM SUSY models is typically hard due to lack of tools - and sometimes theory. However, now is the time to be doing this!
On the theory side,

- Dirac gauginos usually considered in context in gauge mediation; have explored many possibilities.
- Increasing numbers of people interested in this class of models (too many to mention), e.g. effect of Seiberg dualities, lepton number as R-symmetry, detailed studies of naturalness, ...
- We now understand the technical aspects well: RGEs, how the masses are generated, etc.
- However: despite many different models (not yet mapped out) there are no scenarios appropriate for collider studies such as the CMSSM $\rightarrow$ this talk.

On the tools side:

- We now have the tools for general theories: SARAH, PYR@TE, FeynRules, CalcHEP, MadGraph, MicrOmegas, ...
- Have been several studies (e.g. Martin and Kribs; Heikinheimo, Kellerstein, Sanz '11) of collider bounds for simplified models.
- Since 2012 SARAH has incorporated the possibility of Dirac gauginos.

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:

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

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

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

- 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

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

- $\rightarrow$ Bino mass is important for the Higgs mass, cannot be decoupled!

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RGEs

Since the operator is holomorphic, the RGEs are given by

$$\beta_{m_D^{iA}} = \gamma^i_j m_D^{jA} + \frac{\beta_g}{g} m_D^{iA}$$

We can apply the normal rules for the “standard” terms and then modify by

$$\Delta (m^2)^i_j = 2 m_D^{iA} m_{DjA} \equiv 2 (m_D^2)^i_j$$
$$\Delta B^{ij} = 2 m_D^{iA} m_D^{jA} \equiv 2 (m_D^2)^{ij}$$

at the desired energy scale. However, there is an exception: the tadpole RGE! The Dirac mass term enters explicitly here, so it had to be computed from scratch.

$$\beta^{(i)}_{t^a} \equiv X_S^{(i)} + X_\xi^{(i)} + X_D^{(i)}$$

$$(4\pi)^2 X_\xi^{(1)} = 2\sqrt{2} g_Y m_D^{aY} \mathrm{tr}(\mathcal{Y} m^2)$$
$$(4\pi)^4 X_\xi^{(2)} = 2\sqrt{2} g_Y m_D^{aY} \mathrm{tr}(\mathcal{Y} m^2 (4 g^2 C_2 - Y_2))$$

and

$$(4\pi)^2 X_D^{(1)} = 2\Big[(m_D^2)_{ef} (A^{aef} + M Y^{aef}) + Y_{efk} \mu^{ka} (m_D^2)^{ef} \Big]$$
$$(4\pi)^4 X_D^{(2)} = 4 (\beta^{(1)}_{m_D} / m_D)^f_g \Big[(m_D^2)_{ef} (A^{aeg} + M Y^{aeg}) + Y_{efk} \mu^{ka} (m_D^2)^{eg} \Big]$$

Naturalness

Bottom line:

- Tadpole term naturally generated by running, but not dangerous in size (in fact, it is useful phenomenologically): typically $\sim g_Y m_D \overline{m^2}$ (or smaller if $\mathrm{tr}(\mathcal{Y} m^2) = 0$ and $\mu = 0$).
- Dirac gaugino masses do not enter the Higgs or stop mass RGEs $\rightarrow$ increased naturalness: finite contribution to stop mass from gluino of

$$\delta m_{\tilde{t}}^2 = \frac{(m_{D3})^2 \alpha_s}{2\pi} \log\left(\frac{m_{O_P}}{m_{D3}}\right)^2$$

Singlet and triplet scalar masses drive EWSB, enter into naturalness bounds:

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

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

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

(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

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

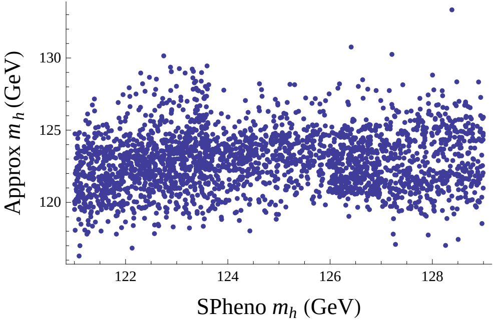
Can enhance the Higgs mass naturally!

