BRANE-TO-BRANE GRAVITY MEDIATION OF SUPERSYMMETRY BREAKING

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- SUSY breaking and standard scenarios.
  Flavor problem and SUGRA models
  Weakly coupled and sequestered models
- Gravitational loop effects in sequestered models
  Soft scalar squared masses and their sign
  Generalizations of the minimal setup
- Dynamical radius stabilization
  Naturalness of magic value for the radius
  Concrete models and their viability

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SUSY BREAKING

Direct spontaneous supersymmetry breaking at the tree-level implies, in a renormalizable and anomaly-free theory, a sum rule on the mass spectrum:

$$\mathrm{STr}\,[\mathcal{M}^2] = \sum_J (-1)^{2J}(2J+1)\, m_J^2 = 0$$

This predicts rather generically that one of the superparticles is lighter than its ordinary partner, in contradiction with experimental observation.

The standard paradigm to evade this difficulty is to assume that SUSY breaking occurs spontaneously in a hidden sector and is transmitted to the visible sector only indirectly, through some interactions.

The effect of SUSY breaking on the visible sector can be parametrized through super-renormalizable soft breaking terms, which depend both on the details of the hidden sector theory and on the mediation mechanism.

The relevant effective Lagrangian to be considered has then the following general form:

$$\mathcal{L} = \mathcal{L}_{\text{susy}} + \mathcal{L}_{\text{soft}}$$

The supersymmetric part is schematically

$$\mathcal{L}_{\rm susy} = \Big[\Phi_i^\dagger e^{g_a V_a}\Phi_i\Big]_D + {\rm Re}\,\Big[\mathcal{W}_a^2\Big]_F + {\rm Re}\,\Big[\frac{1}{2}\,\mu_{ij}\,\Phi_i\Phi_j + \frac{1}{3}\,y_{ijk}\,\Phi_i\Phi_j\Phi_k\Big]_F$$

whereas the breaking part has correspondingly the general form

$$\mathcal{L}_{\rm soft} = -\,m_{0\,ij}^2\,\phi_i^*\phi_j - m_{1/2\,a}\,\bar\lambda_a\lambda_a - {\rm Re}\Big[\frac{1}{2}B_{ij}\,\mu_{ij}\,\phi_i\phi_j + \frac{1}{3}\,A_{ijk}\,y_{ijk}\,\phi_i\phi_j\phi_k\Big]$$

This general situation is however generically unacceptable both from the phenomenological and the theoretical point of view. In fact, the two main delicate points are the following:

- Flavor: m_0^2 must be positive and have a flavor structure that is nearly universal or aligned to that of y, in order to preserve the GIM mechanism.
- Hierarchy: μ, $\sqrt{B\mu}$, m_0 and $m_{1/2}$ must all be close to the electroweak scale $M_{\rm EW}$, in order to naturally explain electroweak symmetry breaking.

The main issue is then to find a microscopic setup providing these needed peculiarities in a natural and robust way. This has led to two main classes of models, which each have distinct appealing properties.

Gauge mediation

SUSY breaking occurs at a scale M in a hidden sector and is transmitted to the visible sector through heavy messenger fields with mass $M_{\rm M} \gg M$, direct couplings to the hidden sector and gauge couplings to the visible one.

Dine, Nelson, Shirman

The microscopic theory is a renormalizable theory. But below $M_{\rm M}$, and in particular at M, it is effectively described by a non-renormalizable SYM theory. Soft terms originate from higher-dimensional operators with scale $M_{\rm M}$.

The gaugino and squared scalar masses are generated at the 1-loop and 2-loop levels:

$$m_{1/2} \sim \frac{g^2}{16\pi^2}\frac{M^2}{M_{\rm M}}\,,\quad m_0^2 \sim \Big(\frac{g^2}{16\pi^2}\Big)^2\Big(\frac{M^2}{M_{\rm M}}\Big)^2$$

The soft masses have therefore a common scale given by

$$m_{\rm soft} \sim \frac{g^2}{16\pi^2}\frac{M^2}{M_{\rm M}}$$

The main characteristics are that:

- m_0^2 comes out automatically universal and turns out to be positive.
- It is difficult to make μ and $\sqrt{B\mu}$ both of the order $m_{\rm soft}$ at the same time.

Gravity mediation

SUSY breaking occurs at a scale M in the hidden sector and is transmitted to the visible sector through gravitational interactions, whose strength is set by $M_{\rm P} \gg M$.

Chamseddine, Arnowitt, Nath
Barbieri, Ferrara, Savoy
Hall, Lykken, Weinberg

The microscopic theory must be some fundamental theory. But below $M_{\rm P}$, and in particular at M, it is effectively described by a non-renormalizable SUGRA theory. Soft terms originate from higher-dimensional operators with scale $M_{\rm P}$.

