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Supersymmetric unification without low energy supersymmetry and signatures for fine-tuning at the LHC

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Abstract: The cosmological constant problem is a failure of naturalness and suggests that a fine-tuning mechanism is at work, which may also address the hierarchy problem. An example — supported by Weinberg's successful prediction of the cosmological constant — is the potentially vast landscape of vacua in string theory, where the existence of galaxies and atoms is promoted to a vacuum selection criterion. Then, low energy SUSY becomes unnecessary, and supersymmetry — if present in the fundamental theory — can be broken near the unification scale. All the scalars of the supersymmetric standard model become ultraheavy, except for a single finely tuned Higgs. Yet, the fermions of the supersymmetric standard model can remain light, protected by chiral symmetry, and account for the successful unification of gauge couplings. This framework removes all the difficulties of the SSM: the absence of a light Higgs and sparticles, dimension five proton decay, SUSY flavor and CP problems, and the cosmological gravitino and moduli problems. High-scale SUSY breaking raises the mass of the light Higgs to $\sim 120-150$ GeV. The gluino is strikingly long lived, and a measurement of its lifetime can determine the ultraheavy scalar mass scale. Measuring the four Yukawa couplings of the Higgs to the gauginos and higgsinos precisely tests for high-scale SUSY. These ideas, if confirmed, will demonstrate that supersymmetry is present but irrelevant for the hierarchy problem — just as it has been irrelevant for the cosmological constant problem — strongly suggesting the existence of a fine-tuning mechanism in nature.

Keywords: Superstring Vacua, Supersymmetry Breaking.

Contents

1. Naturalness and its discontents

1.1 Naturalness

The Standard Model of particle physics is our most successful physical theory, providing an excellent description of experiments up to energies of order $\sim 100\,\mathrm{GeV}$. It is also a consistent theoretical structure that can be extrapolated by itself up to energies Λ_{SM} far above the weak scale. Yet, ever since the mid 1970's, there has been a widely held expectation that the SM must be incomplete already at the $\sim$TeV scale. The reason is the principle of naturalness: if Λ_{SM} is too large, the Higgs mass must be fine-tuned to an accuracy of order $(m_W/\Lambda_{SM})^2$ to explain the weak scale. Solving the naturalness problem has provided the biggest impetus to contructing theories of physics beyond the Standard Model, leading to the proposal of technicolor [1] and the supersymmetric standard model [2], and more recently, the idea of extra dimensions with low-scale quantum gravity [3, 4] and the little Higgs mechanism [5].

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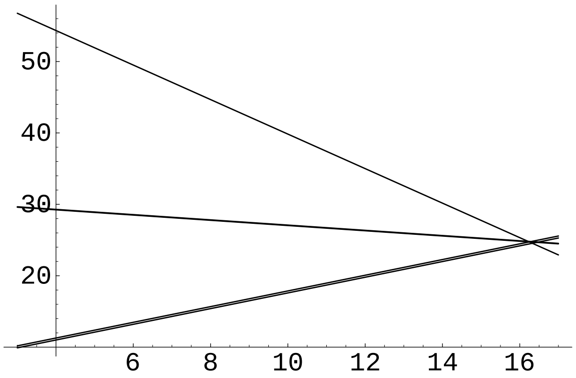

Figure 1: Running couplings in the SSM at one-loop, matching to the full SUSY running at m_t; from top to bottom, $\alpha_1^{-1}, \alpha_2^{-1}, \alpha_3^{-1}$. We use $\alpha_1^{-1}(M_Z) = 58.98 \pm .04$, $\alpha_2^{-1}(M_Z) = 29.57 \pm .03$, and $\alpha_3^{-1}(M_Z) = 8.40 \pm .14$.

Taking naturalness as a principle seriously has had one impressive concrete success, within the minimal supersymmetric standard model (SSM): the prediction of gauge coupling unification [2, 6, 7] at a scale $M_G \sim 2 \times 10^{16}\,\mathrm{GeV}$ tantalizingly near the Planck scale [6]. Another success is that many supersymmetric theories find a good dark matter candidate in the lightest neutralino [2, 8].

Despite these successes, the supersymmetric standard model has also had a number of difficulties, mostly having to do with the fact that the SM explanation for the conservation of baryon, lepton number and the absence of FCNC's as a consequence of accidental symmetries disappears in SUSY. There are the well-known dimension four R-parity violating couplings in the superpotential that give rise to large proton decay rates and neutrino masses. Imposing matter or R-parity to forbid these couplings is quite natural, though, and further ensures the stability of the LSP, making it a good dark matter candidate. However, there are a litany of other well-known problems that can not be dispensed with so elegantly. There are dimension five operators of the form $qq\tilde{q}\tilde{l}$, that give proton decay. There are new flavor violations in the the dimension four couplings of fermions to gauginos and sfermions, that give rise to the SUSY flavor problem. New CP violating phases have to be significantly suppressed to avoid large electron and neutron electric dipole moments. There are also corrections to quantities that do not violate symmetries any more than in the SM, but which receive significant contributions from superpartner loops, ranging from $(g-2)_\mu$ to $B - \bar{B}$ mixing and $b \to s\gamma$. Most important, the SSM strongly favors a light Higgs, as well as some light sparticles; their absence is troubling and indicates that there is already some tuning at the few percent level. Finally, the new gravitational particles in supersymmetric theories, the gravitino and moduli, are associated with a variety of cosmological difficulties.

1.2 A Failure of naturalness

Of course none of these challenges are insurmountable, and indeed attacking them has defined the program of supersymmetric model-building for the last twenty years. Leaving the basic structure of the SSM unaltered, various mechanisms have been invented to address these problems.

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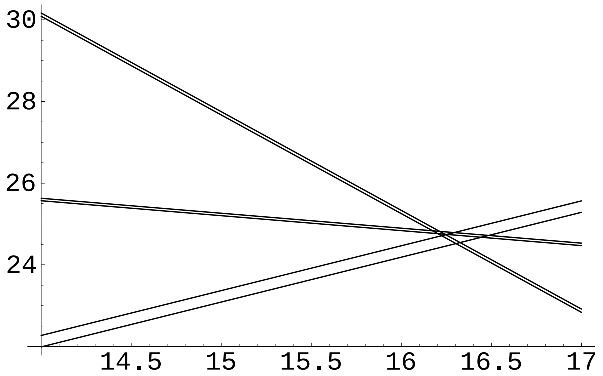

Figure 2: Close-up of the one-loop SSM running couplings near the unification scale. Two-loop contributions tend to push up the predicted value of $\alpha_3(M_Z)$ to about .130, away from the central value .119, requiring somewhat large compensating threshold corrections.

In this paper, we will instead suggest a simple but drastic modification of the usual supersymmetric picture of the world, which will in a single stroke remove all the phenomenological difficulties while automatically preserving the concrete successes of the SSM. In order to motivate our proposal, let us recall the usual logic leading to the prediction of weak-scale SUSY. Nature may well be supersymmetric at short distances, perhaps because SUSY is required for a consistent theory of quantum gravity. However, given that the low-energy theory does not exhibit bose-fermi degeneracy, SUSY must be broken. Let the scale of SUSY breaking in the SSM be m_S; the low-energy theory beneath m_S is non-supersymmetric, and therefore the Higgs mass parameter in this low-energy theory is UV sensitive. Having $m_h \ll m_S$ would require a fine-tuning, and it would be absurd for the world to be both supersymmetric *and* finely tuned! We therefore expect that

$$m_h^2 \sim m_S^2 \tag{1.1}$$

and so $m_S \lesssim 1\,\mathrm{TeV}$.

But there is cause to be suspicious of this logic: all of the theories we study with weak-scale SUSY *are* both supersymmetric and finely tuned, with an enormous fine-tuning for the cosmological constant. The same line of argument as above would predict

$$\Lambda \gtrsim m_S^4 \tag{1.2}$$

which is at least 60 orders of magnitude too large.

The usual attitude to the Cosmological Constant Problem has been one of abhorrence to this fine-tuning, hoping for some deep or exotic mechanism to explain either why the CC appears so small or why an enormous vacuum energy doesn't gravitate. Perhaps the CC is small because of the UV/IR connection, holography and the mysteries of gravity in deSitter space [9], perhaps the graviton is composite at the millimeter scale [10], or maybe gravity is modified in the IR in a way that prevents the large vacuum energy from giving rise to an unacceptable large expansion rate for the universe [11, 12].

Whatever mechanism may be at work, the fact is that in concrete theories, the vacuum energy is cancelled by a fine-tuning. For instance, in supergravity, the positive vacuum energy arising after supersymmetry breaking is cancelled by adding a constant to the superpotential. Indeed the gravitino mass arises precisely from this constant term and therefore is a direct result of the fine-tuning. *Somehow* the enormous $\sim m_S^4$ that we expect from naturalness must be suppressed

$$\Lambda \sim \epsilon_4 m_S^4 \tag{1.3}$$

with $\epsilon_4 \ll 10^{-60}$. Given that this UV sensitive parameter in the low-energy theory beneath m_S is so much smaller than its natural size, why are we so confident that the other UV sensitive parameter, m_h^2, must be $\sim m_S^2$?

