Off-trail SUSY

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II- L. Darmé, G. Policastro, Y. Oz

Program on Particle Physics at the Dawn of the LHC 13

Overview

In this talk, I want to briefly discuss two topics:

1. The FSSM: an attempt to deal with some Split SUSY shortcomings.

2. An example of super-Higgs mechanism for SUSY breaking by neither F nor D term.

SUSY "marked trails"

- SUSY is a symmetry of the UV theory.

- The particles content is that of the MSSM or *a slight modification (few additional states or symmetry enlargement).*

- SUSY breaking happens through an F or a D-term in some sector at an intermediate scale

- SUSY breaking is communicated to the "visible sector" inducing soft masses of order M_s.

The FSSM

A model with high M_S: Split SUSY

- Particle content = MSSM
- SUSY breaking scale = $M_S \equiv$ scale of scalar soft masses
- The two Higgs doublets mass matrix is fine-tuned so that one eigenvalue remains light: $\implies$ the SM higgs
- The gauginos and higgsinos are protected by an R-symmetry from acquiring a mass of order M_S ($\implies$ unification)

In summary:

- @ M_S: squarks + sleptons + heavy Higgs scalars.
- @ electroweak scale: Higgs boson, gauginos and higgsinos

(Arkani-Hamed, Dimopoulos)

Unification of gauge couplings

Searching for

- Extension of the Standard Model that predicts gauge coupling unification
- Has no new light scalars: would require an unjustifiable fine-tuning leading to much more complicated theories
- The only new particles near the electroweak scale are fermions (fermion masses can be easily protected by approximate continuous symmetries)

$\implies$ a simple realisation is Split SUSY (Giudice, Romanino)

Allowed range for M_S?

- In the IR: Visible sector SUSY breaking scale M_S cannot be arbitrarily low: because of LHC bounds, M_s must be bigger than a few TeV.

Allowed range for M_S?

- In the IR: Visible sector SUSY breaking scale M_S cannot be arbitrarily low

- In the UV: can M_S be arbitrarily high?

If M_S is as high as M_{GUT}, then the breaking of SUSY can be at tree-level.

Higgs mass

- Within the MSSM , the lightest Higgs mass is computable
- Here, the lightest SM-like Higgs boson mass M_h governed by two parameters, $M_h^2 = 2\lambda v_h^2$.
- From $v_h = 246 \text{ GeV}$ and $M_h \simeq 125 \text{GeV} \implies$ the value of λ.

- In the MSSM, λ is not arbitrary.
  At the SUSY scale M_S, the quartic Higgs self-coupling is given by the D-term:
  $$\lambda = \frac{1}{4}\left[g^2 + g'^2\right]\cos 2\beta + \Delta\lambda$$
  For a given M_S, is there a value of β corresponding to M_h?

How is the Higgs mass computed?

- Determine the values of gauge couplings at M_s by following their running from M_Z to M_S.
- Write $\lambda(M_s)$ as a function of the gauge couplings.

$$\lambda = \frac{1}{4}\left(g^2 + g'^2\right)\cos^2 2\beta + \Delta\lambda\,,$$

- compute β_λ for the evolution of λ described by the RGEs.
- Run backwards in energies from M_S to M_Z to obtain $\lambda(M_h)$ and compute the pole Higgs mass.

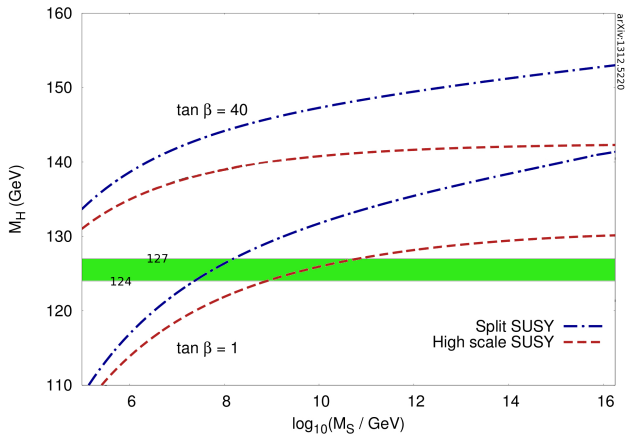For Split and High scale SUSY (N. Bernal, A. Djouadi and P. Slavich - G. F. Giudice and A. Strumia)

Upper bound on M_S in Split SUSY
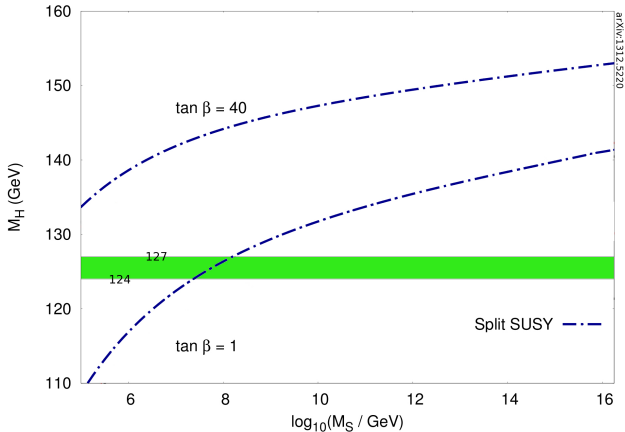
M_s needs to be lower than about 10^8 GeV

Upper bound on M_S in the High scale SUSY

M_s needs to be lower than about 10^{11}GeV

The Higgs quartic coupling evolution

The various contributions to β_λ at one-loop can be roughly classified as:

$$\beta_\lambda = \frac{1}{16\pi^2} \left[\underbrace{12\lambda^2 + \lambda(12y_t^2 + (\cdots \tilde{g}^2 \cdots) - (\cdots g^2 \cdots))}_{\equiv\ \beta_{\text{quartic}}} \right.$$

$$\left. + \underbrace{(\cdots g \cdots)^4}_{\equiv\ \beta_g} \underbrace{-(\cdots \tilde{g} \cdots)^4}_{\equiv\ \beta_{\tilde{g}}} \underbrace{-12y_t^4}_{\equiv\ \beta_t} \right] ,$$

where $(\cdots g \cdots)$ contains gauges contributions, $(\cdots \tilde{g} \cdots)$ contains contributions from Higgs-higgsinos-gauginos couplings.
SUSY fixes λ at M_s and evolve it down to the electroweak scale, positive contributions tend to bring λ towards lower values while negative contributions increases the final tree level Higgs mass.