- 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: $m_h^2 \simeq M_Z^2 c_{2\beta}^2 + \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) s_{2\beta}^2$ $\to \lambda_S \sim 0.7$ 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)

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

Numerical comparison

Comparison of effective potential with a SARAH/SPheno scan over models with heavy singlet and triplet scalars:



Can have light squarks and correct Higgs mass.

$\tan\beta = 50$, $m_{D2} = 600$ GeV, first two generation sfermion mass squareds of $4 \times 10^7\ (\mathrm{GeV})^2$, third generation sfermion mass squareds $4 \times 10^6\ (\mathrm{GeV})^2$ and scanning over $m_{D1} \in [-600, 600]$ GeV, $\mu \in [-750, 750]$ GeV, $B\mu \in [5000, 10^6]\ (\mathrm{GeV})^2$, $\lambda_T \in [0, 1]$ while adjusting λ_S to keep $m_h = 125 \pm 4$ GeV. The expectation values v_S, v_T were set by the tree level minimisation equation with $m_S^2, m_T^2 = 2.5 \times 10^7\ (\mathrm{GeV})^2$

Unification

- MSSM one-loop beta-function coefficients are $(b_3, b_2, b_1 = (5/3)b_Y) = (3, -1, -11)$, lead to unification of couplings at 10^{16} GeV with perturbative couplings $\alpha_{GUT} \sim 1/24$.

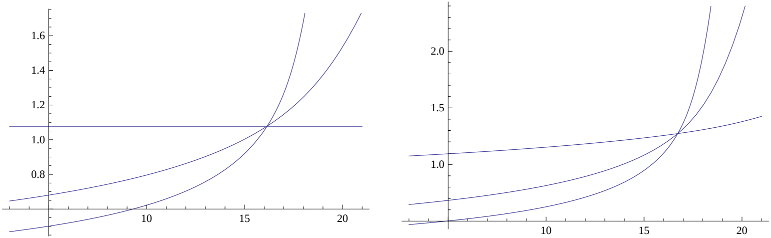
$$\frac{1}{g_i^2(\mu)} = \frac{1}{g_i^2(M_{SUSY})} + \frac{b_i}{8\pi^2} \log \mu/M_{SUSY}$$

- Triumph of the MSSM (modulo two-loop discrepancy...) that we might like to preserve!!
- Adding adjoint fields does (except for $\mathbf{S}$, a singlet): $\mathbf{T}$ decreases b_2 by 2, $\mathbf{O}_g$ decreases b_3 by 3

Our choice: add

$$(\mathbf{1},\mathbf{2})_{1/2} + (\mathbf{1},\mathbf{2})_{-1/2} + 2\times(\mathbf{1},\mathbf{1})_{\pm 1}$$

This could come from $(SU(3))^3$ (would need also four SM singlets).



One loop   Two loops

Toward a GUT scenario

- This configuration is particularly interesting. Let us add to the minimal Dirac gaugino MSSM a pair of doublets R_u, R_d and two non-chiral selectrons $\hat{E}_i, \hat{\hat{E}}_i, i = 1, 2$. The most general higgs potential is then

$$
\begin{aligned}
W \supset & (\mu + \lambda_S S) H_d H_u + 2\lambda_T H_d T H_u \\
& + (\mu_R + \lambda_{SR} S) R_u R_d + 2\lambda_{TR} R_u T R_d + \mu_{\hat{E}\,ij} \hat{E}_i \hat{\hat{E}}_j \\
& + (\mu_u + \lambda_{Su} S) R_u H_u + 2\lambda_{Tu} R_u T H_u + (\mu_d + \lambda_{Sd} S) R_d H_d + 2\lambda_{Td} R_d T H_d \\
& + Y_{\hat{E} i} R_u H_d \hat{E}_i + Y_{\hat{\hat{E}} i} R_d H_u \hat{\hat{E}}_i
\end{aligned}
$$

We can now take one of two directions:

- An extended MRSSM $\rightarrow$ removing $\mu, \mu_R, \lambda_S, \lambda_T$ and related couplings, where an R-symmetry is preserved by the Higgs sector.
- Charge the new fields under lepton number, so that we have new heavy vector-like leptons and sleptons. The superpotential becomes

$$
\begin{aligned}
W \supset & (\mu + \lambda_S S) H_d H_u + 2\lambda_T H_d T H_u \\
& + (\mu_R + \lambda_{SR} S) R_u R_d + 2\lambda_{TR} R_u T R_d + (\mu_{\hat{E}\,ij} + \lambda_{SEij} S) \hat{E}_i \hat{\hat{E}}_j \\
& + Y_{\hat{E} i} R_u H_d \hat{E}_i + Y_{\hat{\hat{E}} i} R_d H_u \hat{\hat{E}}_i \\
& + Y^{ij}_{LFV} L_i \cdot H_d \hat{E}_j + Y^{j}_{EFV} R_u H_d E_j
\end{aligned}
$$

Introducing the CMDGSSM

We can now specify a minimal set of boundary conditions at the GUT scale:

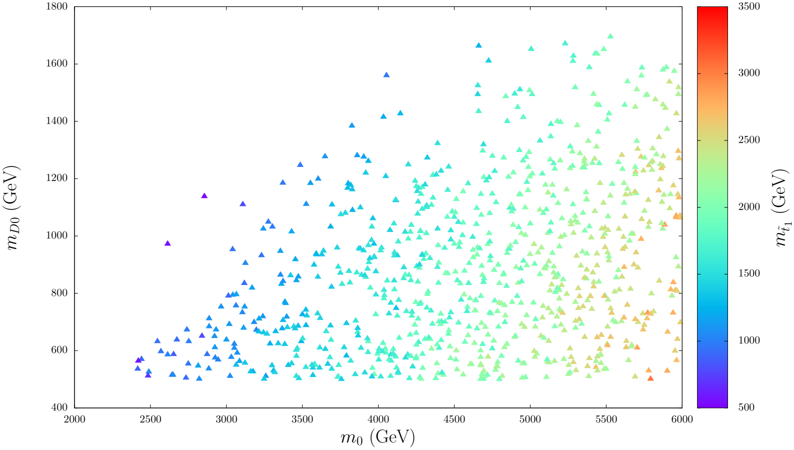
- As in the CMSSM/mSUGRA, we have $m_0, \tan\beta$ but instead of $m_{1/2}$ we have m_D. We set $A_0 = 0$ due to ~~SUSY~~ preserving R-symmetry.
- We also choose to take non-universal Higgs masses, and so specify μ, B_μ.
- Since we have two new tadpole conditions from v_S, v_T we specify m_{S0} (singlet scalar mass) and m_{T0} (triplet scalar mass) at the GUT scale. We set the octet scalar mass equal to the triplets, and take $B_T = B_S = B_O = 0$ for minimality.
- We have the Yukawa couplings $Y_{\hat{E}i}, Y_{\hat{\hat{E}}i}, Y^{ij}_{LFV}, Y^{j}_{EFV}$ which are equivalent to lepton Yukawas; they are constrained to be $\lesssim 0.01$ and so irrelevant for spectrum-generator purposes.
- We have a choice of $\mu_R, \mu_E \rightarrow$ can either adjust for precision gauge unification; set to be equal to the Higgs mu; set at convenient values. The Higgs mass and coloured sparticle spectrum is largely independent of this choice.
- We have a choice of couplings $\lambda_S, \lambda_T, \lambda_{SR}, \lambda_{TR}, \lambda_{SEij}$: can take $N = 2$ values, or $(SU(3))^3$ values, or choose freely.

First forays

- We can now start to explore the parameter space using SPheno code produced by SARAH for this model, which calculates two-loop RGEs and one-loop pole masses.
- One important technical limitation is due to the Higgs mass: if we enhance it using heavy stops, then the accuracy of the spectrum generator is no longer trustworthy.
- We instead choose to explore the corner of parameters space with $\lambda_S \sim 0.7$, small $\tan\beta$ so that no sparticle contributing significantly to the Higgs mass is heavier than about 2 TeV.
- For convenience we take $\lambda_T \sim 0$ and set $\mu_R \sim \mu_E \sim$ TeV, scan over $\lambda_S, \tan\beta$ within a narrow range and otherwise scan randomly over $\mu, B\mu, m_0, m_D, m_{S0}, m_{T0}$.
- We keep only points with the correct Higgs mass satisfying the constraints from HiggsBounds.