Again, the usual attitude is that there must be some deep new physics associated with the CC, since it has to do with gravity, with all of its associated theoretical mysteries. There doesn't seem to be anything similarly special about the Higgs mass parameter. Thus, the philosophy has been to keep the Higgs mass as natural as possible, while continuing to look for new mechanisms to solve the cosmological constant problem.

In this paper we wish to explore an orthogonal possibility. What if the observation of a tiny cosmological constant is telling us that UV sensitive parameters in the low-energy theory beneath the SUSY breaking scale will appear incredibly finely tuned? This leads us to imagine that SUSY is broken in the SSM at very high scales, far above the weak scale, with the Higgs mass parameter appearing finely-tuned in the low-energy effective theory, just as the CC appears finely tuned

$$m_h^2 \sim \epsilon_2 m_S^2 \,. \tag{1.4}$$

1.3 Cosmological constant problem and the emergence of the landscape

A possible explanation for such a pattern of fine-tunings can be found within the context of Weinberg's anthropic resolution of the CC problem [13]: if the CC was bigger than about ~ 100 times its observed value, then structure could never form in our universe; the accelerated expansion due to the CC would rip apart galaxies before they had a chance to form and the universe would quickly become empty of everything except the deSitter Hawking radiation. If there are many different vacua with different values of the CC, together with a cosmological mechanism to populate all of them, it is not surprising that we should find ourselves in a universe with a small enough CC to allow structure to form, any more than it is surprising that in our own universe we find ourselves on a tiny planet rather than in the vastly larger volume of empty space. Note that there is nothing "anthropic" about this argument, it is really invoking the "structure" principle (or "galactic" principle), the entirely reasonable statement that we shouldn't expect find ourselves in an empty universe.

This resolution of the CC problem correctly predicted a small cosmological constant, and has gained more momentum given that (A) string theory may well have an enormous "landscape" of metastable vacua required to be able to scan the CC finely enough [14 – 18], and (B) eternal inflation [19] provides a mechanism by which to populate this landscape. Both of these ingredients remain controversial [20]. Even granting these, there are many potential loopholes to the argument; for instance, if parameters other than the CC vary significantly in the landscape, then there may be bigger regions with much larger CC capable of supporting structure. Nevertheless, it is not implausible that the only parameters that can be efficiently scanned are the ones that are UV sensitive in the low-energy theory, and as such can not be controlled by symmetries.

If the structure principle and the landscape indeed explains the fine-tuning of the CC, what should we expect for the scale of SUSY breaking m_S? One might think that low-energy SUSY with $m_S \sim$ TeV is preferred, since this does not entail a large fine-tuning to keep the Higgs light. However, this conclusion is unwarranted: the enormity of the CC fine-tuning means that there are much larger factors in the measure at play. Suppose, for instance, that we have two regions in the landscape with the structure of the SSM; in one m_S is $\sim$ TeV and the Higgs mass is natural, while in the other, $m_S \sim 10^{10}$ GeV and we have to fine tune by a factor of $\sim 10^{-15}$ for the light Higgs. But suppose that in the first region there are "only" $\sim 10^{40}$ vacua (not enough to be able to find one with a small enough CC for structure formation), while in second there are $\sim 10^{200}$ vacua (which *is* enough for the tuning of the CC). Although in the first region the Higgs can be naturally light without any fine-tuning, there are simply not enough vacua to find a small enough CC, while in the second region, there are so many vacua that the additional $\sim 10^{-15}$ tuning to keep the Higgs light is a small factor in the measure. The point is clear — in the landscape picture, the measure is dominated by the requirement of getting a small enough CC, and since numbers of order 10^{60} are involved, these can dwarf the tuning required to keep the Higgs light. Without a much better understanding of the structure of the landscape, we can't decide whether low-energy SUSY breaking is preferred to SUSY broken at much higher energies. While it has been argued that low-energy SUSY may be prefered [20], this has been questioned in [21]. Furthermore, both the early proposals of [14] and the recent concrete explorations of flux vacua [15, 17] do seem to favor very high scale SUSY breaking.

If the Higgs mass has to be tuned, there must be some extension of the "structure principle" that explains why $m_h^2 \ll m_S^2$. If in addition to structure we require the existence of atoms, both Hydrogen and some atom heavier than hydrogen, this "atomic principle" can explain the need for the Higgs fine-tuning [22]. If the Higgs vev decreases by a factor of a few, the proton becomes heavier than the neutron and Hydrogen decays. If the vev increases by a factor of a few, the neutron-proton mass difference becomes far greater the the nuclear binding energy per nucleon and nuclei heavier than hydrogen decay. Adopting the "Carbonic principle", that Carbon must form, gives an even more precise determination of the Higgs vev. For very large vevs, the mass difference between the up and down quarks exceeds the color energy penalty required to have three identical quarks in a baryon, and the Δ^{++} becomes the lightest baryon. The large coulomb barriers and short-range of the strong interactions prevent the formation of nuclei with multiple Δ^{++}'s, and the only atoms in the universe would be chemically identical to Helium. The authors of [22] performed a systematic analysis of the SM varying only m_h^2, starting from a Higgs vev near the SM value and going all the way up to M_{Pl}, and found that the "atomic principle" restricts the Higgs vev to be within about a factor of ~ 5 of its observed value. It is notable that this line of reasoning also explains one of the striking facts about nature that is never addressed in conventional theories of physics beyond the SM: the remarkable proximity of the QCD and electroweak scales.

With these motivations, we will consider theories in which SUSY is broken at scales much higher than a TeV, and the fine tuning required to make the Higgs light happens by some unspecified mechanism, possibly related to whatever addresses the CCP — using, for example, the structure and atomic principles as a selection criterion for the neighborhood of the landscape that we can find ourselves in.

2. SUSY without scalars

Suppose that SUSY is broken (in the SSM sector) at a high energy m_S far above the TeV scale. The scalars of the SSM will then all be at m_S, except for one combination of the two Higgs doublets that must be finely-tuned to be light. What about the new *fermions* of the SSM, the gauginos and higgsinos? There are two possibilities: they can also be at the scale m_S, or, because they can be protected by chiral symmetries, they can survive beneath m_S. This is the possibility we wish to pursue.

One reason is that, if these fermions are also near the TeV scale, gauge coupling unification works essentially identically as in the SSM. This is because our model differs from the SSM by missing the squarks and sleptons at low energies, but these scalars come in complete SU(5) multiplets and do not affect unification at one-loop order. We are also missing the second Higgs doublet of the SSM, but this makes a comparatively small contribution to the beta function, and as we will see, not having it likely *improves* our unification prediction over the SSM when two-loop corrections are included.

An unrelated reason to expect the gauginos and higgsinos to be near a TeV, is that this mass scale is independently selected by requiring the lightest neutralino to be a good dark matter candidate in our model. Note that, since we are triggering the weak scale by a fine-tuning, there is no longer a direct link between Higgs vev and the mass of the gauginos and higgsinos. The rough link of the dark matter particle mass and the electroweak vev, which happens naturally in the SSM, is an accident in our framework. Of course, the accident is not severe; the SM itself is filled with several "accidents" in its spectrum, ranging from the proximity of the QCD and EW scales to the near equality of the muon and pion, charm and proton, etc. masses. But as we will see, in a generic class of models for supersymmetry breaking, we will in fact *predict* the gauginos and Higgsinos to be near the weak scale, following from the two other high-energy scales M_G and M_{Pl} we know of from a "see-saw" relation of the form

$$m_{1/2} \sim \frac{M_G^9}{M_{Pl}^8}\,. \tag{2.1}$$

In another class of models, $m_{1/2}$ will be generated by dimensional transmutation and again come out naturally near the TeV scale.

We now come to some phenomenological aspects of the low-energy theory.

2.1 Finely tuned Higgs

The most general structure of the low-energy lagrangian we are proposing is as follows. All the scalars of the MSSM get ultraheavy soft masses of order m_S. However, one linear combination of the two Higgs scalars,

$$h = \sin\beta h_u + \cos\beta h_d^* \tag{2.2}$$

is fine-tuned to be light.

In more detail, the two Higgs doublets $H_{u,d}$ have soft masses as well as a μB term, so the Higgs boson mass matrix is of the form

$$\begin{pmatrix} m_u^2 & \mu B \\ \mu B & m_d^2 \end{pmatrix} \tag{2.3}$$

(where we have removed a possible phase in μB by a field redefinition). Having a single light Higgs near the weak scale requires a fine-tuning. The eigenvalues of this mass matrix are

$$\bar{m}^2 \pm \sqrt{\Delta^2 + (\mu B)^2}\,, \qquad \text{where } \bar{m}^2 = \frac{m_u^2 + m_d^2}{2}\,, \qquad \Delta = \frac{m_u^2 - m_d^2}{2} \tag{2.4}$$

and we require that the smaller of these eigenvalues is negative but not larger in magnitude than $\sim -m_{EW}^2$.This requires e.g.

$$(\bar{m}^2)^2 < \Delta^2 + (\mu B)^2\,, \qquad (\bar{m}^2 + m_{EW}^2)^2 > \Delta^2 + (\mu B)^2 \tag{2.5}$$

We assume $\bar{m}^2$ and Δ randomly vary over a range $\sim m_S^2$. As for μB, it is possible that it too varies randomly over a range of size $\sim m_S^2$, however, it may be that since μB also further breaks a PQ symmetry on $H_{u,d}$, it randomly ranges over a range $\sim \epsilon m_S^2$, where ϵ is a small parameter characterizing the PQ breaking.