A source of the obstruction for high M_S

- The β-function for the quartic Higgs coupling reads

$$\beta_\lambda = \frac{1}{16\pi^2} \left[\frac{9}{4}g^4 + \frac{3}{4}g'^4 - \frac{3}{2}g^2 g'^2 \right.$$
$$-12h_t^4$$
$$-5(\tilde{g}_u^4 + \tilde{g}_d^4) - (\tilde{g}_u'^2 + \tilde{g}_d'^2)^2$$
$$\left. -2\tilde{g}_u'^2 \tilde{g}_d'^2 - 2(\tilde{g}_u' \tilde{g}_u + \tilde{g}_d \tilde{g}_d')^2 \right]$$

- The Yukawa couplings Higgsino-gaugino-Higgs are fixed in the Split SUSY model to be those of the MSSM at M_S, and this leads for $M_S = M_{GUT}$ to:

$$\beta_\lambda = \frac{1}{16\pi^2}(-3.14g^4 - 12h_t^4)$$

- so there is a monotonic increase of λ in the IR.

Getting arbitrarily high M_S

- We saw that the value of λ increases monotonically when running from the UV to the IR.
- Even starting with the lowest value, $\lambda = 0$, it can reach unacceptably high values at the electroweak scale.

- Solution: change the RGE of λ
- There is a simple way to do it and allow for arbitrary M_S : the FSSM

Fake Split SUSY Model $\equiv$ FSSM

- For unification of couplings, only the quantum numbers of the particles count.
- Swap the higgsinos with copies that have Yukawa couplings ($\tilde{g}_{u,d} \sim 0$)
- The β-function for the quartic Higgs coupling reads

$$\beta_\lambda = \frac{1}{16\pi^2}\left[\frac{9}{4}g^4 + \frac{3}{4}g'^4 - \frac{3}{2}g^2 g'^2 - 12h_t^4\right]$$

- This leads for $M_s = M_{GUT}$ to:

$$\frac{1}{16\pi^2}(-3.14g^4 - 12h_t^4)|_{SplitSUSY} \rightarrow \frac{1}{16\pi^2}(3.42g^4 - 12h_t^4)|_{FSSM}$$

- Not monotonic increase of λ in the IR but careful to the instabilities $\lambda < 0$.
- An experimental discovery of a Split SUSY like fermionic spectrum at the TeV range would ask for measurement of the Yukawa coupling of the new fermions.

Fake Split SUSY Model $\equiv$ FSSM

- **FSSM-I**: both the higgsinos and gauginos are swapped for fake gauginos and fake higgsinos. (K.B, Darmé, Goodsell and Slavich — Dudas, Goodsell, Heurtier, Tziveloglou, 2014)

  1. The fermions remain light because of a $U(1)$ flavour symmetry
  2. The f-gauginos are Dirac partners of the gauginos
  3. Yukawa couplings $\tilde{g}_{u,d}, \tilde{g}'_{u,d}$ are suppressed by $(\text{TeV}/M_S)^2$
  4. Two pairs of vector-like electron superfields need to be added at M_S to insure unification at two-loop

- **FSSM-II**: Only the higgsinos are swapped for fake higgsinos. (K.B, Darmé, Goodsell 2015)

  1. The fermions remain light because of R-symmetry charges
  2. Yukawa couplings $\tilde{g}_{u,d}, \tilde{g}'_{u,d}$ are suppressed by (TeV/M_S)
  3. Two pairs of Higgs-like doublets needed for the f-higgsinos
  4. Two pairs of $(\mathbf{3}, \mathbf{1})_{1/3} \oplus (\bar{\mathbf{3}}, \mathbf{1})_{-1/3}$. In total we have added a vector-like pair of $\mathbf{5} + \bar{\mathbf{5}}$ of $SU(5)$

Higgs mass in the FSSM
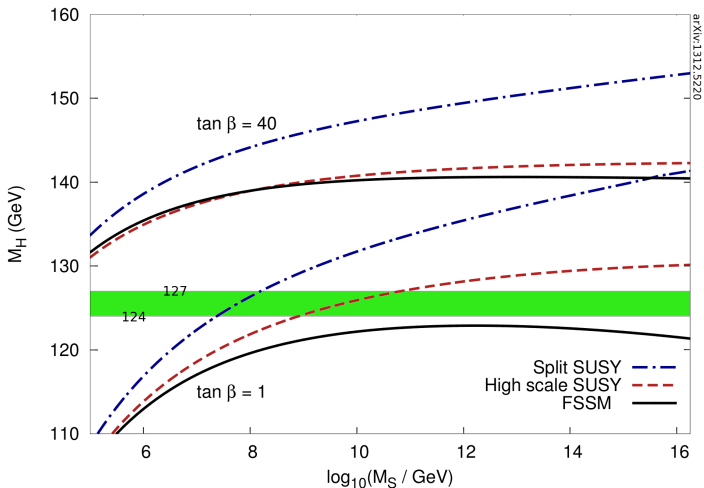
$M_{\tilde{g}'} = \mu = 2$ TeV and $\tan \beta = 1$ or 40

Split for Higgs

1. Both FSSM and High scale SUSY have vanishing $(\cdots \tilde{g} \cdots)$ terms. This decreases $\beta_\lambda(M_s)$ in the High scale SUSY and FSSM cases.

2. High scale SUSY has smaller gauge couplings than the FSSM and Split SUSY who have extra fermions below M_s pushing their couplings towards higher values.