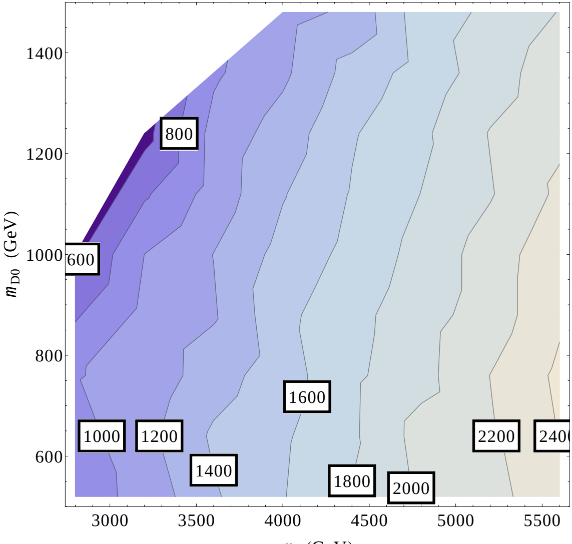
$m_0 - m_D$ plane and stop masses

We find no models in the with small m_0 and large m_D since the bino mass is important in the Higgs mass calculation. We do find many models with light stops:



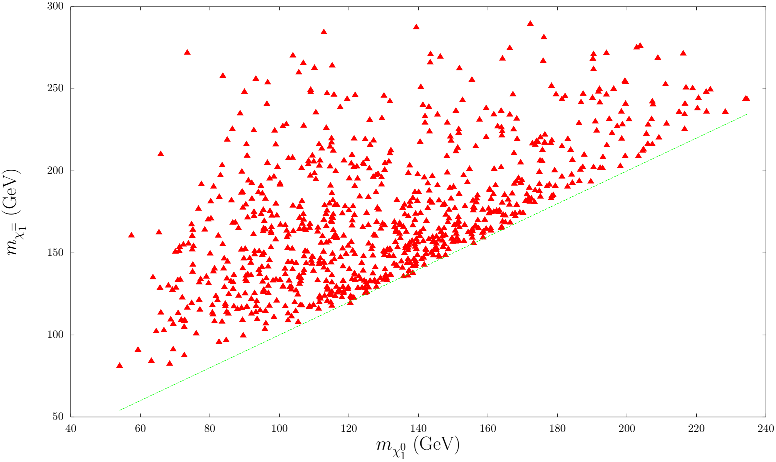
$m_0 - m_D$ plane II

Same plot as on previous slide, clearly showing contours of stop mass:



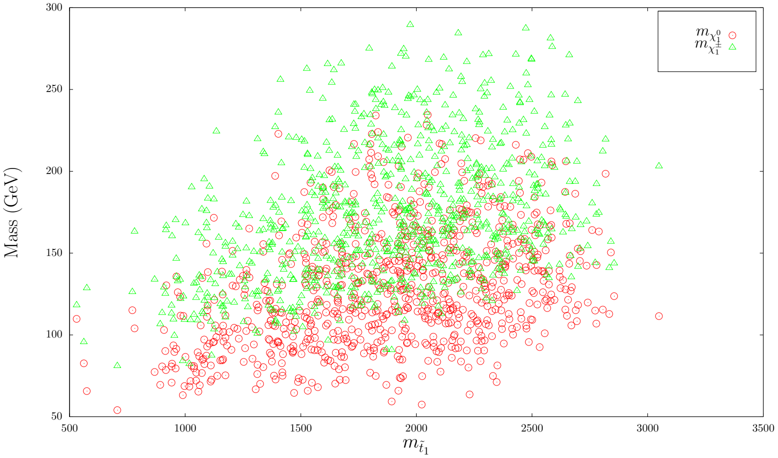
Charginos and neutralinos

The neutralinos are typically light due to the restriction on the bino mass from the Higgs mass. Also the scans preferentially find models with light μ, leading to light charginos.