To see the tuning explicitly, let us fix μB at ϵm_S^2, and randomly vary $\bar{m}^2, \Delta$. The volume of the region in $(\bar{m}^2, \Delta)$ given above, where the light Higgs is in the tuned range, is then

$$\frac{V_{\text{tuned}}}{V_{\text{total}}} \sim \frac{m_{EW}^2 \int d\Delta}{m_S^4} \sim \frac{m_{EW}^2}{m_S^2} \tag{2.6}$$

exhibiting the $\sim (m_{EW}^2/m_S^2)$ tuning needed for the light Higgs. In this form it is clear that the measure of the tuned region is dominated by $\Delta \sim m_S^2$. For small ϵ, getting a light eigenvalue requires $m_u^2 m_d^2 \sim \epsilon^2 m_S^4$, but since the volume of the tuned region is dominated by $\Delta \sim m_S^2$, in most of region one of $m_{u,d}^2$ is $\sim m_S^2$ while the other is $\sim \epsilon^2 m_S^4$. If m_u^2 is the small one, the mass matrix has a "see-saw" form and $\tan\beta$ must be large

$$\tan\beta \sim \frac{1}{\epsilon} \tag{2.7}$$

which can help explain the top-bottom mass hierarchy.

There may be natural explanations for why of all the scalars in the SSM it is only the Higgs that can be light in the low-energy theory beneath m_S. For instance, suppose that the $m^2\phi^\dagger\phi$ type masses for the scalars stay positive and $\sim m_S^2$ over the whole range they are scanned. The only scalars that can even be finely tuned to be light are the ones that can have μB-type terms, and in the SSM, these are only the Higgs doublets.

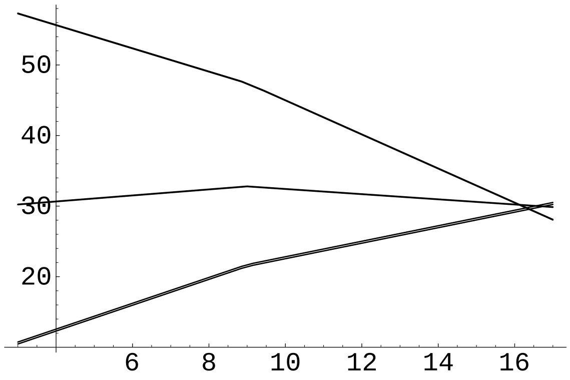

Figure 3: Running couplings in our model at one-loop, with the scalars at 10^9 GeV.

2.2 Gauge coupling unification as a signal of high-scale SUSY

In our model the gauge couplings unify essentially exactly as in the SSM. Relative to the SSM, we are missing the squarks and sleptons which come in complete SU(5) multiplets, and therefore do not affect the unification of couplings at 1-loop. We are also missing the extra scalar Higgs doublet, which as we will see does not make a significant contribution to the running.

As we will see later, cosmology favors m_S lighter than $\sim 10^{12} - 10^{13}$ GeV, and in a simple class of models for SUSY breaking we find m_S near 10^9 GeV. In all cases therefore some part of the running beneath the GUT scale reverts to the usual SUSY case. We present the 1-loop evolution of the gauge couplings for scalars at 10^9 GeV in figures 3 and 4. If as usual we use the scale where $\alpha_{1,2}^{-1}$ unify to determine the GUT scale and extrapolate back to predict $\alpha_3(M_Z)$, our one-loop prediction for $\alpha_3(M_Z) = .108$ is somewhat lower than in the usual SSM. This is welcome, because in the SSM, the two-loop running corrections push up $\alpha_3(M_Z)$ to around .130, somewhat higher than the measured central value of .119. Of course the discrepancy is parametrically within the uncertainties from GUT scale threshold corrections, although numerically these have to be somewhat large to compensate for the discrepancy. While the two-loop corrections in our case are different than in the SSM and have yet to be calculated, we expect that they will go in the same direction, pushing our somewhat low 1-loop value for $\alpha_3(M_Z)$ higher, into better agreement with experiment, requiring smaller compensating threshold corrections than in the SSM.

2.3 Effective lagrangian

The particle content in the effective theory beneath m_S consists of the Higgs h, as well as the higgsinos $\psi_{u,d}$, and the gauginos $\tilde{g}, \tilde{b}, \tilde{w}$. The most general renormalizable effective lagrangian for these fields consists of mass terms for the fermions, Yukawa couplings between the Higgs and the fermions and the Higgs quartic coupling:

$$\begin{aligned}\Delta L = M_3\tilde{g}\tilde{g} + M_2\tilde{w}\tilde{w} + M_1\tilde{b}\tilde{b} + \mu\psi_u\psi_d + \sqrt{2}\kappa_u h^\dagger\tilde{w}\psi_u + \sqrt{2}\kappa_d h\tilde{w}\psi_d + \\ +\sqrt{2}\frac{1}{2}\kappa_u' h^\dagger\tilde{b}\psi_u - \sqrt{2}\frac{1}{2}\kappa_d' h\tilde{b}\psi_d - -m^2 h^\dagger h - \frac{\lambda}{2}(h^\dagger h)^2 \,.\end{aligned} \quad (2.8)$$

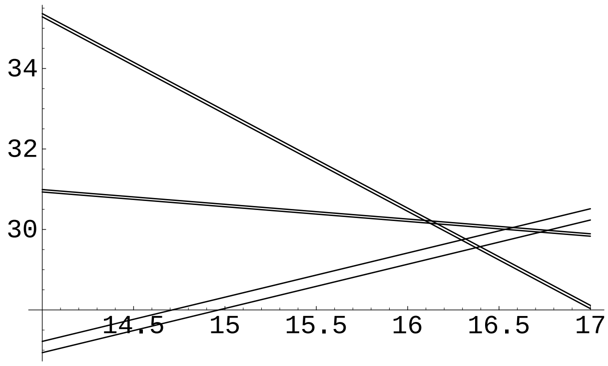

Figure 4: Close-up of the one-loop couplings near the unification scale with the heavy scalars at 10^9 GeV. Note that the prediction for $\alpha_3(M_Z)$ is lower than in the SSM. We expect two-loop corrections to push up $\alpha_3(M_Z)$ to better agreement with experiment.

Here we have assumed an analog of $R-$ parity, under which all the new states are odd. As usual, this will ensure that the lightest of the new fermions is stable and, if it is a neutralino, will be an good dark matter candidate. Note that even without imposing R-parity, there are no dimension four baryon number violating operators in the theory. The reason for imposing R-parity is not proton decay, but neutrino masses: operators of the form $lh\tilde{b}$, together with the Majorana mass term for the gauginos, do violate lepton number, giving rise to unacceptably heavy neutrinos after EWSB.

2.4 High scale SUSY boundary conditions and Higgs mass prediction

At the high scale m_S, the four dimensionless couplings $\kappa_{u,d}, \kappa'_{u,d}$ are determined at tree-level by the supersymmetric gauge Yukawa couplings of $h_{u,d}$ as

$$\kappa_u(m_S) = g_2(m_S)\sin\beta,\ \kappa_d(m_S) = g_2(m_S)\cos\beta \tag{2.9}$$

$$\kappa'_u(m_S) = \sqrt{\frac{3}{5}}g_1(m_S)\sin\beta,\ \kappa'_d(m_S) = \sqrt{\frac{3}{5}}g_1(m_S)\cos\beta \tag{2.10}$$

where we are using SU(5) normalization for hypercharge. The Higgs quartic coupling λ is determined by the supersymmetric D terms as usual

$$\lambda(m_S) = \frac{\frac{3}{5}g_1^2(m_S) + g_2^2(m_S)}{4}\cos^2 2\beta\,. \tag{2.11}$$

There can of course be threshold corrections to these relations from integrating out the heavy scalars at the scale m_S.

In order to determine the low-energy parameters, we have to run down from the high scale m_S using the RGE's for this low-energy effective theory. Note that since the theory is not supersymmetric beneath m_S, the usual supersymmetric relations between the Yukawa, quartic and gauge couplings will no longer hold.

In particular, we will have a significantly different prediction for the Higgs mass than in the SSM [23]. Usually in the SSM, there are two corrections to the Higgs quartic coupling. First, integrating out the stops generates a threshold correction to λ parametrically of order

$(3/8\pi^2)(A_t/m_{\tilde{t}})^4$, where A_t is the A- parameter associated with the top Yukawa coupling, which can be a large correction. Second, there is a log enhanced contribution to λ from the top loop in the low-energy theory beneath $m_{\tilde{t}}$. Normally with all the scalars near the TeV scale, these effects are comparable in size, the logarithm is not particularly big, a full 1-loop analysis is needed, and the Higgs mass prediction depends on the details of the A-terms and stop spectrum.

The situation is different in our case. First, the same physics that suppresses the gaugino masses will inevitably also suppress the A terms so that $A_t \ll m_S$, and the threshold correction to the Higgs quartic coupling from integrating out the stops at m_S is tiny, therefore the boundary value for $\lambda(m_S)$ is accurately given by the tree-level value. Second, with very high m_S, the low-energy Higgs quartic coupling is controlled by the logarithmically enhanced contribution given by the running the RGE for λ to low energies. This running is quickly dominated by the contribution from the top Yukawa couplings, and we obtain a prediction for the Higgs mass.