A "domino" effect: at one-loop β_t has a positive contribution from Yukawas $(\cdots \tilde{g} \cdots)$ and a negative one from g_3. Since y_t is fixed at the electroweak scale, a smaller β_{y_t} means a smaller y_t at M_s. To summarise

- since a decrease in $\beta_{y_t}(M_s)$ implies a lower Higgs mass, we expect the FSSM to predict much lower Higgs mass than Split SUSY.

- while this effect is also present in High scale SUSY, it is mitigated by smaller gauge couplings . the difference between the pure gauge contribution being smaller than the contribution from extra $(\cdots \tilde{g} \cdots)$ terms. This explains why High scale SUSY still predict Higgs tree mass somehow lower than Split SUSY.

Split for Higgs

$\beta_\lambda(M_s)$

Split SUSY $-\cdot-\cdot-$
High scale SUSY $-\cdot-\cdot-$
FSSM ———

$\log_{10}(M_S \,/\, \mathrm{GeV})$
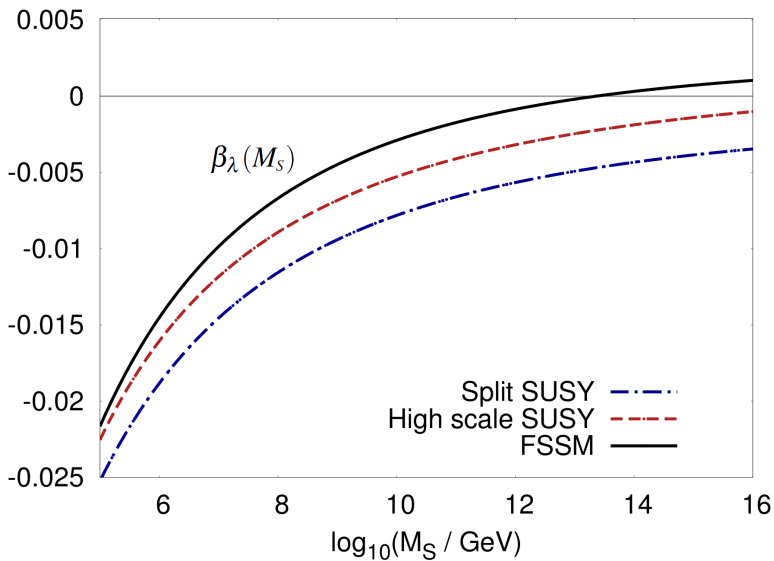
Split for Higgs
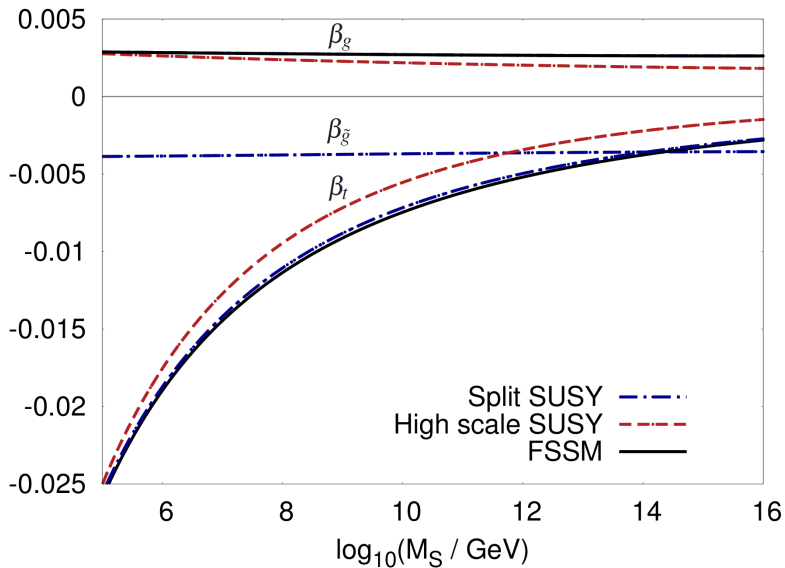
Mini-Split vs Mega-Split

At the SUSY scale, $\tan \beta$ is defined fixed by the requirement that one Higgs has a mass at the electroweak scale. In Split SUSY this fixes

$$\tan \beta \,=\, \sqrt{\frac{m_{H_d}^2 + |\mu|^2}{m_{H_u}^2 + |\mu|^2}} \,\simeq\, \sqrt{\frac{m_{H_d}^2}{m_{H_u}^2}}$$

since higgsinos are light. Renormalisation group evolution of $m_{H_d}^2$ and $m_{H_u}^2$ then naturally generates $\tan \beta \neq \mathcal{O}(1)$ and leads to increased Higgs mass. On the contrary, in FSSMs, we have

$$\tan \beta \,=\, \sqrt{\frac{m_{H_d}^2 + |\mu_d|^2}{m_{H_u}^2 + |\mu_u|^2}} \,,$$

with $\mu_d = \mu_u$ in the FSSM-I and $\mu_u, \mu_d \simeq M_s$ since they are unrelated to the low energy spectrum. We showed that this implies $\tan \beta = \mathcal{O}(1)$, further reducing the predicted Higgs mass compared to Split SUSY.

The gluino decay -I

Since F-gluinos decay must proceed via mixing with the usual gluinos, their decay rates are even more suppressed. $\tau_{\tilde{g}'}$, the F-gluino life-time is given by

$$\tau_{\tilde{g}'} \sim 4 \text{ sec} \times \left(\frac{M_S}{10^7 \text{GeV}} \right)^6 \times \left(\frac{1 \text{ TeV}}{m_{fg}} \right)^7 , \tag{1}$$

in the FSSM-I, where m_{fg} is the F-gauginos mass scale. While in the FSSM-II, since the gauginos are not fake, this enhancement does not occur and one is left instead with the Split SUSY gluino life-time

$$\tau_{\tilde{g}'} \sim 4 \text{ sec} \times \left(\frac{M_S}{10^9 \text{GeV}} \right)^4 \times \left(\frac{1 \text{ TeV}}{M_{1/2}} \right)^5 , \tag{2}$$

where $M_{1/2}$ is the gauginos mass scale.