The gaugino and squared scalar masses are both generated at the tree-level:

$$m_{1/2} \sim \frac{M^2}{M_{\rm P}}\,,\quad m_0^2 \sim \left(\frac{M^2}{M_{\rm P}}\right)^2$$

The soft masses have therefore a common scale given by

$$m_{\rm soft} \sim \frac{M^2}{M_{\rm P}}$$

The main characteristics are that:

- m_0^2 is in general neither universal nor positive.
- It is easy to make μ and $\sqrt{B\mu}$ both of the order $m_{\rm soft}$ at the same time.

Extra symmetries

To naturally solve both problems, one can try to introduce new symmetries. Main options:

- Gauge mediation + complications for hierarchy.
  No really simple and compelling model so far.
- Gravity mediation + constraints for flavor.
  Difficult to forbid mixing of the two sectors.

Extra dimensions

An interesting possibility, which is natural in string theory, is to separate the visible and the hidden sectors along an extra dimension. This framework has very peculiar characteristics going beyond symmetries:

- Geographic distinction between visible sector, hidden sector, and mediating interactions.
- New physical scale M_C acting as cut-off for the mixing between the two sectors.

SUSY FLAVOR PROBLEM

The flavor structure of Yukawa couplings should ideally be explained by some physical mechanism, associated to some physical scale $M_{\rm F}$, where the distinction between different species of fermions takes place.

If $M_{\rm F}$ is not larger than the scale at which soft terms are generated, these will in general acquire a non-trivial and independent flavor structure. This is the case of gravity mediation ($M_{\rm F} \lesssim M_{\rm P}$), in contrast to gauge mediation ($M_{\rm F} \gtrsim M_{\rm M}$).

There exist very severe experimental constraints on the relative size of non-universal versus universal scalar soft masses. The strongest bounds come from ϵ_K for squarks and from $\mathrm{Br}(\mu \to e + \gamma)$ for sleptons. The present limits are

$$\left(\frac{\delta m^2_{\rm u\!\!\!/niv}}{m^2_{\rm univ}}\right)_{\rm squarks} \lesssim 3 \times 10^{-4} \left(\frac{m_{\rm univ}}{M_{\rm EW}}\right)$$

$$\left(\frac{\delta m^2_{\rm u\!\!\!/niv}}{m^2_{\rm univ}}\right)_{\rm sleptons} \lesssim 3 \times 10^{-3} \left(\frac{m_{\rm univ}}{M_{\rm EW}}\right)^2$$

The accuracy of the relevant experimental measurements will further improve in the next few years, and we clearly need at least a partial theoretical explanation for these small values.

SUGRA MODELS

Consider a general SUGRA theory with:

Visible: $\Phi_0 = (\phi_0, \chi_0; F_0)$, $V_0 = (A_0^\mu, \lambda_0; D_0)$.

Hidden: $\Phi_\pi = (\phi_\pi, \chi_\pi; F_\pi)$, $V_\pi = (A_\pi^\mu, \lambda_\pi; D_\pi)$.

Interactions: $C = (e_\mu^a, \psi_\mu; a_\mu, b_\mu)$, $S = (s, \psi_S; F_S)$.

After SUPCONF gauge-fixing, $b_\mu = 0$, $s = 1$, $\psi_S = 0$, and:

$$\mathcal{L} = \left[\Omega(\Phi, \Phi^\dagger e^V) S S^\dagger\right]_D + \mathrm{Re}\left[P(\Phi) S^3\right]_F + \mathrm{Re}\left[\tau(\Phi)\mathcal{W}^2\right]_F$$

The functions Ω, $\tau\mathcal{W}^2$ and P have expansions of the type:

$$\Omega = -3M_\mathrm{P}^2 + \Phi_0^\dagger e^{g_0 V_0}\Phi_0 + \Phi_\pi^\dagger e^{g_\pi V_\pi}\Phi_\pi + \frac{h}{M_\mathrm{P}^2}\Phi_0^\dagger e^{g_0 V_0}\Phi_0\, \Phi_\pi^\dagger e^{g_\pi V_\pi}\Phi_\pi + \dots$$

$$P = \frac{1}{16\pi^2}\Lambda_\pi^3 + M_\pi^2 \Phi_\pi + \dots$$

$$\tau\mathcal{W}^2 = \mathcal{W}_0^2 + \mathcal{W}_\pi^2 + \frac{k}{M_\mathrm{P}}\Phi_\pi \mathcal{W}_0^2 + \dots$$