At 1-loop, the RGE for λ in the theory beneath m_S is

$$\begin{aligned} 16\pi^2 \frac{d\lambda}{dt} &= 12\lambda^2 + \lambda\left(12\lambda_t^2 + 6\kappa_u^2 + 6\kappa_d^2 + 2\kappa_u'^2 + 2\kappa_d'^2\right) - 3\lambda\left(3g_2^2 + \frac{3}{5}g_1^2\right) + \\ &+ \frac{3}{4}\left(2g_2^4 + \left(g_2^2 + \frac{3}{5}g_1^2\right)^2\right) - \left(12\lambda_t^4 + 5\kappa_u^4 + 5\kappa_d^4 + \kappa_u'^4 + \kappa_d'^4\right) - \\ &- \left(2\kappa_u^2\kappa_u'^2 + 2\kappa_d^2\kappa_d'^2 + 2\kappa_d^2\kappa_u^2 + 2\kappa_u'^2\kappa_d'^2 + 4\kappa_u\kappa_u'\kappa_d\kappa_d'\right). \end{aligned} \tag{2.12}$$

A complete analysis of the Higgs mass prediction at one loop would require solving the coupled RGE's for λ together with the top Yukawa coupling λ_t and the κ's. But the largest contribution to λ come from the top Yukawa and are $\propto \lambda_t^4$, so the Higgs mass depends most sensitively on the top Yukawa coupling. Extracting λ_t from the top mass at tree-level gives $\lambda_t = 1$, however, when the 1-loop QCD corrections to the top mass are taken into account, this is reduced to $\lambda_t = .95$, decreasing the Higgs mass prediction by ~ 20 %. Meanwhile, the recent CDF/D0 measurements of the top mass point to a somewhat heavier top quark $\sim 178\,\text{GeV}$ [24]. These uncertainties involving the top Yukawa have a bigger impact on our Higgs mass prediction than all the terms involving the κ and $g_{2,1}$ couplings in the λ RGE's. Nonetheless as a preliminary analysis that we believe will capture the most important effects, we numerically solve the 1-loop RGE's for λ, λ_t and the κ's, using the boundary condition $\lambda_t(m_t) = .95$ for $m_t = 174\,\text{GeV}$ and $\lambda_t(m_t) = .97$ for $m_t = 178\,\text{GeV}$, while keeping only the contributions from the largest couplings λ_t, g_3 in the λ_t, κ runnings

$$16\pi^2 \frac{d\lambda_t}{dt} = \lambda_t\left(\frac{9}{2}\lambda_t^2 - 8g_3^2\right) + \cdots \tag{2.13}$$

$$16\pi^2 \frac{d\kappa_{u,d}^{(\prime)}}{dt} = 3\lambda_t^2\kappa_{u,d}^{(\prime)} + \cdots \tag{2.14}$$

and using the high-scale SUSY boundary conditions for all the couplings. The prediction for the low-energy Higgs mass

$$m_h \sim \sqrt{\lambda}v \tag{2.15}$$

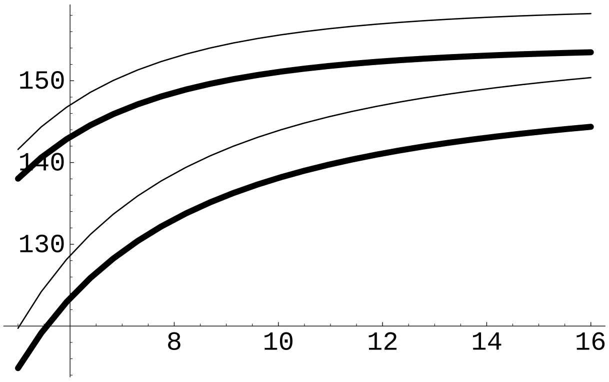

Figure 5: The Higgs mass in GeV, as a function of $\log_{10}(m_S/\mathrm{GeV})$. The thick line is for $m_t = 174\,\mathrm{GeV}$, the thin line for $m_t = 178\,\mathrm{GeV}$, while lower lines are for $\cos 2\beta = 0$ and the upper lines for $\cos 2\beta = 1$.

with λ evaluated at $\sim m_t$, is plotted in figure 5, as a function of the the scale m_S and $\tan\beta$. We find m_h in the range between $\sim$ 120-150 GeV, light but above the LEPII limits.

The Higgs mass can give a first, very rough estimate of the scale m_S. But in principle, all the couplings $\lambda, \kappa_{u,d}, \kappa'_{u,d}$ can be measured at low energies, and running them to high energies should show that all five of them hit their supersymmetric values at the same scale m_S. A convenient digramatic representation of this is to group $\kappa_{u,d}$ into a two-dimensional vector $\vec{\kappa} = \kappa_d \hat{x} + \kappa_u \hat{y}$, and $\kappa'_{u,d}$ into the vector $\sqrt{3/5}\vec{\kappa}' = \kappa'_d \hat{x} + \kappa'_u \hat{y}$. Running these vectors to high energies, the SUSY boundary conditions tell us that at the scale m_S these two vectors must be aligned in the same direction, with angle from the horizontal β, and that the lengths of $\vec{\kappa}, \vec{\kappa'}$ should be $g_2(m_S), g_1(m_S)$ respectively. Having determined β, one can also check that the running Higgs quartic $\lambda(m_S)$ hits its supersymmetric value. These checks are illustrated in figure 6. Clearly if all of these measurements were made and these predictions confirmed, it would be striking quantitative evidence for *high scale* supersymmetry at a scale m_S, with a finely tuned Higgs in the theory beneath m_S.

We have discussed the Higgs mass and κ, κ' predictions in our minimal model, but these can change in less minimal models by a number of new factors absent in the usual SSM. Since the evolution of the couplings beneath m_S is non-supersymmetric and the supersymmetric link between $\lambda, \kappa_{u,d}, \kappa'_{u,d}$ and the gauge couplings $g_{1,2}$ no longer holds, the presence of additional vector-like matter multiplets (say a number of $(5+\bar{5})$'s) in the theory beneath m_S can affect the Higgs mass prediction. There may also be new contributions to the Higgs quartic coupling at m_S, coming from additional superpotential or D-term couplings. These are now only constrained by the requirement of perturbativity from m_S to M_G, and may therefore give larger corrections to the Higgs mass than in the usual SSM. These interesting issues deserve further exploration.

2.5 Long-lived gluino as a probe of fine-tuning

A striking qualitative prediction of our new framework, decisively differentiating it from the usual SSM, is the longevity of the gluino. Because the scale of supersymmetry breaking is now high, the squarks are heavy and the lifetime for the gluino to decay into a quark,

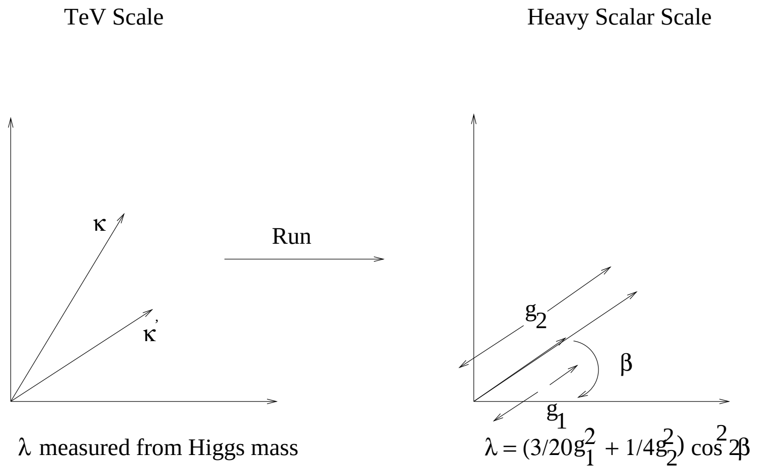

Figure 6: Evidence for high scale SUSY from running the couplings $\kappa_{u,d}, \kappa'_{u,d}$, grouped into 2-d vectors, from low to high energies. The two vectors align at a scale m_S where they must have lengths $g_{2,1}$. This must agree with the m_S extracted from the gluino lifetime. β is determined, and fixes $\lambda(m_s)$, which can be checked against the λ determined by the Higgs mass.

antiquark and LSP — which is mediated by virtual squark exchange — becomes very long, of order

$$\tau = 3 \times 10^{-2}\mathrm{sec}\left(\frac{\mathrm{m_S}}{10^9\,\mathrm{GeV}}\right)^4\left(\frac{1\,\mathrm{TeV}}{\mathrm{m}_{\tilde{g}}}\right)^5, \tag{2.16}$$

where m_S is the squark mass, $m_{\tilde{g}}$ the gluino mass. We have included a QCD enhancement factor of ~ 10 in the rate, as well as another factor ~ 10 for the number of decay channels. The gluino lifetime can easily range from 10^{-6} sec to the age of the universe, as m_S ranges from 10^8 GeV to 10^{13} GeV, and $m_{\tilde{g}}$ from 100 GeV to 1 TeV. As long as its lifetime is much longer than 10^{-6} sec, a *typical* gluino produced at the LHC will decay far outside the detector. This is a key difference between our theories and the SSM, fundamentally tied to the large-scale breaking of supersymmetry, and, once the gluino is produced at the LHC, can immediately experimentally distinguish our model from a conventional hierarchy-motivated SUSY theory with scalars just barely too heavy to be produced (say at ~ 10 TeV), where the gluino would still decay well inside the detector.