The gluino decay -II

If one sticks with a standard cosmology the (F)-gluino lifetime is severely constrained, $\tau_{\tilde{g}} < 100$s A. Arvanitaki et al.. Big-Bang Nucleosynthesis (BBN), the CMB spectrum and the gamma-ray background ruled out relic (F)-gluinos with lifetime between 10^2 s until 10^{17} s. When the (F)-gluino is stable at the scale of the age of the universe, heavy-isotope searches also rule out such relic (F)-gluinos. This translates into limiting the SUSY scale to be below $5 \cdot 10^8$ GeV for the FSSM-I and $5 \cdot 10^{10}$ GeV for the FSSM-II.

A late time reheating occuring before BBN could for instance dilute gluino relic. In such case, heavy-isotope searches are so stringent that they still constraint $\tau_{\tilde{g}} \lesssim 10^{16}$s but one can avoid constraints from the CMB spectrum and the gamma-ray background, allowing therefore SUSY scales up to 10^{10} GeV for the FSSM-I and 10^{14} GeV for the FSSM-II.

Cosmology of FSSM

DM types	Inelastic scattering	Relic density	Gluino lifetime
$W|_{DM}$	None	$m_{\tilde{W}} \subset [2350, 2410]\text{GeV}$	For multi-TeV gluinos :
$B/H|_{DM}$	$\mu_{pole} \lesssim 900\text{GeV}$	$m_{\tilde{B}} \simeq \mu_{pole}$	$\begin{cases} M_S \lesssim 10^{10}\text{GeV (FSSM-II)} \\ M_S \lesssim 10^8\text{GeV (FSSM-I)} \end{cases}$
$H|_{DM}$	$\begin{cases} M_S \lesssim 10^8\text{GeV (FSSM-II)} \\ M_S \lesssim 10^6\text{GeV (FSSM-I)} \end{cases}$	$\mu_{pole} \subset [1110, 1140]\text{GeV}$	
FSSM-II = fake higgsinos		/	**FSSM-I = fake higgsinos+ fake gauginos**

Table: We impose a splitting between fake Higgsinos bigger than 300 keV to avoid direct detection through inelastic scattering, we require a gluino life-time smaller than 100 s to avoid hampering BBN and finally constraint the relic density to be $\Omega h^2 \subset [0.1172, 0.1226]$. When considering constraints on M_S, gauginos masses where taken in the multi-TeV range.

The slow gravitino

Lorentz symmetry and SUSY breaking

- Das and Kaku 1978: SUSY is broken by temperature
- Girardello, Grisaru, Solomonson 1981: SUSY is broken by temperature explicitly: non-periodic boundary conditions in imaginary time.
- Tesima 1983 Boyanovsky 1984: Goldstone fermion working in real-time.
- Aoyoma and Boyanovsky prove the Ward identities in 1984: Supersymmetry is broken spontaneously. Imaginary time explicitly breaks Lorentz symmetry therefore will show explicit breaking of SUSY.
- Klimov 1982, Weldon 1982, Aoyama 1986 and Ojima 1986: identification of pole with propagating mode in non-Lorentz invariant background.
- Further arguments for the goldstino: Leigh, Rattazzi 1995, Kotvun 2003 ...
- The spontaneous breaking of Lorentz symmetry comes with a Goldstone boson: the phonon.
- SUSY is also broken spontaneously (Ward identities of SUSY) with the appearance of a goldstino, a massless fermionic mode: the phonino.
- This spontaneous breaking does NOT happen through an F or a D-term.

SUSY breaking in a fluid

- A fluid background breaks Lorentz symmetry spontaneously (a vev for the energy-momentum tensor $T^{\mu\nu}$)

- SUSY is also broken spontaneously (Ward identities of SUSY) with the appearance of a goldstino, a massless fermionic mode: "the phonino" (D. Boyanovsky — R. G. Leigh, R. Rattazzi — P. Kovtun and L. G. Yaffe — C. Hoyos, B. Keren-Zur, Y. Oz).

- The spontaneous breaking does NOT happen through an F or a D-term.

- How does the super-Higgs mechanism work?

Done

- Written a Lagrangian of gravitino-goldstino interactions at the quadratic level

- Unitary gauge for a massive gravitino: super-Higgs mechanism

- Formula for the gravitino mass, dependence on ρ, p and M_P

- Identification in ψ_μ of the propagating d.o.f (constraint equations written)

- Equations of motion for longitudinal and transverse modes of spin 3/2?

K.B, Y. Oz, G. Policastro ⎯ K.B, L. Darmé, Y. Oz

To be done

- Written a Lagrangian of gravitino-goldstino interactions: higher orders ?
- Superfield formulation? (working on it)
- Fluid as a "hidden sector" for phenomenology: mediation of both breaking of SUSY and Lorentz
- Some eats are the same other similar to the ones of "slow gravitinos" in FRW (inflation, cosmology): (R. Kallosh, L. Kofman, A. D. Linde, A. Van Proeyen — G. F. Giudice, I. Tkachev, A. Riotto — Y. Kahn, D. A. Roberts, J. Thaler)
- ...

Phonino-gravitino interaction

The phonino equation of motion:

$$T^{\mu\nu}\gamma_\mu\partial_\nu G = 0 .$$

can be derived from the Lagrangian:

$$\mathcal{L}_G = \frac{i}{2\mathcal{T}^4} T^{\mu\nu} \bar{G}\gamma_\mu\partial_\nu G .$$

$\mathcal{T} \neq 0$ can be chosen to be $\mathcal{T} = |\mathrm{Tr}\ T^{\mu\nu}|^{\frac{1}{4}}$ or $\mathcal{T} = |\det\ T^{\mu\nu}|^{\frac{1}{16}}$.