We assume that SUSY breaking originates only from a singlet Φ_π, and none of the V_π. To have a vanishing CC, we need to tune $\Lambda_\pi^3 \sim 16\pi^2 M_\pi^2 M_\mathrm{P}$. The SUSY breaking VEVs are then:

$$|F_\pi| \sim M_\pi^2\,, \quad |F_S| \sim \frac{\Lambda_\pi^3}{16\pi^2 M_\mathrm{P}^2} \sim \frac{M_\pi^2}{M_\mathrm{P}}$$

Therefore, there are two correlated sources of SUSY breaking:

$$|F_S| \sim \frac{|F_\pi|}{M_{\rm P}} \sim m_{3/2}$$

with the scale

$$m_{3/2} \sim \frac{M_\pi^2}{M_{\rm P}}$$

Classical theory

The leading soft masses at classical level are:

$$m_{1/2} \sim k\,\frac{|F_\pi|}{M_{\rm P}} \sim k\,m_{3/2}$$

$$m_0^2 \sim h\,\frac{|F_\pi|^2}{M_{\rm P}^2} \sim h\,m_{3/2}^2$$

The coefficients h and k are model-dependent. The former is generically expected to be of order unity and non-universal, since the fundamental theory is supposed to explain flavor.

Quantum corrections

The corrections coming from gravity loops give

$$m_{1/2} \sim 0$$

$$m_0^2 \sim \frac{\Lambda^2}{16\pi^2 M_{\rm P}^2}\,\frac{|F_\pi|^2}{M_{\rm P}^2} \sim \frac{\Lambda^2}{16\pi^2 M_{\rm P}^2}\,m_{3/2}^2$$

They are not reliably calculable, but their size can be estimated by using a cut-off scale Λ that depends on what is assumed for the fundamental theory.

Assuming nothing special corresponds to taking the highest possible range of validity of the effective theory, with a cut-off given by $\Lambda \sim 4\pi M_{\mathrm{P}}$.

In that case, the loop corrections to $m_{1/2}$ are negligible, due to R-symmetry, but those to m_0 are generically large. This means that it makes sense to choose k, but not h.

There are also corrections coming from gauge loops. These induced an anomalous logarithmic dependence on S through the superconformal anomaly, which yields:

$$m_{1/2} \sim \frac{g^2}{16\pi^2} |F_S| \sim \frac{g^2}{16\pi^2} m_{3/2}$$

$$m_0^2 \sim \left(\frac{g^2}{16\pi^2}\right)^2 |F_S|^2 \sim \left(\frac{g^2}{16\pi^2}\right)^2 m_{3/2}^2$$

These are reliably calculable, model-independent and smaller than the tree-level effects. They are universal and have a definite sign: positive for squarks and negative for sleptons.

Randall, Sundrum

Giudice, Luty, Murayama, Rattazzi

To reach a satisfactory situations for the soft masses, one needs therefore to introduce some selection mechanism going beyond symmetries and suppressing non-universal effects.

This calls for some specific assumption about the fundamental theory. These should however concern only its general structure, not its details, to represent a robust situation.

SPECIFIC SCENARIOS

Weakly coupled models

A first possibility is to assume weak coupling. Loops are then cut off at $\Lambda \sim M_{\rm S} \ll 4\pi M_{\rm P}$ and suppressed by the factor

$$\frac{M_{\rm S}^2}{16\pi^2 M_{\rm P}^2} \ll 1$$

In this situation, it is technically possible to assume that classical effects are for some reason universal and that non-universal loop corrections are small.

The requirement that the underlying theory should provide a unified description of interactions however naturally suggests that (gravity loop) $\sim$ (gauge loop) at the scale $M_{\rm S}$, that is:

$$\frac{M_{\rm S}^2}{16\pi^2 M_{\rm P}^2} \sim \frac{g^2}{16\pi^2} \;\Rightarrow\; \frac{M_{\rm S}}{M_{\rm P}} \sim g \sim 7 \times 10^{-1}$$

This applies, for instance, to standard superstring models.

Louis, Nir

Brignole, Ibañez, Muñoz

In this setup, the leading contributions to m_0 and $m_{1/2}$ come thus from tree-level gravitational effects that are assumed to be universal and have the standard scale:

$$m_{\rm univ} \sim m_{3/2}$$

The leading non-universal corrections come from gravity loops and have a relative suppression factor given by:

$$\frac{\delta m^2_{\not{\rm univ}}}{m^2_{\rm univ}} \sim \frac{g^2}{16\pi^2} \sim 3 \times 10^{-3}$$

The main issue in this scenario is the possibility of having at the same time non-universal Yukawa couplings and universal scalar soft masses at the classical level. This seems to be possible in certain string models, thanks to special superselection rules.