The only trace of a typical gluino decaying outside the detector will be the energy that it deposits in the detector [25 – 27]. However, at peak luminosity of 30 fb^{-1} per year, the LHC may well be a gluino factory producing roughly a gluino per second (for $m_{\tilde{g}} \sim 300$ GeV). It is therefore possible to get statistically important information by relying on *atypical* events involving:

Displaced gluinos: these are simply gluinos which decay in the detector, even though their lifetime is longer than the size of the detector. The number of these events will become too small once the lifetime becomes longer than roughly one second. For $m_S \sim 1000$ TeV or so, all the gluinos will decay inside the detector, but may live long enough to have displaced vertices.

Stopped, long-lived gluinos: these are gluinos which lose energy and stop in the detector or in the surrounding earth, and decay much later — seconds, days, or months later, perhaps even when there is no beam in the accelerator! The lifetime sensitivity can extend to $\sim 10^7$ years, corresponding to a SUSY breaking scale m_S of up to $\sim 10^{13}\,\mathrm{GeV}$, and depends on the fraction of gluinos that stop in the detector; this is involved and in turn depends on the fraction of time the gluino dresses into a charged hadron and loses energy electromagnetically. Such events can be spectacular and give charged tracks, displaced vertices, delayed decays, and possibly even intermittent tracks, all at the same time. It is also possible to just have displaced vertices, and delayed decays, without charged tracks. Since the final decays can occur much after the collision that created the gluino, triggering on these poses interesting challenges. For long lifetimes, a good time to look for such events is when there is no beam, but the detectors are on. A particularly good place to look is at the endpoints of charged tracks.

2.6 Gluino cosmology

In our framework, there are two particles that are potentially important for cosmology. One is the LSP neutralino, a natural candidate for the DM particle [2], as in most versions of the SSM. Another is the gluino, which is now long lived.

Gluinos can be cosmologically excluded either because their abundance today is unacceptably large or, if their lifetime is shorter than the age of the universe, their decay products can distort the photon background or destroy nuclei synthesized during primordial nucleosynthesis, which began when the universe was one second old. A gluino that decays in less than a second is harmless, as its decay products thermalize and the heat bath erases any trace of its existence. Gluinos that live longer than a second can be safe, as long as their abundance is small.

We now turn to an estimate of the abundance of gluinos before they decay. When the temperature of the universe drops below $m_{\tilde{g}}$, the gluino's abundance is maintained in thermal equilibrium by their annihilation into gluons. Eventually, their abundance becomes so low that they cannot find each other in the expanding universe, they stop annihilating, and their abundance "freezes out". There are three stages of gluino annihilation, characterized by three different processes by which gluinos can annihilate. In chronological order, they are: perturbative annihilation, annihilation via recombination, and annihilation in the QCD era. The first ends at a temperature of order $T_F \sim 1/27\, m_{\tilde{g}}$ and, as long as $T_F \gg \Lambda_{QCD}$, leads to the canonical fractional abundance today of order,

$$\frac{\rho_{\tilde{g}}}{\rho_T} = \frac{m_{\tilde{g}}^2}{N\,\alpha_s^2\,(100\,\mathrm{TeV})^2} \tag{2.17}$$

where N is a numerical factor, depending on the number of decay channels, color factor etc., which we estimate to be of order 100. Next, when the universe cools down to a temperature below the gluinium (a bound state of two gluinos) binding energy E_B,

$$T \lesssim E_B = \frac{1}{2}\, m_{\tilde{g}}\, \alpha_s^2 \sim 5\,\mathrm{GeV}\left(\frac{m_{\tilde{g}}}{1\,\mathrm{TeV}}\right) \tag{2.18}$$

a new period of possible gluino annihilation begins. Gluino pairs can now "recombine" (via gluon emission) into a gluinium which is no longer broken apart by the ambient thermal gluons. The recombination cross section, however, is comparable to the perturbative annihilation cross section, so no significant further gluino annihilation occurs.

Next, the temperature drops to Λ_{QCD} and the strongly interacting particles organize themselves into a dilute gas of color-singlet baryons, R-hadrons and gluinia. What happens beyond this is difficult to analyze quantitatively, as it involves hadron dynamics [25, 26].

One scenario that we consider plausible is that when two slow R-hadrons collide they recombine into a bound R-molecule (by emitting a pion), containing two gluino "nuclei", with a cross section of order $\sigma \sim 30$ mb. Subsequently, the two gluinos inside this small molecule rapidly find and annihilate each other into gluons, before the molecule has a chance to be dissociated by collisions with the dilute gas particles. This avoids the suppression $\sim m_{\tilde{g}}^{-2}$ in the perturbative annihilation cross-section, and results in a small gluino abundance which we estimate by equating expansion and reaction rates,$n\,\sigma\,v \sim T^2/M_{Pl}$ where $T \sim \Lambda_{QCD}$. This translates to

$$\frac{n_{\tilde{g}}}{n_\gamma} = 10^{-18}\left(\frac{m_{\tilde{g}}}{1\,\text{TeV}}\right)^{1/2} \tag{2.19}$$

$$m_{\tilde{g}}\,\frac{n_{\tilde{g}}}{n_\gamma} = 10^{-15}\left(\frac{m_{\tilde{g}}}{1\,\text{TeV}}\right)^{3/2}\ \text{GeV} \tag{2.20}$$

The last quantity measures the destructive power of the decaying gluino gas, as it depends on both the mass and the concentration of gluinos. The abundance of gluinos with lifetimes comparable to the age of the universe is constrained by the the negative searches for abnormally heavy hydrogen, helium, and lithium that would have formed during primordial nucleoynthesis, as well as the limits on stable strongly interacting massive particles [28] that would result from the pairing of a gluino and a gluon. These limits are much stronger than the abundance of equation(2.19), in the case of heavy hydrogen by a factor of 10^{22} [29]. So, the gluino lifetime must be shorter than roughly 10^{16} seconds, corresponding to an upper limit of about $m_S \lesssim 3\times 10^{13}\,\text{GeV}$, for a $1\,\text{TeV}$ gluino. Evading this cosmological limit on m_S is possible in theories where the reheat temperature is much lower than the gluino mass, so that gluinos are not produced after reheating.

The abundance of gluinos with lifetime up to 10^{13} sec must be small to avoid spectral distortions of the CMBR [30]. This constraint is mild, and equation (2.20) easily satisfies it. The abundance of gluinos with lifetime in the range from 10^{-1} sec to 10^{12} sec must also be small to avoid the destruction of the light nuclei synthesized during the BBN [31, 32]. Although this constraint is strong, especially for lifetimes between 10^4 sec to 10^7 sec, equation (2.20) satisfies it. Other constraints from possible distortions of the diffuse photon background are easily satisfied.

The problem of computing the gluino abundance through the QCD era is important and should be revisited [25]. We stress that our picture for gluino annihilation after the QCD phase transition is rough and may be missing important effects that suppresses the annihilation and increases the abundance, which may lead to better limits for the gluino mass and the scale of SUSY breaking in our framework.

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2.7 Addressing the problems of the SSM

The SSM has many phenomenological problems associated with the 110 independent parameters in the flavor sector alone [33]. These problems originate in the 97 parameters that reside in the scalar sector, in the mass- and A-matrices of squarks and sleptons (the rest are just the usual KM parameters); it is hard to hide all 97 parameters of the flavor sector of the SSM from low energy physics and avoid problems with a large number of rare processes such as, FCNCs, CP-violation, b-decays — when the scalars are near a TeV, as required by the naturalness. Starting with the original universality hypothesis [2], much of the model building in the last 23 years has been targeted to solving these flavor problems by attempting to derive universality from some specific dynamics – such as gravity [34], gauge [35, 36], anomaly [37, 38] and gaugino [39] mediation — in spite of the violations of flavor in the Yukawa couplings of quarks. In addition there are difficulties associated with dimension five proton decay operators and CP violating SUSY phases. Meanwhile, the absence of a light Higgs at LEPII has raised new problems, necessitating tunings at the few percent level for electroweak symmetry breaking [41].

All these problems evaporate as soon as we raise the scale of sparticles to $\sim 100 - 1000\,\text{TeV}$. The physical relevance of all the 97 parameters connected to the flavor problem disappears because they are linked to the scalars that now decouple. Similarly for the proton lifetime via the dimension 5 operators. As we have seen the light Higgs is naturally heavier; of course there is tuning to get the Higgs mass light, but unlike the usual SSM, naturalness is not our guiding principle, and we have argued that this tuning is taken care of by the "atomic principle". Finally, while there may be phases in μ and the M_i, these first affect only the Higgs sector at 1-loop, and only much more indirectly feed into the electron and neutron edms, which are naturally small enough.

In addition the SSM has problems of cosmological origin, the gravitino and moduli problems. As soon as the scale of SUSY breaking is raised to over $\sim 100\,\text{TeV}$, the gravitino and moduli decay with lifetimes less than a second, and these problems also evaporate.