In the flat Minkowski space-time limit, the Lagrangian describing the system phonino-gravitino at the quadratic order and the lowest order of an expansion in powers of the dimensionless parameter $\frac{\mathcal{T}}{M_p}$ reads:

$$\mathcal{L} = -\frac{i}{2}\epsilon^{\mu\nu\rho\sigma}\bar{\psi}_\mu\gamma^5\gamma_\nu\partial_\rho\psi_\sigma - \frac{1}{4}\epsilon^{\mu\nu\rho\sigma}n_{\sigma\lambda}\bar{\psi}_\mu\gamma^5\gamma_\rho\gamma^\lambda\psi_\nu - \frac{i}{\sqrt{2}}\frac{\mathcal{T}^2}{M_P}\frac{T^{\mu\nu}}{\mathcal{T}^4}\bar{\psi}_\mu\gamma_\nu G$$

$$+i\frac{T^{\mu\nu}}{2\mathcal{T}^4}\bar{G}\gamma_\mu\partial_\nu G + \frac{1}{4}\frac{T^{\mu\nu}n_{\mu\nu}}{\mathcal{T}^4}\bar{G}G .$$

SUSY transformations

This Lagrangian is invariant under the supersymmetry transformations with
Lorentz violating coefficients

$$\delta G = \sqrt{2}\mathcal{T}^2 \varepsilon \, ,$$
$$\delta \psi_\mu = -M_P(2\partial_\mu \varepsilon + i n_{\mu\nu}\gamma^\nu \bar{\varepsilon}) \, ,$$

if $n_{\mu\nu}$ satisfies:

$$-\frac{1}{2}\epsilon^{\mu\nu\sigma\rho}\epsilon_\rho^{\ \lambda\gamma\kappa}n_{\nu\lambda}n_{\sigma\gamma} = \frac{T^{\mu\kappa}}{M_P^2}$$

In the unitary gauge, G is set to zero through the supersymmetry
transformation:

$$\psi_\mu \to \psi_\mu + \frac{\sqrt{2}M_P}{\mathcal{T}^2}\partial_\mu G + i\frac{M_P}{\sqrt{2}\mathcal{T}^2}n_{\mu\nu}\gamma^\nu \bar{G}.$$

Case of a perfect fluid

Ideal fluid background specified by the energy density, pressure and velocity vector u^μ normalised as $u^\mu u_\mu = -1$,

$$T_{\mu\nu} = [p\eta_{\mu\nu} + (\rho + p)u_\mu u_\nu]$$

We will use the equation of state

$$w = \frac{p}{\rho}$$

Gamma matrices projections along time and space directions:

$$r^\mu = \mathbf{r}^{\mu\nu}\gamma_\nu$$
$$t^\mu = \mathbf{t}^{\mu\nu}\gamma_\nu$$

A mass formula for the slow gravitino

The Lagrangian describing the gravitino field takes the form:

$$\mathcal{L} = \frac{1}{2}\bar{\psi}_\mu\big[(\gamma^\mu\gamma^\nu + \eta^{\mu\nu})(-i\slashed{\partial} + m) + i\gamma^\nu\partial^\mu - i\gamma^\mu\partial^\nu$$
$$+ \frac{3\epsilon_{LV}m}{4 - 3\epsilon_{LV}}(r^\mu t^\nu + t^\mu r^\nu)\big]\psi_\nu$$

$$m = \frac{\sqrt{3\rho}}{4M_P}\,|\frac{1}{3} - w|$$

- For $w \neq -1$ both supersymmetry and invariance under Lorentz boosts are spontaneously broken
- Longitudinal mode propagates at a speed $|w|c$ slower than c.
- $\epsilon_{LV} = 0 \rightarrow F$ term breaking i.e. $w = -1$ and $T^{\mu\nu} = -|F|^2\eta^{\mu\nu}$

Spin 3/2

We construct a spin-3/2 field from the product of spin-1/2 and spin-1 states (a spinor-vector) denoted as ψ^μ. This is a reducible representation of the rotation group that can be decomposed into spin representation as

$$(\frac{1}{2}, \frac{1}{2}) \otimes (\frac{1}{2}, 0) = \frac{1}{2} \oplus (1 \otimes \frac{1}{2}) = \frac{1}{2} \oplus \frac{1}{2} \oplus \frac{3}{2} \ .$$

In general, we will therefore decompose ψ_μ into four spinors corresponding to the helicity-$\frac{3}{2}$ states $\psi^\mu_{\frac{3}{2}}$, the helicity-$\frac{1}{2}$ states $\psi_{\frac{1}{2}}$, and two remaining un-physical spinors that are projected out by two constraints. They are obtained from the equations of motion by contracting with either u_μ to extract the temporal part, or by the derivative operator $\nabla_\mu - N_\mu/2$ and read

$$[r^\mu r^\nu - \mathbf{r}^{\mu\nu}] \, \partial_\mu \psi_\nu \ = \ -\frac{m}{1 - \frac{3}{4}\epsilon_{LV}} r^\rho \psi_\rho \ , \tag{3}$$

and

$$(w r^\nu - t^\nu)\psi_\nu \ = \ 0 \tag{4}$$

Spin 3/2

One can use these constraints to obtain the physical degrees of freedom $\psi_{\frac{3}{2}}^{\mu}$ and $\psi_{\frac{1}{2}}$. We can obtain them directly from ψ_{μ} by

$$\psi_{\frac{1}{2}} = \sqrt{\frac{3}{2}} \frac{m}{k(1 - 3\epsilon_{LV}/4)} \not{u}\, r^{\rho}\psi_{\rho} \tag{5}$$

$$\psi_{\frac{3}{2}}^{\mu} = \mathcal{P}_{\frac{3}{2}}^{\mu\nu}\psi_{\nu} \equiv \left[\eta^{\mu\nu} - \frac{1}{3}r^{\mu}r^{\nu} - t^{\mu}t^{\nu} - \frac{1}{6}(r^{\mu} - 3\frac{\not{k}k^{\mu}}{k^2})(r^{\nu} - 3\frac{\not{k}k^{\nu}}{k^2})\right]\psi_{\nu} .$$