Sequestered models

A second possibility is to assume the existence of an extra dimension, along which the visible and hidden sectors are separated and interact only gravitationally. The relevant loops are then cut off at $\Lambda \sim M_{\rm C} \ll 4\pi M_{\rm P}$ and suppressed by

$$\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2} \ll 1$$

In this situation, higher-dimensional locality forbids any local classical effect and allows only non-local quantum effects, which are generated at the scale $M_{\rm C}$ and are thus universal and reliably computable within a 5D SUGRA description.

It is then conceivable that all the soft masses are induced by universal loop effects.

To get a satisfactory situation, the gravity loop contribution to m_0^2 should be positive, and large enough to overcome the negative gauge loop contribution for sleptons. This requires that (gravity loop) $\sim$ (gauge loop)^2 at M_C, that is:

$$\frac{M_\mathrm{C}^2}{16\pi^2 M_\mathrm{P}^2} \sim \left(\frac{g^2}{16\pi^2}\right)^2 \Rightarrow \frac{M_\mathrm{C}}{M_\mathrm{P}} \sim \frac{g^2}{4\pi} \sim 4\times 10^{-2}$$

This is reasonable and interesting.

Randall, Sundrum

In this setup, the leading contributions to m_0 and $m_{1/2}$ come thus from tuned 1-loop gravity effects and 1-loop gauge effects that are automatically universal, with a common scale:

$$m_\mathrm{univ} \sim \frac{M_\mathrm{C}}{4\pi M_\mathrm{P}} m_{3/2} \sim \frac{g^2}{16\pi^2} m_{3/2}$$

Generic non-universal corrections come from gravity loops involving couplings with two extra derivatives. This implies two extra inverse-powers of the bulk gravity scale $(4\pi M_\mathrm{P})^{2/3} M_\mathrm{C}^{1/3}$. By dimensional analysis, these effects are thus suppressed by

$$\frac{\delta m_{\not\mathrm{univ}}^2}{m_\mathrm{univ}^2} \sim \left(\frac{M_\mathrm{C}}{4\pi M_\mathrm{P}}\right)^{4/3} \sim \left(\frac{g^2}{16\pi^2}\right)^{4/3} \sim 5\times 10^{-4}$$

The main issues in this scenario are the possibility to get the right sign and size for the gravitational effect and to keep possible extra model-dependent non-universal effects related to radius stabilization small enough.

MINIMAL SEQUESTERED MODELS

The simplest example of sequestered models can be constructed by compactifying a 5D SUGRA theory on the orbifold $S^1/\mathbf{Z_2}$, with a flat background geometry, and locating the visible and the hidden sectors at the two fixed-point or branes.

The geometry is parametrized by the radius R of the internal dimension and the 5D Lagrangian has the general form:

$$\mathcal{L}(x,y) = \mathcal{L}_5(x,y) + \mathcal{L}_0(x)\,\delta(y - 0R) + \mathcal{L}_\pi(x)\,\delta(y - \pi R)$$

where

$$e_5^{-1}\mathcal{L}_5 = -\frac{1}{2}M_5^3\mathcal{R}_5 + \cdots$$

$$e_4^{-1}\mathcal{L}_{0,\pi} = -\frac{1}{2}M_5^3 L_{0,\pi}\mathcal{R}_4 + |\partial^\mu\phi_{0,\pi}|^2 - \frac{1}{4g_{0,\pi}^2}(F_{0,\pi}^{\mu\nu})^2 + \cdots$$

The bulk fields can be expanded in towers of KK modes with masses proportional to

$$M_\mathrm{C} = \frac{1}{\pi R}$$

The low-energy effective theory for $E \ll M_\mathrm{C}$, which is obtained by integrating out the heavy KK modes, is a 4D SUGRA theory, with an extra chiral multiplet:

Radion: $T = (t, \psi_T; F_T)$.

This new multiplet describes the dynamics associated to the extra dimension. The VEV of t controls the size of the extra dimensions, $\mathrm{Re}\, t = \pi R$, whereas the VEV for F_T can in general give additional SUSY-breaking effects.

The functions Ω, $\tau\mathcal{W}^2$ and P have now the form

$$\Omega = -3M_5^3\left(T+T^\dagger+L_0+L_\pi\right) + \Phi_0^\dagger e^{g_0 V_0}\Phi_0 + \Phi_\pi^\dagger e^{g_\pi V_\pi}\Phi_\pi + \ldots$$

$$P = \frac{1}{16\pi^2}\Lambda_\pi^3 + M_\pi^2\Phi_\pi + f(T) + \ldots$$

$$\tau\mathcal{W}^2 = \mathcal{W}_0^2 + \mathcal{W}_\pi^2 + \ldots$$

From the matter-independent terms we deduce that:

$$M_{\mathrm{P}}^2 = M_5^3\left(2\,\mathrm{Re}\, t + L_0 + L_\pi\right)$$

We assume as before that SUSY breaking comes only from a singlet Φ_π and no V_π, and also that the radion potential is not significantly influenced by gravitational effects.