Charginos, neutralinos and stops

Here we show the correlation of the neutralino and chargino masses with the lightest stop mass:



Predictions

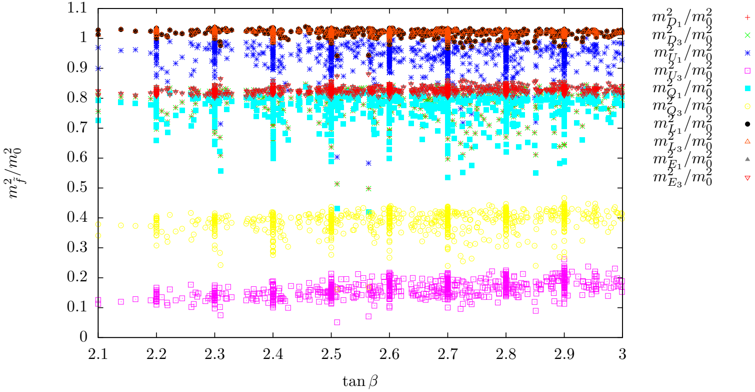
- Unification takes place at $(1.8 \pm 0.4) \times 10^{17}$ GeV
- We have a compressed pattern of soft masses (with deviations of a few percent):

$$m^2_{U33} : m^2_{Q33} : m^2_{Q11} : m^2_{Dii} : m^2_{Eii} : m^2_{U11} : m^2_{Lii}$$
$$= 0.16 : 0.39 : 0.77 : 0.79 : 0.83 : 0.93 : 1.02$$

- Hence sleptons are heavy and quasi-degenerate with the first two generations of squarks. This is because the Dirac gaugino masses do not enter into the squark RGEs.
- While the lightest stop masses are $1.9 \pm 0.5, 2.9 \pm 0.6$ TeV.
- The gaugino masses are in the ratio $0.22 : 0.9 : 3.5$, i.e. the Wino barely runs from m_D (as can be seen from the one-loop RGE, which is zero for small λ_T).

Squark masses

Over the range of $\tan\beta$ scanned, the squark masses vary little:



Conclusions

- Dirac gauginos have many attractive phenomenological and theoretical advantages over their Majorana counterparts, and can arise naturally in many different contexts (strong dynamics, higher dimensions, string theory, ...)
- There now exists a tool (`SARAH` with `SPheno`) to seriously study many aspects of their phenomenology which can interface with other tools.
- We have embedded DG models into a unified scenario and conducted a first study.
- Long program of research into these and other beyond-MSSM theories

Future Possibilities

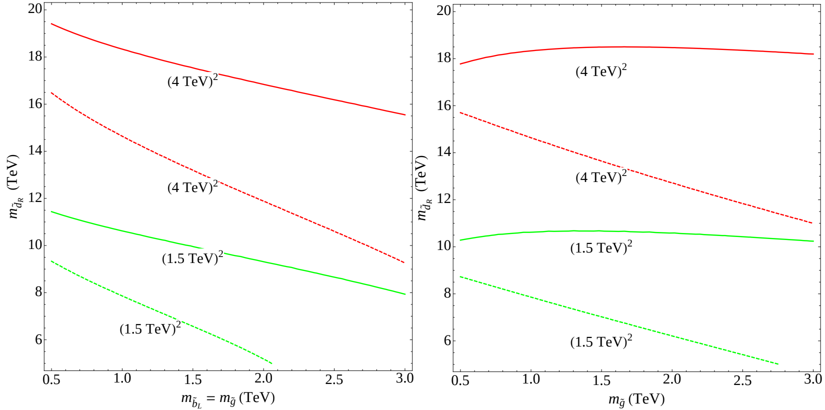
Many possible avenues for future work:

- Modifications of Higgs sector
- Calculation of two-loop effects and implementation in codes
- Connection with collider limits
- Models to realise messenger mass patterns
- Explicit D-term ~~SUSY~~ sectors (e.g. $4-1$ model)
- Warped models
- Gauge messengers
- Gravity mediation, embedding in string models, Dirac gravitinos,....