Note that the spin-1/2 degrees of freedom are proportional to $r^{\mu}\psi_{\mu}$ up to a normalisation factor and a gamma matrix. These factors are crucial in obtaining a correctly normalised fermionic kinetic term for $\psi_{\frac{1}{2}}$. The equations of motion for these fields then take the form

$$\begin{aligned}
(t^{\rho}\partial_{\rho} + r^{\rho}\partial_{\rho} + m)\psi_{\frac{3}{2}}^{\mu} &= 0 , \\
(t^{\rho}\partial_{\rho} - wr^{\rho}\partial_{\rho} + m)\psi_{\frac{1}{2}} &= 0 .
\end{aligned} \tag{6}$$

Spin 3/2

The propagator for the gravitino and polarisations

$$G^{\mu\nu} = \frac{\Pi_{\frac{3}{2}}^{\mu\nu}}{p^2 + m^2} + \frac{\Pi_{\frac{1}{2}}^{\mu\nu}}{w^2 k^2 + q^2 + m^2} - \frac{3}{4}\epsilon_{LV}\frac{\slashed{k}}{mk^2}(t^\mu k^\nu - k^\mu t^\nu) . \qquad (7)$$

$$\Pi_{\frac{3}{2}}^{\mu\nu} = (m - i\slashed{p})\mathcal{P}_{\frac{3}{2}}^{\mu\nu} ,$$

and

$$\Pi_{\frac{1}{2}}^{\mu\nu} = \frac{2}{3}\, \Lambda^\mu \,(i\slashed{p} - \epsilon_{LV}i\slashed{k} + m)\, \Lambda^\nu ,$$

where

$$\Lambda^\mu = \gamma^\mu - i\frac{p^\mu}{n} - \frac{3}{2}(r^\mu - \frac{\slashed{k}k^\mu}{k^2}) - \frac{3}{4}\epsilon_{LV}t^\mu ,$$

In the limit of gravitino high momentum where we have the hierarchy

$$m \ll |p| \ll f ,$$

the propagator simplifies to

$$G^{\mu\nu} \to -\mathcal{P}_{\frac{3}{2}}^{\mu\nu}\frac{i\slashed{p}}{p^2} - \frac{2}{3}\frac{p^\mu p^\nu}{n^2}\frac{i\slashed{q} - iw\slashed{k}}{q^2 + w^2 k^2} . \qquad (8)$$

Conclusions

- We presented a framework (FSSM) where the Supersymmetry breaking scale can be high and unification of couplings is preserved. It does not suffer from some of the drawbacks of Split SUSY.

- Generalisation of F, D term breaking to breaking. A new Lagrangian describing the propagation of a fluid in a non-Lorentz invariant background is presented.

Acknowledgments

In this talk:

- the FSSM
  JHEP 1405 (2014) 113
  K.B., L. Darmé, M.D. Goodsell, P. Slavich.
  arXiv:1508.02534 [hep-ph] JHEP XXXX
  K.B., L. Darmé, M.D. Goodsell

- Beyond F and D term SUSY breaking: The Slow Gravitino
  JHEP 1410 (2014) 121   K.B., L. Darmé, Y. Oz
  JHEP 1402 (2014) 015   K,B., Y. Oz, G. Policastro

FSSM- I

The adjoint chiral superfields are called "fake gauginos" (henceforth F-gauginos). They consist of a set of chiral multiplets, namely a singlet $\mathbf{S} = S + \sqrt{2}\theta\chi_S + \dots$; an $SU(2)$ triplet $\mathbf{T} = \sum_a \mathbf{T}^a \sigma^a/2$, where $\mathbf{T}^a = T^a + \sqrt{2}\theta\chi_T^a + \dots$ where σ^a are the three Pauli matrices; and an $SU(3)$ octet $\mathbf{O} = \sum_a \mathbf{O}^a \lambda^a/2$, where $\mathbf{O}^a = O^a + \sqrt{2}\theta\chi_O^a + \dots$ and λ^a are the eight Gell-Mann matrices.

Unification is jeopardised if one does not add further fields since the F-gaugino multiplets do not fill complete representations of a GUT group. An easy way to recover unification is to add two pairs of vector-like right-handed electron superfields ($\mathbf{E}'_{1,2}$ in $(\mathbf{1},\mathbf{1})_1$ and $\mathbf{E}'_{1,2}$ in $(\mathbf{1},\mathbf{1})_{-1}$) and one pair of $SU(2)$ doublets ($\mathbf{H}'_d$ in $(\mathbf{1},\mathbf{2})_{1/2}$ and $\mathbf{H}'_u$ in $(\mathbf{1},\mathbf{2})_{-1/2}$). In this work, the latter become fake Higgs doublets (henceforth F-Higgs) and their fermionic components fake higgsinos (henceforth F-higgsinos) rather than, for example, assigning them lepton number.

FSSM- I

An essential difference with usual Dirac gaugino models is that we do not impose an *R*-symmetry which forbids Majorana gaugino masses leading to the same mass for gauginos and F-gauginos. Instead, we keep only the F-gauginos light thanks to an approximate $U(1)_F$ flavour symmetry with the following charge assignments

Superfield	$U(1)_F$ charge
$\mathbf{H'_u}, \mathbf{H'_d}; \mathbf{S}, \mathbf{T}, \mathbf{O}$	1
$\mathbf{E'}_{1,2}, \mathbf{\bar{E}'}_{1,2}$	0

All other (MSSM) multiplets are neutral under $U(1)_F$. We parametrise the breaking of this symmetry by a small number ε which could be considered, as standard in flavour models, to come from the expectation value of a field (divided by some UV scale); in this case we can suppose it to have charge -1 under $U(1)_F$.