3. Models

We now give examples of models where there is a natural separation of scales between the scalar and the gaugino/higgsino masses, with chiral symmetries keeping the fermions light relative to the scalars. We will begin by considering very standard sorts of models where the low-energy theory beneath the cutoff contains supergravity. In such a theory, the only way to cancel the vacuum energy after supersymmetry breaking is to add a constant c to the superpotential, which breaks R-symmetry and makes the gravitino massive. Since R is neccessarily broken, at some level we must induce a gaugino mass.

3.1 Anomaly mediated gaugino masses with scalars at $\sim 1000\,\text{TeV}$

If we assume that the field Z breaking SUSY ($F_Z \neq 0$) carries some symmetry so that the operator $\int d^2\theta ZWW$ is forbidden, then the leading source for a gaugino mass is from anomaly mediation [37, 38]. Since R is broken, the F component of the chiral compensator field ϕ in supergravity can be non-zero, yielding gaugino masses of order

$$m_{1/2} \sim \frac{g^2}{16\pi^2} F_\phi \,. \tag{3.1}$$

If F_ϕ is $\sim m_{3/2}$, this limits $m_{3/2} \lesssim 50\,\text{TeV}$ for gauginos near the TeV scale. In [38] and more recently in [42], the phenomenology of scalars with mass $\sim m_{3/2}$ has been explored. From our point of view, however, this is not heavy enough — scalars at $\sim 50\,\text{TeV}$ are in an uncomfortable no-man's land between being natural and tuned, and do not in themselves solve e.g. the SUSY flavor problem outright. Fortunately, the scalars can be much heavier than $m_{3/2}$, since the operators in the Kähler potential giving the scalar masses can be suppressed by a fundamental scale $M_* \sim M_G$ much smaller than the Planck scale, so that

$$m_{3/2} \sim \frac{F_Z}{M_{Pl}}, \sim 50\,\text{TeV}\,, \qquad m_S \sim \frac{F_Z}{M_*} \sim 500 - 5000\,\text{TeV}\,. \tag{3.2}$$

For μ and μB, we can simply write down the usual Giudice-Masiero [43] operators

$$\int d^4\theta \frac{Z^\dagger}{M_*} H_u H_d \,, \qquad \int d^4\theta \frac{Z^\dagger Z}{M_*^2} H_u H_d \,. \tag{3.3}$$

If both of the above operators have an additional suppression by a factor of ϵ since they break a PQ symmetry, then we have

$$(\mu B) \sim \epsilon (100\,\text{TeV})^2 \,, \qquad \mu \sim \epsilon (100\,TeV)\,. \tag{3.4}$$

Recall that small ϵ leads to large $\tan\beta \sim 1/\epsilon$, which is natural to be $\sim m_t/m_b$; in this case we can find μ close to the TeV scale

$$\mu \sim \frac{m_b}{m_t} \times 100\,\text{TeV} \sim \,\text{TeV}\,. \tag{3.5}$$

This set-up is very generic, and it can push scalars up to masses $\gtrsim 1000\,\text{TeV}$, high enough to evade all phenomenological problems. This is also in an interesting range for gluino collider phenomenology: all the gluinos can decay inside the detector, but with a long enough lifetime to have observable displaced vertices. Together with the gaugino masses, this is a smoking gun for anomaly mediation with ultraheavy scalars. In addition, as usual in anomaly mediation, with $m_{3/2} \sim 50\,\text{TeV}$ the gravitino and moduli problems disappear.

3.2 Theories with $\sim 100\,$GeV Gauginos and Higgsinos: generalities

It is possible to consider a large class of theories where the $F_\phi \ll m_{3/2}$, and the anomaly mediated gaugino masses are negligible relative to other sources of SUSY breaking. The leading R-invariant operator we can write down that generates a gaugino mass directly from R and SUSY breaking is

$$\int d^4\theta \frac{Z^\dagger Z c^\dagger}{M_{Pl}^5} W_\alpha W^\alpha \,. \tag{3.6}$$

Similarly, the leading R- invariant operator leading to a non-zero μ (given $H_{u,d}$ have R-charge 0) is

$$\int d^4\theta \frac{Z^\dagger Z c}{M_{Pl}^5} D_\alpha H_u D_\alpha H_d\,. \tag{3.7}$$

Recalling that the gravitino mass is $m_{3/2} \sim |c|/M_{Pl}^2$, and that we must have $3|c|^2/M_{Pl}^2 = F_X^2$ to cancel the vacuum energy, leads to a gaugino/Higgsino mass

$$m_{1/2} \sim \mu \sim \frac{m_{3/2}^3}{M_{Pl}^2}\,. \tag{3.8}$$

If we assume that $c \sim M_G^3$, the most natural value in a theory where the fundamental scale is near M_G, we find

$$m_{3/2} \sim \frac{M_G^3}{M_{Pl}^2}, \sim 10^{13}\,\mathrm{GeV},\;\; m_{1/2} \sim \frac{m_{3/2}^3}{M_{Pl}^2} \sim \;\mathrm{TeV} \tag{3.9}$$

so the gauginos and higgsinos are very naturally near the TeV scale! While these estimates are rough, we have found a new link between the TeV scale, here setting the dark matter mass, and the GUT/Planck hierarchy. The scalar masses are more dependent on the details of SUSY breaking. We next discuss a concrete model implementing these ideas.

3.3 Scherk-Schwarz models

Models with $F_\phi \ll m_{3/2}$ arise naturally in the context of no-scale models [44], which can arise from Scherk-Schwarz SUSY breaking in extra dimensions [45], or equivalently from SUSY breaking by the $F-$ component of a radion chiral superfield $T = r + \theta^2 F_T$ [46]. We follow the discussion of [47]. Consider an extra dimension which is an interval, and add a constant superpotential localized on one of the boundaries, say the right boundary. The tree-level low-energy effective lagrangian for T and the chiral compensator $\phi = 1 + \theta^2 F_\phi$ is of the form

$$L = \int d^4\theta M_5^3 \phi^\dagger \phi (T + T^\dagger) + \int d^2\theta c M_5^3 \phi^3 + \mathrm{h.c} \tag{3.10}$$

leading to the scalar potential

$$V = M_5^3 \left[r|F_\phi|^2 + F_T^* F_\phi + 3cF_\phi + \mathrm{h.c.} \right]. \tag{3.11}$$

The F_T equation of motion fixes

$$F_\phi = 0 \tag{3.12}$$

while the F_ϕ equation of motion fixes

$$F_T = -3c\,. \tag{3.13}$$

Therefore supersymmetry is broken, with vanishing tree-level potential. This is the famous "no-scale" structure. The goldstino is the fermionic component of T, which is eaten by the gravitino, and the mass is

$$m_{3/2} \sim \frac{c}{r} \tag{3.14}$$

while F_ϕ vanishes at this level. Hereafter we will assume $c \sim 1$.

At 1-loop, a non-zero potential is generated-the gravitational Casimir energy-which arises from a contribution to the Kähler potential of the form

$$\int d^4\theta \frac{1}{16\pi^2} \frac{1}{(T+T^\dagger)^2} \tag{3.15}$$

yielding

$$V_{\text{grav}} \sim -\frac{1}{16\pi^2} \frac{1}{r^4} \,. \tag{3.16}$$

This tends to make the radius shrink. However, the addition of N bulk hypermultiplets of mass M can give rise to a repulsive Casimir energy of the form

$$V_{Hyper} \sim +\frac{N}{16\pi^2} \frac{1}{16\pi^2} e^{-Mr} \,. \tag{3.17}$$

Therefore there can be a minimum of the potential around $r \sim M^{-1}$. The F_T equation of motion now forces

$$F_\phi \sim \frac{1}{16\pi^2} \frac{F_T}{M_5^3 r^4} \sim \frac{1}{16\pi^2} \frac{1}{M_5^3 r^4} \tag{3.18}$$

so clearly $F_\phi \ll m_{3/2}$.

The value of the potential is negative at the minimum; in order to cancel the cosmological constant, we have to have an additional source of SUSY breaking on one of the branes, a superfield X with $F_X \neq 0$ and

$$|F_X|^2 \sim \frac{1}{16\pi^2} \frac{1}{r^4} \,. \tag{3.19}$$

To be concrete, suppose we have a chiral superfield X localized on the left boundary with a superpotential

$$W = \phi^3 m^2 X \tag{3.20}$$

and a Kähler potential

$$K = \phi^\dagger \phi \left(X^\dagger X - \frac{(X^\dagger X)^2}{M_5^2} + \text{higher powers of } X^\dagger X \right) . \tag{3.21}$$

This form of W and K can be guaranteed by an R-symmetry under which X has charge 2, although this is not necessary. If we instead only assume that X carries a spurious U(1) charge -2 under which m has charge +1, so that X only appears in the combinations $X^\dagger X$ and $m^2 X$, our conclusions are unaltered.