Flavour constraints

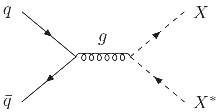
Work with E. Dudas, L. Heurtier, P. Tziveloglou

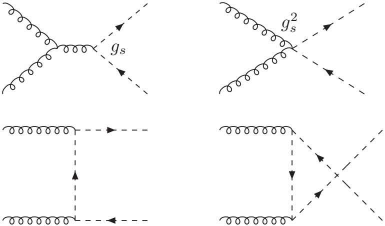
Dirac gauginos typically suppress flavour-chaging neutral currents. As an example, consider the model of Dudas, von Gersdorff, Pokorski, Ziegler 13]:



Collider Signatures

- May be difficult to distinguish directly Dirac and Majorana gauginos except at $e^- e^-$ collider
- Indirectly we would obtain spectacular signals from the adjoint scalars





- Decay at tree level (1210.4826 excludes $150 - 287$ GeV)

$$X \to \tilde{g}\tilde{g} \to qq\tilde{q}\tilde{q} \to qqqq + \tilde{\chi}\tilde{\chi}$$
$$X \to \tilde{q}\tilde{q} \to qq + \tilde{\chi}\tilde{\chi}$$

and (one loop):

$$X \to t\bar{t}$$

Higgs mass matrix

- Recall holomorphic coupling $\int d^2\theta 2\sqrt{2} m_D \theta^\alpha \mathrm{tr}(W_\alpha \Sigma) \supset -m_D(\lambda_a \chi_a) + \sqrt{2} m_D \Sigma_a D_a$ gives new D-term interactions
- This translates into a shift in D-term Higgs potential, masses and also the sfermion masses! Higgs mass matrix in the basis $\{h, H, S_R, T_R^0\}$ is

$$\begin{pmatrix} M_Z^2 + \Delta_h s_{2\beta}^2 & \Delta_h s_{2\beta} c_{2\beta} & \Delta_{hs} & \Delta_{ht} \\ \Delta_h s_{2\beta} c_{2\beta} & M_A^2 - \Delta_h s_{2\beta}^2 & \Delta_{Hs} & \Delta_{Ht} \\ \Delta_{hs} & \Delta_{Hs} & \tilde{m}_S^2 & \lambda_S \lambda_T \frac{v^2}{2} \\ \Delta_{ht} & \Delta_{Ht} & \lambda_S \lambda_T \frac{v^2}{2} & \tilde{m}_T^2 \end{pmatrix}$$

where

$$\Delta_h = \frac{v^2}{2}(\lambda_S^2 + \lambda_T^2) - M_Z^2$$
$$\Delta_{hs} = v[v_S \lambda_S^2 - g' m_{D1} c_{2\beta}]$$
$$\Delta_{Hs} = g' m_{1D} v s_{2\beta}$$

Higgs mass is not independent of bino mass! Although triplet scalar is heavy $\rightarrow$ roughly independent of Wino mass.

Scenarios

- Recall that in the MSSM we have two non-trivial tadpole conditions (from $\partial_{H_u^0} V = \partial_{H_d^0} V = 0$)
- In the NMSSM we have three (including the singlet vev $\partial_S V = 0$)
- Here we have four, from including the triplet scalar vev ($\partial_T V = 0$)
- $\rightarrow$ these should be traded with four soft parameters.

Can have several different scenarios, e.g. :

1. MSSM in disguise: here we shall allow a μ-term, and assume that the only source of R-symmetry violation arises in the supersymmetry-breaking sector, but permit only a B_μ term.
2. MSSM without μ term: this is the scenario of [Nelson, Rius, Sanz, Unsal], taking $\mu = 0$, essentially now ruled out.
3. Dynamical μ models: take $\mu = 0$ but allow substantial non-zero expectation value for the singlet, possibly via non-zero t_S.
4. Dynamical μ and $B\mu$ models: we allow a non-zero $W \supset \frac{1}{3}\kappa S^3$, breaking R-symmetry in the visible sector, but allowing μ and B_μ to be generated via a non-zero singlet vev as in the NMSSM.