FSSM- I

The superpotential contains a hierarchy of couplings due to suppressions by different powers of ε:

$$
\begin{aligned}
W \supset\ & W_{\text{unif}} + \mu_0\,\mathbf{H_u}\cdot\mathbf{H_d} + Y_u\,\mathbf{U^c}\,\mathbf{Q}\cdot\mathbf{H_u} - Y_d\,\mathbf{D^c}\,\mathbf{Q}\cdot\mathbf{H_d} - Y_e\,\mathbf{E^c}\,\mathbf{L}\cdot\mathbf{H_d} \\
& + \varepsilon\left(\hat{\mu}'_d\,\mathbf{H_u}\cdot\mathbf{H'_d} + \hat{\mu}'_u\,\mathbf{H'_u}\cdot\mathbf{H_d} + \hat{Y}'_u\,\mathbf{U^c}\,\mathbf{Q}\cdot\mathbf{H'_u} - \hat{Y}'_d\,\mathbf{D^c}\,\mathbf{Q}\cdot\mathbf{H'_d} - \hat{Y}'_e\,\mathbf{E^c}\,\mathbf{L}\cdot\mathbf{H'_d}\right) \\
& + \varepsilon\left(\hat{\lambda}_S\,\mathbf{S}\,\mathbf{H_u}\cdot\mathbf{H_d} + 2\,\hat{\lambda}_T\,\mathbf{H_d}\cdot\mathbf{T}\,\mathbf{H_u}\right) \\
& + \varepsilon^2\left(\hat{\lambda}'_{Sd}\,\mathbf{S}\,\mathbf{H_u}\cdot\mathbf{H'_d} + \hat{\lambda}'_{Su}\,\mathbf{S}\,\mathbf{H'_u}\cdot\mathbf{H_d} + 2\,\hat{\lambda}'_{Tu}\,\mathbf{H_d}\cdot\mathbf{T}\,\mathbf{H'_u} + 2\,\hat{\lambda}'_{Td}\,\mathbf{H'_d}\cdot\mathbf{T}\,\mathbf{H_u}\right) \\
& + \varepsilon^2\,\hat{\mu}''\,\mathbf{H'_u}\cdot\mathbf{H'_d} + \varepsilon^2\left[\frac{1}{2}\,\hat{M}_S\,\mathbf{S}^2 + \hat{M}_T\,\mathrm{Tr}(\mathbf{TT}) + \hat{M}_O\,\mathrm{Tr}(\mathbf{OO})\right] + \mathcal{O}(\varepsilon^3)\,, \quad (9)
\end{aligned}
$$

where $\mathbf{Q}$, $\mathbf{U^c}$, $\mathbf{D^c}$, $\mathbf{L}$ and $\mathbf{E^c}$ are the quarks and leptons superfields, $\mathbf{H_u}$ and $\mathbf{H_d}$ the usual MSSM two Higgs doublets. We have explicitly written the ε factors so that all mass parameters are expected to be generated at M_S and all dimensionless couplings are either of order one or suppressed by loop factors. The additional superpotential W_{unif} contains the interactions involving the pairs $\mathbf{E}'_{1,2}$ and $\mathbf{E}'_{1,2}$; these fields are irrelevant for the low energy theory because their masses are not protected, so are of order M_S.

FSSM- I

We shall not explicitly write all of the soft terms in the model for reasons of brevity, since they can simply be inferred from the flavour assignments. For example, for the gauginos, allowing all terms permitted by the symmetries we have unsuppressed Majorana masses for the gauginos, and then the suppressed Majorana masses for the F-gauginos $\varepsilon^2 \hat{M}_{S,T,O}$ – and Dirac masses mixing the two suppressed only by ε, giving a generic mass matrix of

$$\mathcal{M}_{1/2} \sim \mathcal{O}(M_s) \begin{pmatrix} 1 & \mathcal{O}(\varepsilon) \\ \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon^2) \end{pmatrix}.$$

We have a heavy eigenstate of mass $\mathcal{O}(M_s)$ and a light one, the F-gaugino at leading order, of mass $\mathcal{O}(\varepsilon^2 M_s)$. Requiring that the F-gauginos have a mass at the TeV scale (for unification and, as we shall later see, dark matter) then fixes ε:

$$\varepsilon = \mathcal{O}(\sqrt{\frac{\text{TeV}}{M_s}}) . \tag{10}$$

For the adjoint scalars we shall define the explicit soft terms:

$$-\mathcal{L}_{\text{scalar soft}} \supset m_S^2 |S|^2 + 2m_T^2 \text{tr } T^\dagger T + 2m_O^2 \text{tr } O^\dagger O$$
$$+ \frac{1}{2}\varepsilon^2 B_S[S^2 + h.c] + \varepsilon^2 B_T[\text{tr } TT + h.c] + \varepsilon^2 B_O[\text{tr } OO + h.c] .$$

FSSM- I

The Higgs mass matrix can be written in terms of the four-vector
$v_H \equiv (H_u, H_d^*, H'_u, H'^*_d)$ as

$$-\frac{1}{M_S^2} \mathcal{L}_{soft} \supset v_H^\dagger \begin{pmatrix} \mathcal{O}(1) & \mathcal{O}(1) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) \\ \mathcal{O}(1) & \mathcal{O}(1) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) \\ \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(1) & \mathcal{O}(\varepsilon^2) \\ \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon^2) & \mathcal{O}(1) \end{pmatrix} v_H. \tag{12}$$

In the spirit of Split SUSY we tune the weak scale to its correct value and
define the SM-like Higgs boson H as

$$H_u \approx \sin\beta\, H + \dots, \qquad H_d \approx \cos\beta\, i\sigma^2 H^* + \dots,$$
$$H'_u \approx \varepsilon\, H + \dots, \qquad H'_d \approx \varepsilon\, i\sigma^2 H^* + \dots,$$

where β is a mixing angle and the ellipses represent terms at higher order in
ε. In particular, we see that at leading order H only has components in the
original Higgs doublet. This means that the matter Yukawa couplings will
have the same structure as in Split-SUSY at low energy. Furthermore, the
presence of a light SM-like Higgs implies at first order in ε

$$B_\mu \simeq \sqrt{(m_{H_u}^2 + \mu_0^2)(m_{H_d}^2 + \mu_0^2)} + \mathcal{O}(\varepsilon). \tag{13}$$