Evidently $F_X \sim m^2$, in order to cancel the vacuum energy, we have to choose

$$m^2 \sim \frac{1}{4\pi} \frac{1}{r^2} \,. \tag{3.22}$$

The $(X^\dagger X)^2$ term in K give the scalar component of X a positive mass squared $\sim m^4/M_5^2$. The non-vanishing F_ϕ gives a linear term to X from the superpotential coupling, so X gets a vev. We then have a local minimum with

$$F_X \sim m^2 \,, \qquad X \sim \frac{m^2}{M_5} \,. \tag{3.23}$$

The combination of the non-vanishing F_X and X also gives the fermionic component of X, ψ_X, a mass of order

$$m_{\psi_X} \sim \frac{m^4}{M_5^3}\,. \tag{3.24}$$

From the $(X^\dagger X)^2$ part of K.

At this point we have broken SUSY, stabilized the various moduli and fine-tuned away the vacuum energy. The masses of all fields can be expressed in terms of the 5D and 4D Planck scales by using the usual flat space relationship $M_4^2 \sim M_5^3 r$, and we find

$$m_{3/2} \sim \frac{M_5^3}{M_4^2}\,, \qquad m_{\rm radion} \sim \frac{M_5^6}{4\pi M_4^5}\,, \qquad m_X \sim \frac{M_5^5}{4\pi M_4^4}\,, \qquad m_{\psi_X} \sim \frac{M_5^9}{16\pi^2 M_4^8}\,. \tag{3.25}$$

The spectrum of the rest of the superpartners now depends on their location in the bulk. We will assume that the SSM fields are localized on the same brane as X, and therefore direct mediation of SUSY breaking to the SSM scalars through operators of the form

$$\int d^4\theta \frac{1}{M_5^2} X^\dagger X Q^\dagger Q \tag{3.26}$$

are unsuppressed, leading to scalars masses of the same order as m_X:

$$m_S \sim \frac{|F_X|}{M_5} \sim \frac{M_5^5}{4\pi M_4^4}\,. \tag{3.27}$$

What about the gaugino masses? An irreducible source of R- breaking is through the gravitino mass $m_{3/2}$. There is then a finite 1-loop diagram, involving a propagator stretching between the two boundaries of the extra dimension, giving a gaugino mass. The magnitude can be estimated by drawing the diagram in the low-energy theory cut-off off the scale $1/r$. The result is equal for all gauginos and is

$$M_{3,2,1}^{\rm grav} \sim \frac{1}{16\pi^2 M_4^2}\left(\frac{1}{r}\right)^3 \sim \frac{M_5^9}{16\pi^2 M_4^8} \tag{3.28}$$

of the same order as m_{ψ_X}. Note that the operator corresponding to this gaugino mass must is of the general form we considered in the previous subsection,

$$\sim \int d^4\theta \frac{c^\dagger}{M_5^3 (T+T^\dagger)^2} W_\alpha W^\alpha\,. \tag{3.29}$$

This dominates over the anomaly mediated contribution by a perturbative loop factor

$$M_{3,2,1}^{\rm anom} \sim \frac{g^2}{16\pi^2} F_\phi \sim \frac{g^2}{16\pi^2} M^{\rm grav}\,. \tag{3.30}$$

In addition, we can have contact interactions on our brane which can give rise to gaugino masses. The leading allowed operators are of the form

$$\int d^2\theta \frac{m^2 X}{M_5^3} WW\,; \qquad \int d^4\theta \frac{X^\dagger X}{M_5^3} WW + \text{h.c.} \tag{3.31}$$

which generate gaugino masses of order

$$M_i \frac{|F_X|^2}{M_5^3} \sim M^{\text{grav}} \tag{3.32}$$

which are comparable to the gravitationally induced masses, and can be different for $M_{3,2,1}$.

We now turn to μB and μ. As in the MSSM, something must have suppressed the $M_5 H_u H_d$ term in the superpotential. This could for instance be due to an accidental R-symmetry at the level of the renormalizable couplings of the theory, under which $H_{u,d}$ carry R- charge 0. The leading allowed couplings are then

$$\int d^2\theta \frac{m^2 X}{M_5^2} H_u H_d\,; \qquad \int d^4\theta X^\dagger X H_u H_d\,; \qquad \int d^4\theta \frac{m^2 X^\dagger}{M_5^3} H_u H_d \tag{3.33}$$

these generate a μB term of the appropriate size

$$\mu B \sim \frac{|F_X|^2}{M_5^3} \sim m_S^2\,. \tag{3.34}$$

As well as a μ term of the same order as the gaugino masses

$$\mu \sim M_i\,. \tag{3.35}$$

So, we have presented a simple model which breaks supersymmetry with a stabilized extra dimension, producing an interesting hierarchy of scales for the gravitino, scalars, gauginos and Higgsinos of the theory:

$$m_{3/2} \sim \frac{\pi M_5^3}{M_4^2}\,; \qquad m_S \sim \frac{\pi M_5^5}{M_4^4}\,; \qquad M_i, \mu, m_{\psi_X} \sim \frac{\pi M_5^9}{M_4^8} \tag{3.36}$$

where in the above we have been more careful about the 2π factors involved in the ratio of 5D and 4D Planck scales.

Let us get an idea for the scales involved. It is natural to use this extra dimension to lower the higher dimensional Planck scale down to the GUT scale, a la Horava-Witten [48], $M_5 \sim M_G \sim 3 \times 10^{16}\,\text{GeV}$. Then we have

$$m_{3/2} \sim 10^{13}\,\text{GeV}\,; \qquad m_S \sim 10^9\,\text{GeV}\,; \qquad m_{\text{radion}} \sim 10^7\,\text{GeV}\,; \qquad M, \mu \sim 100\,\text{GeV}\,. \tag{3.37}$$

Note that even though there was no a priori reason for the gauginos and Higgsinos to end up anywhere near the $\sim 100\,\text{GeV}$ scale, they are in the right ball-park from this simple estimate.

We can also contemplate other sorts of theories of SUSY breaking on the SSM brane where R is more badly broken, such that X and F_X are set by the same scale $X \sim \sqrt{F_X}$. In this case the operator $X^\dagger X H_u H_d$ also generates a μ term of the order of

$$\mu \sim \frac{X^* F_X}{M_5^2} \tag{3.38}$$

which, if we make the reasonable assumption that $|X| \sim \sqrt{|F_X|}$, gives rise to the estimate

$$\mu \sim \left(\frac{1}{4\pi}\right)^{3/2} \frac{M_5^7}{M_4^6} \tag{3.39}$$

which is much bigger than $M^{\rm grav}$ or M_i. This is also a potentially interesting scenario, since for the same parameters as as above this μ is $\sim 100\,$TeV. Again, if the coefficient of the $H_u H_d$ are suppressed by a factor ϵ, this leads naturally to a large $\tan\beta \sim 1/\epsilon$ which can be $\sim 10^{-2}$, so that μ can be suppressed by a further $\tan\beta$ factor to be near the $\sim$TeV scale.

It should be clear that these theories are not particularly engineered; we are breaking SUSY in one of the simplest possible ways, and then simply following our nose to cancel the vacuum energy and stabilize all the moduli. It certainly seems possible that this sort of mechanism could be "generic" within a large neighborhood of the landscape.

Note that in addition to the SSM fermions, we have an additional light fermion ψ_X of comparable mass. Unlike the gravitino of the usual SSM, however, this new particle does not give rise to cosmological difficulties. If we assume that it is heavier than the LSP, it can rapidly decay to it via the couplings giving rise to the μ term; for instance the $(m^2/M_5^2) X H_u H_d$ operator gives rise to a decay width for ψ_X of order

$$\Gamma \sim \frac{m^4}{M_5^4} m_{\psi_X}, \tau \sim 10^{-10}{\rm s} \tag{3.40}$$

for $m_{\psi_X} \sim$TeV. This decay happens well before nucleosynthesis and poses no cosmological dangers.

3.4 Non-SUGRA models with gauginos/Higgsinos $\sim 100\,$GeV

It is also possible that supersymmetry is so badly broken in the gravitational sector of the theory that supergravity is not a good low-energy approximation, but that nevertheless in the $M_{Pl} \to \infty$ limit a globally supersymmetric field theory sector is recovered. This can certainly be done within a consistent effective theory, and amounts to working with a fixed cutoff M_* which we will take to be $\sim M_G$, and including hard SUSY breaking spurions suppressed by appropriate powers of (M_G/M_{Pl}) in the effective theory. In this case, we can write an effective action of the form e.g.

$$\int d^4x \frac{1}{\epsilon}\sqrt{-g} M_G^2 R + \mathcal{L}_{SSM} + \epsilon M_G^2 \left(\tilde{q}^* \tilde{q} + \tilde{l}^* \tilde{l} + h_{u,d}^* h_{u,d} + h_u h_d\right) + \cdots \tag{3.41}$$

where the spurion ϵ is

$$\epsilon \sim \left(\frac{M_G}{M_{Pl}}\right)^2 \tag{3.42}$$

and we can also expect corrections to the dimensionless SUSY couplings of $O(\epsilon)$. The scale m_S is then naturally

$$m_S \sim \frac{M_G^2}{M_{Pl}} \sim 10^{13} - 10^{14}\,{\rm GeV}\,. \tag{3.43}$$

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In such a model, there is no a priori need to break R in order to cancel the vacuum energy. It is amusing to contemplate the R- symmetric limit of the SSM [49, 50]. The most immediate problem is the massless gluino; although perhaps this can be hidden in the QCD muck [51]. However, in the usual SSM, the R-symmetric limit also suffers from problems in the electroweak sector: the charginos and one of the neutralinos do get masses from electroweak symmetry breaking, but they are too light: the sum of the chargino masses is smaller than $2m_W$, and the one of the neutralinos is degenerate with the Z, both of which have been ruled out at LEP II. However in our model with very high m_S, the low-energy Yukawa couplings $\kappa_{u,d}, \kappa'_{u,d}$ are no longer forced to equal the gauge couplings g_2, g_1 by SUSY, and in fact grow relative to $g_{2,1}$ by RG scaling, which is also sensitive to the possible presence of additional $(5 + \bar{5})$ multiplets beneath the GUT scale. Thus the chargino/neutralino masses can become heavier, and it may be possible to evade these direct detection limits on the electroweak-inos even in the R-symmetric limit.