Will mostly consider (1) and (3) scenarios in the following, leaving plenty of scope for future work.

Unification

- MSSM one-loop beta-function coefficients are $(b_3, b_2, b_1 = (5/3)b_Y) = (3, -1, -11)$, lead to unification of couplings at 10^{16} GeV with perturbative couplings $\alpha_{GUT} \sim 1/24$.

$$\frac{1}{g_i^2(\mu)} = \frac{1}{g_i^2(M_{SUSY})} + \frac{b_i}{8\pi^2} \log \mu/M_{SUSY}$$

- Triumph of the MSSM (modulo two-loop discrepancy...) that we might like to preserve!!
- Adding complete GUT multiplets (as in gauge mediation) does not alter this (beta-function coefficients decreased by $(1, 1, 1)$ per pair of $SU(5)$ messengers).
- Adding adjoint fields does (except for $\mathbf{S}$, a singlet): $\mathbf{T}$ decreases b_2 by 2, $\mathbf{O}_g$ decreases b_3 by 3

Four alternatives

1. Abandon matter and gauge unification
2. Modify our definition of "unification" ...
3. Add extra "bachelor" states to make up complete GUT adjoint multiplets [Fox, Nelson and Weiner, 02], allows matter and gauge unification
4. Add minimal extra states to restore gauge unification

Messengers to the Rescue

- Gauge mediation requires messenger fields - these could also restore gauge unification!
- Require at least 2 pairs of messengers in (anti) fundamental of $SU(2)$ and $SU(3)$ for adjoint scalar masses (see later)
- Easy to find sets of messengers that satisfy this, e.g.

$$\begin{array}{rcl} 4 \times [(1,1)_1 + (1,1)_{-1}] & \text{at} & m_1 = 3\ 10^{12}\text{GeV} \\ 4 \times [(1,2)_{1/2} + (1,\overline{2})_{-1/2}] & \text{at} & m_2 = 1.3\ 10^{13}\text{GeV} \\ 2 \times [(3,1)_{1/3} + (\overline{3},1)_{-1/3}] & \text{at} & m_3 = 10^{13}\text{GeV} \\ M_U \sim 9.9 \cdot 10^{17}\text{GeV} & & \alpha_U^{-1} \sim 4.77 \end{array}$$

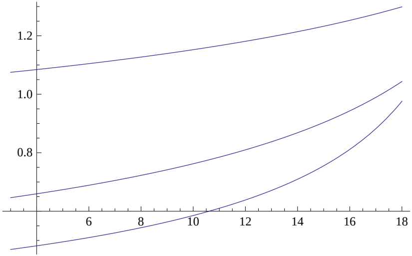
- High messenger scale required to allow perturbativity up to GUT scale

F-theory Unification

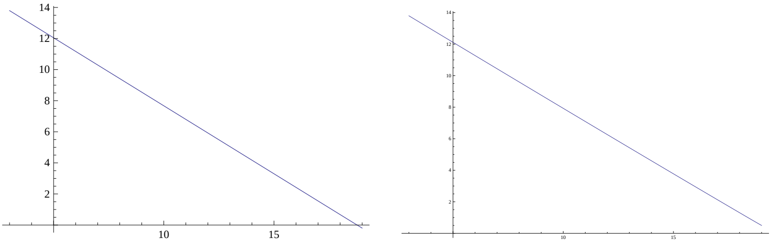
- If we modify our definition of unification, in the heterotic string, we could "unify" the hypercharge at a different Kac-Moody level.
- Alternatively, we can modify our definition of unification to the F-theory criterion

$$5\alpha_1^{-1} - 3\alpha_2^{-1} - 2\alpha_3^{-1} = 0$$

at a GUT scale of $\sim 1.7 \times 10^{17}$ GeV. This entails simply adding a vector-like pair of electron fields $(\mathbf{1},\mathbf{1})_{\pm 1}$ to the model, following [Davies, 2012]



F-theory Unification II



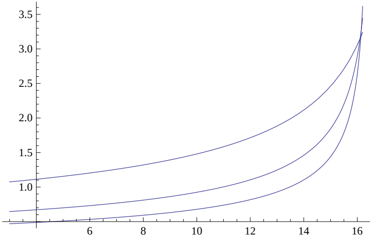
One loop Two loops

Unfortunately at two loops the couplings seem to unify beyond the Planck scale.