A vanishing CC again requires the tuning $\Lambda_\pi^3 \sim 16\pi^2 M_\pi^2 M_{\mathrm{P}}$. The SUSY breaking VEVs are then

$$|F_\pi| \sim M_\pi^2\ , \quad |F_S| \sim \frac{\Lambda_\pi^3}{16\pi^2 M_{\mathrm{P}}^2} \sim \frac{M_\pi^2}{M_{\mathrm{P}}}$$

and, generically assuming that $f'(t)/f''(t) \sim t$,

$$|F_T| \sim \left|\frac{f'(t)}{f''(t)}\right| \frac{\Lambda_\pi^3}{16\pi^2 M_{\mathrm{P}}^2} \sim |t|\, \frac{M_\pi^2}{M_{\mathrm{P}}}$$

In this situation, there are therefore three correlated sources of SUSY breaking

$$|F_S| \sim \frac{|F_\pi|}{M_\mathrm{P}} \sim \frac{|F_T|}{|t|} \sim m_{3/2}$$

with the scale

$$m_{3/2} \sim \frac{M_\pi^2}{M_\mathrm{P}}$$

There exist no local operators that can mixing the visible and hidden sectors at the classical level, and the only effect comes from non-local operators generated at the quantum level.

The dominant effects come from the 1-loop correction to Ω. This has a divergent T-independent part, renormalizing the operators that were already present, and a finite T-dependent part, containing new operators.

The effect under consideration is similar to a Casimir effect. The relevant bulk-brane interactions are the field-dependent localized kinetic terms for the bulk fields, which are determined by the brane Kähler functions:

$$\Omega_{0,\pi} = -3L_{0,\pi}M_5^3 + \Phi_{0,\pi}^\dagger e^{g_{0,\pi}V_{0,\pi}}\Phi_{0,\pi}$$

By power counting and dimensional analysis, the correction must then have the form:

$$\Delta\Omega(\Omega_{0,\pi}, T+T^\dagger) = \sum_{m,n=0}^{\infty} \frac{c_{m,n}\,\Omega_0^m\,\Omega_\pi^n}{M_5^{3(m+n)}(T+T^\dagger)^{2+m+n}}$$

By explicit computation, one finds:

$$\Delta\Omega(\Omega_{0,\pi}, T+T^\dagger)$$
$$= -\frac{9}{\pi^2}M_5^2\int_0^\infty dx\, x\ln\left[1-\frac{1+\frac{\Omega_0}{M_5^2}x}{1-\frac{\Omega_0}{M_5^2}x}\frac{1+\frac{\Omega_\pi}{M_5^2}x}{1-\frac{\Omega_\pi}{M_5^2}x}e^{-6(T+T^\dagger)M_5 x}\right]$$

Rattazzi, Scrucca, Strumia

Expanding this expression in powers of $\Omega_{0,\pi}/M_5^2$, one deduces that all the $c_{m,n}$ are positive. The first few ones are:

$$c_{0,0} = \frac{\zeta(3)}{4\pi^2}\,,\quad c_{1,0} = c_{0,1} = \frac{\zeta(3)}{6\pi^2}\,,\quad c_{1,1} = \frac{\zeta(3)}{6\pi^2}\,,\quad \ldots$$

The first three coefficients have been independently derived also in other ways.

Buchbinder et al.
Gherghetta, Riotto
Ponton, Poppitz
Falkowski

The relative sign of the first three terms can be checked by noticing that the leading-order dependence on $\Omega_{0,\pi}$ of both the kinetic terms and the boundary conditions of the bulk modes can be obtained through the replacement:

$$T+T^\dagger \to T+T^\dagger - \frac{\Omega_0}{3M_5^3} - \frac{\Omega_\pi}{3M_5^3}$$

This is related to the fact that a displacement of the branes can be accounted for by a flow of the local operators that they support.