Nevertheless, since the massless gluino is so problematic [52], it is more reasonable to imagine that gaugino and Higgsino masses are generated from some source of spontaneous R- breaking in the low-energy theory. For instance, we can have a hidden sector gauge group G with fermions ψ, ψ^c with R-charge -1 (like the Higgsinos), and R- symmetric higher dimension operators linking that sector to ours via

$$\frac{\epsilon}{M_G^2}\psi^c\psi\lambda\lambda\,, \qquad \frac{\epsilon}{M_G^2}\bar{\psi}^c\bar{\psi}\psi_u\psi_d \tag{3.44}$$

then if the ψ, ψ^c condense at a scale Λ we generate gaugino/Higgsino masses

$$m_{1/2} \sim \mu \sim \epsilon\frac{\Lambda^3}{M_G^2}\,. \tag{3.45}$$

Note that there is no need to worry about R-axions associated with the breaking of R; like the η', the would-be Goldstone can get a mass from its anomaly with the hidden sector gauge group.

In order to be able to make a prediction for these masses, we need to know the particle content and Λ scale for the hidden sector. A natural assumption is that the hidden group is a unified group like SU(5) or SO(10), and that the value of the coupling at the GUT scale is equal to the SSM unified coupling, with $\alpha_{GUT} = 1/33$ for scalars near $\sim 10^{13}$ GeV. If we take SU(5) with a single $(5 + \bar{5})$ in the hidden sector, then $m_{1/2}$ comes out to be too small, about 10^{-3} GeV. However if we use SO(10), then $m_{1/2}$ is naturally near the weak scale! For SO(10) with N_T 10's, we find $m_{1/2} \sim \mu \sim 1\,\mathrm{TeV}$ for $N_T = 1$, and decreasing as N_T increases, down to $\sim 10\,\mathrm{GeV}$ for $N_T = 8$ or equivalently a single adjoint of SO(10). Once again, making a minimal set of assumptions, the gaugino/Higgsino masses again end up "accidentally" near the weak scale.

3.5 SUSY unification in non-SUSY natural theories

For the readers who continue to pine for natural theories, it is perhaps worth mentioning that it is possible to construct theories with natural electroweak symmetry breaking without low-energy SUSY, but with essentially supersymmetric gauge coupling unification. Let us suppose that SUSY is broken and that now *all* the scalars are heavy, but the gauginos and Higgsinos remain light. Gauge coupling unification will still work well, but another sector is needed for electroweak symmetry breaking. We can imagine triggering this with strong dynamics as in technicolor or composite Higgs models, or via AdS duals [53] of such theories. In order to preserve gauge coupling unification, the EWSB sector must have a global SU(5) or $\mathrm{SU}(3)^3/Z_3$ symmetry, into which the SM is gauged in the usual way.

A sketch of an AdS representation of such an idea is as follows. Consider a slice of AdS with SUSY broken in the bulk. There is an $\mathrm{SU}(3)^3/Z_3$ gauge symmetry in the bulk, broken by boundary conditions to $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ on the Planck and IR branes. We have the gauginos and Higgsinos on the Planck brane. Meanwhile, we have an elementary Higgs doublet on the IR brane, so that in 4D CFT language we have a composite Higgs model (we do this for ease of discussion; such models suffer from some tunings and some extra model-building is needed to preserve custodial SU(2), by preserving the full $\mathrm{SU}(2)_R$ on the IR brane [54], but we ignore these details here). In the bulk, we have 3 copies of the standard trinification $(3, \bar{3}, 1)+$ cyclic multiplets, which contain (in SO(10) language) the SM 16 together with an additional $10 + 1$. We can decouple the extra $10 + 1$'s by marrying them off with elementary fermions on the Planck brane, and the SM Yukawa couplings can be generated by writing down Yukawa couplings between the appropriate components of the bulk fermion and Higgs localized on the IR brane; the fermion mass hierarchy can be generated by giving the different generation fermions different bulk mass terms, which localize them by varying amounts to the IR brane, this also avoids large FCNC's [55, 54]. Note that $\mathrm{SU}(3)^3$ was chosen instead of SU(5) since in the SU(5) case, bulk X/Y gauge boson exchange would give rise to unacceptably large rates for proton decay.

It would be interesting to flesh out this construction. The gluino lifetime will continue to be large, however, one would never know whether the high-energy theory is really supersymmetric: while the Higgsinos and gauginos survive to low energies, the scalar Higgs does not, and there are therefore no dimensionless couplings that bear an imprint of the high-scale supersymmetry, unlike the finely tuned examples that have been the focus of this paper.

4. Open problems

There are a number of computations of immediate phenomenological importance in the high-scale supersymmetry scenario we have outlined in this paper.

- **Higgs Mass Prediction.** It is important do a full analysis for the Higgs mass in this model; we have included part of the full 1-loop running and the largest effects from threshold corrections (most notably the 1-loop QCD correction to the top mass), but a fully systematic analysis inclduing 2-loop running and 1-loop threshold corrections is needed.

- **Gluino Phenomenology.** The gluino is perhaps the most important particle of this framework, as its lifetime is a direct probe of the SUSY breaking scale. Moreover, the LHC can be a gluino factory, and therefore an ideal place to study its properties.

Understanding the gluino energy loss and looking for charged tracks, intermittent tracks, displaced vertices and, especially, delayed off-time decays, can help us measure the gluino lifetime and open a window into scales of supersymmetry breaking as high as 10^{13} GeV.

- **Higgsino-Gaugino Phenomenology.** Winos and higgsinos will be produced through Drell-Yan — and not, as typical in the SSM, by squark/gluino production followed by a cascade of decays down to the LSP. It is important to investigate how accurately we can measure the Higgs-Higgsino-Gaugino Yukawa couplings at colliders, possibly at the LHC but more likely at a linear collider. These measurements extrapolated to high energies can give striking evidence for high-scale SUSY.

- **Two-Loop Corrections to Unification.** As we have seen, our 1-loop prediction for $\alpha_3(M_Z)$ is somewhat lower than in the SSM; it is important to perform the full two loop analysis, as this will likely push up our $\alpha_3(M_Z)$ into better agreement with experiment.

- **Dark Matter Detection and Abundance.** Because of the absence of scalars, the collision and annihilation cross sections of the lightest neutralino depend on fewer (than in the SSM) parameters. So, a proper computation of these processes is important, as it can help pin down the interesting parameter ranges in our low-energy theory. It may also predict more precisely the DM detection cross sections.

- **Gluino Cosmology.** As we argued, the gluinos can undergo a second stage of annihilation around the QCD phase transition that further depletes their abundance relative to the standard perturbative freeze-out calculation. It is important to understand this in detail, as this can determine the allowed ranges for the gluino lifetime, and thereby affect the allowed masses for the heavy scalars.

5. Travel guide to a finely-tuned world

Although the cosmological constant problem casts a giant shadow on the principle of naturalness, the prevailing view has been that the LHC will reveal a natural theory for electroweak symmetry breaking, and that gauge coupling unification favors this to be *low-energy* SUSY, despite its nagging problems and the accompanying epicyclic model-building needed to address them.

Here we have outlined an alternate viewpoint, where the usual problems of SUSY vanish, unification is evidence for *high-energy* SUSY, and where accelerators can convincingly demonstrate the presence of fine tuning in the electroweak sector.

The first sign of this proposal at the LHC should be the Higgs, in the mass range of $\sim 120 - 150\,\text{GeV}$. No other scalar should be present, since it would indicate a second, needless, fine-tuning. Next will be the gluino, whose long lifetime will be crucial evidence that the scale of supersymmetry breaking is too large for the hierarchy problem, and a fine-tuning is at work. A measurement of the gluino lifetime can yield an estimate for the large SUSY breaking scale m_S. Next will come the electroweak gauginos and higgsinos, whose presence will complete the picture, and give supporting evidence that the colored octets of the previous sentence are indeed the gluinos. Further precise measurements of the gaugino-higgsino-higgs couplings, presumably at a linear collider, will accurately determine m_S and provide several unambiguous quantitative cross-checks for high-scale supersymmetry.

If this scenario is confirmed experimentally, it will be a striking blow against naturalness, providing sharp evidence for the existence of supersymmetry in nature, as may have been expected for a consistent UV theory of gravity, but not at low enough scales to solve either the hierarchy or cosmological constant problems. This will strongly point to a very different set of ideas to explain these fine-tunings — such as the "galactic" and "atomic" principles, selecting the vacuum of our finely tuned world from a small neighborhood in a landscape of vacua. This may be the closest we will ever come to direct experimental evidence for this vast landscape.

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