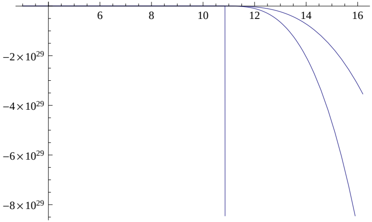
Other unification schemes

Can consider other schemes for unification with gravity mediation. E.g. Adding "bachelor" states at low scale

$$\mathbf{24} \rightarrow \mathbf{8_0} + \mathbf{3_0} + \mathbf{1_0} + (\mathbf{3}, \mathbf{2})_{-5/6} + (\bar{\mathbf{3}}, \mathbf{2})_{5/6}$$



One loop



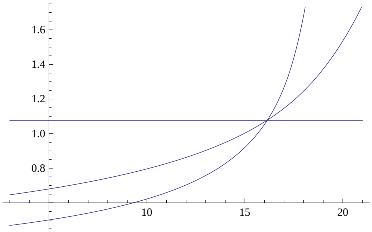
Two loops

Other unification schemes II

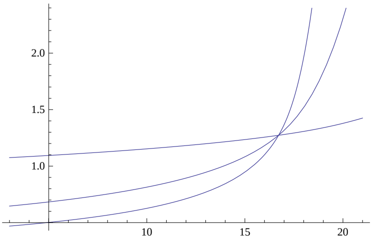
Or just add

$$(\mathbf{1},\mathbf{2})_{1/2}+(\mathbf{1},\mathbf{2})_{-1/2}+2\times(\mathbf{1},\mathbf{1})_{\pm 1}$$

This could come from $(SU(3))^3$ (would need also four SM singlets).



One loop



Two loops

RGEs cont'd

Bottom line:

- Tadpole term naturally generated by running, but not dangerous in size (in fact, it is useful phenomenologically): typically $\sim g_Y m_D \overline{m^2}$ (or smaller if $\mathrm{tr}(\mathcal{Y} m^2) = 0$ and $\mu = 0$).
- Dirac gaugino masses do not enter the Higgs or stop mass RGEs $\rightarrow$ increased naturalness: finite contribution to stop mass from gluino of

$$\delta m_{\tilde{t}}^2 = \frac{(m_{D3})^2 \alpha_s}{2\pi} \log\left(\frac{m_{O_P}}{m_{D3}}\right)^2$$

However, singlet and triplet scalar masses enter and help drive EWSB:

$$\begin{aligned}
16\pi^2 \frac{d}{dt} m_{H_u}^2 =& 6|y_t|^2[m_{Q_3}^2 + m_{U_3}^2 + m_{H_u}^2] \\
&+ 2\lambda_S^2[m_{H_u}^2 + m_S^2 + m_{H_d}^2] + 6\lambda_T^2[m_{H_u}^2 + m_T^2 + m_{H_d}^2] \\
&+ g_Y^2 \mathrm{Tr}(Y m^2) \\
16\pi^2 \frac{d}{dt} m_{H_d}^2 =& 6|y_b|^2[m_{Q_3}^2 + m_{D_3}^2 + m_{H_d}^2] \\
&+ 2|y_\tau|^2[m_{L_3}^2 + m_{E_3}^2 + m_{H_d}^2] \\
&+ 2\lambda_S^2[m_{H_u}^2 + m_S^2 + m_{H_d}^2] + 6\lambda_T^2[m_{H_u}^2 + m_T^2 + m_{H_d}^2] \\
&- g_Y^2 \mathrm{Tr}(Y m^2)
\end{aligned}$$

$$16\pi^2 \frac{d}{dt} B_\mu = B_\mu[3|y_t|^2 + 3|y_b|^2 + |y_\tau|^2 - 3g_2^2 - y_Y^2] + 2B_\mu\lambda_S^2 + 6B_\mu\lambda_T^2$$

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

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

NB for light stops

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