In the general case with non-vanishing localized kinetic terms, $L_0, L_\pi \neq 0$, all the operators with coefficients $c_{m,n}$ give a non-vanishing contribution to m_0^2. The result depends on the dimensionless variables:

$$r_{0,\pi} = \frac{L_{0,\pi}}{2\pi R}$$

One can vary the parameters $r_{0,\pi}$, but one must correspondingly change the bulk gravity scale $M_5(r_{0,\pi})$, to keep a fixed value for the 4D effective Planck scale:

$$M_{\rm P} \sim M_5^{3/2}\, M_{\rm C}^{-1/2}(1+r_0+r_\pi)^{1/2}$$

In terms of the constant $M_{\rm P}$ and the parameters $r_{0,\pi}$, the induced m_0^2 is the sum of RB and BB effects:

$$\begin{aligned} m_0^2 \;\sim\;\; & f_{\rm RB}(r_0, r_\pi)\,(1+r_0+r_\pi)\,\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2}\,\frac{|F_T|^2}{|t|^2} \\ & +f_{\rm BB}(r_0, r_\pi)\,(1+r_0+r_\pi)^2\,\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2}\,\frac{|F_\pi|^2}{M_{\rm P}^2} \\ \sim\;\; & c(r_0, r_\pi)\,\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2}\,m_{3/2}^2 \end{aligned}$$

The functions $f_{\rm RB}$ and $f_{\rm BB}$ are known and reflect the dependence of the masses and the wave functions of the bulk KK modes on the parameters $r_{0,\pi}$. There are then three substantially different regimes.

Regime where $r_0 \ll 1$ and $r_\pi \ll 1$

The KK modes have an integer spectrum and symmetric delocalized wave-functions. The RB and BB effects are both negative and sizable:

$$f_{\rm RB} \sim -1\,, \quad f_{\rm BB} \sim -1 \;\Rightarrow\; c \sim -1$$

Regime where $r_0 \gg 1$ and $r_\pi \ll 1$

The KK have a half-integer spectrum and wave-functions that are localized at the hidden sector. The RB effect is negative and suppressed, whereas the BB effect is negative and sizable:

$$f_{\rm RB} \sim -\frac{1}{r_0^2}\,, \quad f_{\rm BB} \sim -\frac{1}{r_0^2} \;\Rightarrow\; c \sim -1$$

Regime where $r_0 \ll 1$ and $r_\pi \gg 1$

The KK have a half-integer spectrum and wave-functions that are localized at the visible sector. The RB effect is positive and enhanced, whereas the BB effect is negative and sizable:

$$f_{\rm RB} \sim +1\,, \quad f_{\rm BB} \sim -\frac{1}{r_\pi^2} \;\Rightarrow\; c \sim +r_\pi$$

Summarizing, in this minimal setup the sign of the induced m_0^2 tends to be negative, and can become positive only if the bulk modes are significantly more localized on the visible sector than on the hidden sector.

Rattazzi, Scrucca, Strumia

GENERALIZATIONS

Warped models

A natural generalization of the minimal setup can be obtained by switching on a non-trivial curvature and considering warped geometries with an arbitrary warping parameter k.

It is known that warped models share many specific features with flat models in the presence of asymmetric distributions of kinetic terms, because of the similarities of the spectra and the wave functions of their KK modes.

It is then a priori conceivable that a sizable warping could lead, as localized kinetic terms, to a positive sign for the induced m_0^2, and provide a more satisfactory justification for the underlying asymmetry between visible and hidden sectors.

However, a detailed analysis of the general case with non-vanishing warping and arbitrary localized kinetic terms has shown that warping cannot give on its own a positive effect, and localized kinetic terms are still necessary.

These models, though very interesting, do therefore not offer really new possibilities.

Gregoire, Rattazzi, Scrucca, Strumia, Trincherini

Models with both F and D type SUSY breaking

An other possible generalization is to consider models where SUSY breaking occurs not only through a chiral auxiliary field F_π but also a vector auxiliary field D_π.

This is generically the case in models where the fields Φ_π are not all singlets under the fields V_π. One needs however a more complicated superpotential P_π.

The basic point is that F_π and D_π can lead to effects that differ in sign. This is easily seen from the way the various superspace structures contribute to the action. One has:

$$\begin{aligned}\mathcal{L} &\supset \Big[\Phi_\pi^\dagger e^{g_\pi V_\pi}\Phi_\pi\Big]_D + \mathrm{Re}\Big[\mathcal{W}_\pi^2\Big]_F + \mathrm{Re}\Big[P_\pi(\Phi_\pi)\Big]_F \\ &\supset \Big[|F_\pi|^2 + D_\pi g_\pi \Phi_\pi^\dagger T \Phi_\pi\Big] + \Big[\frac{1}{2}D_\pi^2\Big] + \Big[F_\pi P_\pi'(\Phi_\pi) + \mathrm{h.c.}\Big]\end{aligned}$$