FSSM- II

Since we do not have fake gauginos, the ultraviolet model building is much more conservative than the FSSM-I; in particular one does not have to appeal to Dirac gauginos. Instead, we just add two pairs of Higgs-like doublets, $\mathbf{H}'_u$, $\mathbf{H}'_d$ and $\mathbf{R}_u$, $\mathbf{R}_d$. Unification of the gauge couplings at one-loop above M_s is recovered by adding two pairs of supermultiplets in the representations $(\mathbf{3}, \mathbf{1})_{1/3} \oplus (\bar{\mathbf{3}}, \mathbf{1})_{-1/3}$. In total, we have therefore added two vector-like pairs of $\mathbf{5} + \bar{\mathbf{5}}$ of $SU(5)$. This should be reminiscent of gauge mediation scenarios, except that here the doublets mix with the Higgs fields. In order to create a split spectrum, we introduce an approximate R-symmetry with charges:

Superfields	R-charge
$\mathbf{H}_u$, $\mathbf{H}_d$	0
$\mathbf{R}_u$, $\mathbf{R}_d$	2
$\mathbf{H'_u}$, $\mathbf{H'_d}$	$+1, -1$

FSSM- II

Parametrising the breaking of this R-symmetry by a small parameter ε, the part of the superpotential containing the μ terms of the three Higgs-like multiplets is

$$
\begin{aligned}
W \supset\ & \varepsilon^2(\mu \mathbf{H}_u \mathbf{H}_d + \mu_{H'} \mathbf{H}'_u \mathbf{H}'_d) \\
& + [\mu_u \mathbf{H}_u \mathbf{R}_d + \mu_d \mathbf{R}_u \mathbf{H}_d] \\
& + \varepsilon \mu_{fdr} \mathbf{R}_u \mathbf{H}'_d + \varepsilon \mu_{df} \mathbf{H}'_u \mathbf{H}_d + \varepsilon^3 \mu_{uf} \mathbf{H}_u \mathbf{H}'_d.
\end{aligned}
$$

The R-charges have been chosen so that the mixing terms between $\mathbf{H}_{u,d}$ and $\mathbf{R}_{u,d}$ fields are unsuppressed. This allows the particles described mainly by $\mathbf{H}_{u,d}$ and $\mathbf{R}_{u,d}$ to have masses of order M_s, while $\mathbf{H}'_{u,d}$ provide a pair of light F-higgsinos with a mass of $\mathcal{O}(\varepsilon^2 M_s)$. The Yukawa part of the superpotential is given by

$$
\begin{aligned}
W \supset\ & [Y_u \mathbf{U^c Q} \cdot \mathbf{H_u} - Y_d \mathbf{D^c Q} \cdot \mathbf{H_d} - Y_e \mathbf{E^c L} \cdot \mathbf{H_d}] \\
& + \varepsilon[-Y_d \mathbf{D^c Q} \cdot \mathbf{H'_d} - Y_e \mathbf{E^c L} \cdot \mathbf{H'_d}]
\end{aligned}
$$

which allows a successful mass generation for the quarks and leptons, the SM-like Higgs obtained from fine-tuning at the electroweak scale must originate from the $\mathbf{H}_u$ and $\mathbf{H}_d$ multiplets.

FSSM- II

Imposing the R-symmetry on the soft terms leads to the suppression of the Majorana gauginos mass by ε^2 factors (this mechanism is similar to the usual Split SUSY one). In the term of the vector $v_H \equiv (H_u, H_d^*, H_u', H_d'^*, R_u, R_d^*)$, the Higgs mass matrix has the following hierarchy

$$
-\frac{1}{M_S^2} \mathcal{L}_{soft} \supset v_H^\dagger \begin{pmatrix}
\mathcal{O}(1) & \mathcal{O}(1) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon^2) & \mathcal{O}(\varepsilon^2) \\
\mathcal{O}(1) & \mathcal{O}(1) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon^2) & \mathcal{O}(\varepsilon^2) \\
\mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(1) & \mathcal{O}(1) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon^3) \\
\mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(1) & \mathcal{O}(1) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon^3) \\
\mathcal{O}(\varepsilon^2) & \mathcal{O}(\varepsilon^2) & \mathcal{O}(\varepsilon) & \mathcal{O}(\varepsilon) & \mathcal{O}(1) & \mathcal{O}(\varepsilon^4) \\
\mathcal{O}(\varepsilon^2) & \mathcal{O}(\varepsilon^2) & \mathcal{O}(\varepsilon^3) & \mathcal{O}(\varepsilon^3) & \mathcal{O}(\varepsilon^4) & \mathcal{O}(1)
\end{pmatrix} v_H.
$$

$$(14)$$

FSSM- II

We can tune the SM-like Higgs from the scalar components of $\mathbf{H}_u$ and $\mathbf{H}_d$ to get

$$H_u \approx \sin\beta\, H + \ldots\,, \qquad H_d \approx \cos\beta\, i\sigma^2\, H^* + \ldots\,,$$

and the other Higgs-like scalars only enters the linear combination with ε suppression. The fine-tuning condition can therefore be applied on the B_μ term similarly, with the exception that the μ terms are not ε-suppressed compared to the soft masses, leading to

$$B_\mu \simeq \sqrt{(m_{H_u}^2 + \mu_u^2)(m_{H_d}^2 + \mu_d^2)} + \mathcal{O}(\varepsilon)\,. \tag{15}$$

The parameter ε is here also fixed by the requirement that the gauginos obtain a mass at the TeV scale

$$\varepsilon = \mathcal{O}\left(\sqrt{\frac{\text{TeV}}{M_s}}\right). \tag{16}$$