The auxiliary field equations of motion are then:

$$F_\pi = -P_\pi'(\Phi_\pi)\ ,\quad D_\pi = -g_\pi \Phi_\pi^\dagger T \Phi_\pi$$

Substituting back gives:

$$\mathcal{L} \supset \Big[|F_\pi|^2 - D_\pi^2\Big] + \Big[\frac{1}{2}D_\pi^2\Big] + \Big[-2|F_\pi|^2\Big]$$

Thus, F_π and D_π contribute differently to $[\cdots]_D$ and $[\cdots]_F$ terms, although they enter similarly in the total potential:

$$\mathcal{L} \supset -\Big(|F_\pi|^2 + \frac{1}{2}D_\pi^2\Big)$$

Quantum effects are of $[\cdots]_D$ rather than $[\cdots]_F$ type, and are more sensitive to $[\cdots]_D$ than $[\cdots]_F$ classical couplings, due R-symmetry. The effects due to F_π and D_π can then differ.

This applies in particular to gravitational corrections that are relevant for sequestered models. Their form is unchanged:

$$\Delta\Omega(\Omega_{0,\pi}, T+T^\dagger) = \sum_{m,n=0}^{\infty} \frac{c_{m,n}\,\Omega_0^m\,\Omega_\pi^n}{M_5^{3(m+n)}(T+T^\dagger)^{2+m+n}}$$

The m_0^2 that they induce is however different. The RB effect still comes with $|F_T|^2$, but the BB effect now involves the combination $|F_\pi|^2 - D_\pi^2$, which is no longer positive.

It is quite natural to expect that $|F_\pi|^2 \sim D_\pi^2$ and that the sign of the BB effect is not an issue in general. We have verified this by working out various examples. Tuning the CC to zero, one finds then four correlated VEVs:

$$|F_S| \sim \frac{|F_\pi|}{M_{\rm P}} \sim \frac{|D_\pi|}{M_{\rm P}} \sim \frac{|F_T|}{|t|} \sim m_{3/2}$$

This leads finally to

$$\begin{aligned} m_0^2 \sim\; & f_{\rm RB}(r_0,r_\pi)\,(1+r_0+r_\pi)\,\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2}\,\frac{|F_T|^2}{|t|^2} \\ & + f_{\rm BB}(r_0,r_\pi)\,(1+r_0+r_\pi)^2\,\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2}\,\frac{|F_\pi|^2 - D_\pi^2}{M_{\rm P}^2} \\ \sim\; & d(r_0,r_\pi)\,\frac{M_{\rm C}^2}{16\pi^2 M_{\rm P}^2}\,m_{3/2}^2 \end{aligned}$$

Gregoire, Rattazzi, Scrucca

RADIUS STABILIZATION

There exists a natural way to stabilize the radius at the magic value for which gravity-mediated effects are of the same order of magnitude as anomaly-mediated effects in soft terms.

Luty, Sundrum

It relies on an additional bulk vector multiplet, for which gaugino condensation takes place. To preserve sequestering, the boundary fields must be singlets under the new symmetry.

The additional 5D Lagrangian has the form

$$\mathcal{L}(x,y) = \mathcal{L}_5(x,y) + \mathcal{L}_0(x)\,\delta(y - 0R) + \mathcal{L}_\pi(x)\,\delta(y - \pi R)$$

where now

$$e_5^{-1}\mathcal{L}_5 = -\frac{1}{4g_5^2}\,(F^{\mu\nu})^2 + \cdots$$

$$e_4^{-1}\mathcal{L}_{0,\pi} = -\frac{1}{4g_5^2}\,K_{0,\pi}\,(F^{\mu\nu})^2 + \cdots$$

The NDA estimate for the 5D cut-off scale is:

$$\Lambda_{\mathrm{G}} \sim 16\pi^2 g_5^{-2}$$

The classical 4D effective gauge coupling is given by

$$g_{\mathrm{G}}^{-2} = g_5^{-2}\Big(2\,\mathrm{Re}\,t + K_0 + K_\pi\Big)$$

Quantum effects induce a logarithmic running of the 4D effective coupling, which becomes strong at the scale $\Lambda_{\rm G}\, e^{-16\pi^2 g_{\rm G}^{-2}}$. Gaugino condensation induces then the superpotential

$$f(T) \sim \frac{1}{16\pi^2}\Lambda_{\rm G}^3\, e^{-\Lambda_{\rm G}(2T+K_0+K_\pi)}$$

The full theory has the following scales, after the CC tuning: the overall fixed $M_{\rm P}$, the arbitrary $m_{3/2}$ from the SUSY breaking sector, and the arbitrary $\Lambda_{\rm G}$ and $K_{0,\pi}$ from the radion stabilizing sector.

The crucial dimensionless parameter for the minimization of the potential turns out to be:

$$\epsilon \sim \left(\ln\frac{\Lambda_{\rm G}^3}{16\pi^2 M_{\rm P}^2\, m_{3/2}} - \frac{1}{2}K_0\Lambda_{\rm G} - \frac{1}{2}K_\pi\Lambda_{\rm G}\right)^{-1}$$

Under the assumption that $0 \lesssim \epsilon \lesssim 1$, one finds that the radion is stabilized with:

$$M_{\rm C} \sim \epsilon\,\Lambda_{\rm G}\;, \quad m_T \sim \epsilon^{-1}\, m_{3/2}$$

and the SUSY-breaking VEVs are given by

$$|F_S| \sim \frac{|F_\pi|}{M_{\rm P}} \sim m_{3/2}\;, \quad \frac{|F_T|}{|t|} \sim \epsilon\, m_{3/2}$$

At the minimum of the potential, ϵ can be identified with the loop factor for the new bulk gauge interactions at the compactification scale:

$$\epsilon \sim \frac{M_{\rm C}}{\Lambda_{\rm G}} \sim \frac{g_5^2 M_{\rm C}}{16\pi^2}$$

The value of $m_{3/2}$ is fixed by requiring the anomaly-mediated $m_{1/2}$ to be around $M_{\rm EW}$. This means that $m_{3/2}$ should not be much beyond $M_{\rm EW}$.

The value of $\Lambda_{\rm G}$ is practically fixed by requiring that $M_{\rm C}$ has the magic value needed for the gravity-mediated m_0 to also be around $M_{\rm EW}$. This implies that $\Lambda_{\rm G}$ should be close to $M_{\rm P}$.

The values of $K_{0,\pi}$ are instead free and allow to change ϵ.

This setup therefore realizes in a natural way the required situation, independently of the distribution of kinetic terms, which is controlled by the two dimensionless quantities

$$s_{0,\pi} = \frac{K_{0,\pi}}{2\pi R}$$

The stationary values of these can be used to parametrize the quantitative features of the model. In particular, using the fact that the $\ln$ term in ϵ is numerical about $(g^2/4\pi)^{-1}$, one can rewrite the latter as:

$$\epsilon \sim \Big(1 + s_0 + s_\pi\Big)\,\frac{g^2}{4\pi}$$

The bulk fields can mediate non-universal loop effects, due to higher-derivative couplings that are generically expected. These new effects can be dangerous. They have a RB part that can be ignored and a BB part that must be studied.

POSSIBLE MODELS

F-type breaking with a large localized kinetic term

In the universal gravity-mediated effect, RB $\propto r_\pi \cdot \epsilon^2$ must be larger than BB $\propto 1 \cdot 1 \Rightarrow$ assume $r_\pi \sim \epsilon^{-2}$.

The non-universal gauge-mediated effect has an extra ϵ^2 suppression but is $\propto (r_\pi/s_\pi)^2 \cdot 1 \Rightarrow$ assume $s_\pi \sim (g^2/4\pi)^{-1}$ so that $\epsilon \sim 1$ and $r_\pi \sim 1$. We get then:

$$\frac{\delta m^2_{\not{\mathrm{univ}}}}{m^2_{\mathrm{univ}}} \sim s_\pi^{-2} \sim \left(\frac{g^2}{4\pi}\right)^2 \gtrsim 2 \times 10^{-3}$$

D-type breaking with small localized kinetic term

In the universal gravity-mediated effect, RB $\propto 1 \cdot \epsilon^2$ must be smaller than BB $\propto 1 \cdot 1 \Rightarrow$ always fine.

The non-universal gauge-mediated effect has as before an extra ϵ^2 suppression and is now $\propto 1 \cdot 1 \Rightarrow$ assume $s_\pi \sim 0$ so that $\epsilon \sim g^2/4\pi$. We get then:

$$\frac{\delta m^2_{\not{\mathrm{univ}}}}{m^2_{\mathrm{univ}}} \sim \epsilon^2 \sim \left(\frac{g^2}{4\pi}\right)^2 \gtrsim 2 \times 10^{-3}$$

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CONCLUSIONS

- Sequestered models can provide a robust explanation for the flavor problem, the crucial issue being the dynamics that is supposed to stabilize the radius at its magic value.
- The leading non-universal effects are in general controlled by model-dependent physics stabilizing the radius, and viable concrete examples have been shown to exist.
- These scenarios represent a valid alternative to weakly coupled models, where the crucial issue is the selection rule that is supposed to be responsible for the universality of soft terms at the classical level.