THEORETICAL PREDICTIONS FOR COLLIDER SEARCHES

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I review recent developments in extensions of the Standard Model that address the question of electroweak symmetry breaking and discuss how these theories can be tested at future colliders.

In their search for understanding the fundamental laws of elementary particles, high energy physicists are eager to prove that one of the most successful and elegant scientific theories ever formulated – the Standard Model (SM) – actually fails at distances smaller than about a hundred zeptometers (10^{-19} m). A generation of high energy colliders, culminating with the LHC under construction at CERN, has been designed to achieve this goal. Here I will briefly review the present status of the theoretical speculations for new physics within reach of the LHC. In the spirit of this conference, and because of limited space, I will consider only topics in which, in my opinion, there have been recent developments, leaving aside other possibilities which, although interesting, have known only limited progress. The list of references is also largely incomplete.

1. "Big" and "Little" Hierarchy Problems

The hierarchy problem[1] is not necessarily the most puzzling open question of the SM, but it is certainly the most relevant to collider experiments, since its resolution lies – most likely – in the TeV energy range. Its formulation is well known. Treating the SM as an effective theory valid up to a scale $\Lambda_{\rm SM}$, and cutting off momenta in loop integrals at the same scale, we find that the dominant radiative correction to the Higgs mass is given by

$$\delta m_H^2 = \frac{3G_F}{4\sqrt{2}\pi^2}\left(2m_W^2+m_Z^2+m_H^2-4m_t^2\right)\Lambda_{\rm SM}^2$$
$$= -\left(200 \text{ GeV}\frac{\Lambda_{\rm SM}}{0.7 \text{ TeV}}\right)^2 . \qquad (1)$$

The request of no fine-tuning between the tree-level and one-loop contributions to m_H implies $\Lambda_{\rm SM} \lesssim$ TeV. In other words, the SM cannot be valid beyond the TeV, and new physics should appear to modify the ultraviolet behavior. I will refer to this result as the "big" hierarchy, since other fundamental scales, such as the Planck mass $M_{\rm Pl}$, are known to be much larger than $\Lambda_{\rm SM}$.

Table 1. 90% CL limits on the scale $\Lambda_{\rm LH}$ (in TeV) of dimension-six operators $\mathcal{O}$ in the effective Lagrangian $\mathcal{L} = \pm\Lambda_{\rm LH}^{-2}\mathcal{O}$, in both cases of constructive and destructive interference with the SM contribution (±). The limits on the operators relevant to LEP1 are derived under the assumption of a light Higgs.

		−	+
LEP1[3]	$H^\dagger\tau^a H W^a_{\mu\nu}B^{\mu\nu}$	10	9.7
	$\lvert H^\dagger D_\mu H\rvert^2$	5.6	4.6
	$iH^\dagger D_\mu H \bar{L}\gamma^\mu L$	9.2	7.3
LEP2[4]	$\bar{e}\gamma_\mu e\bar{\ell}\gamma^\mu\ell$	6.1	4.5
	$\bar{e}\gamma_\mu\gamma_5 e\bar{b}\gamma^\mu\gamma_5 b$	4.3	3.2
MFV[2]	$\frac{1}{2}(\bar{q}_L\lambda_u\lambda_u^\dagger\gamma_\mu q)^2$	6.4	5.0
	$H^\dagger\bar{d}_R\lambda_d\lambda_u\lambda_u^\dagger\sigma_{\mu\nu}q_L F^{\mu\nu}$	9.3	12.4

As a word of caution, I should recall that, if I apply the same naturalness argument to the quartic divergences, then the present upper bound on the cosmological constant implies that the corresponding cut-off has to be smaller than 10^{-3} eV. Although we cannot fully rule out the possibility that the ultraviolet behavior of gravity is prematurely modified at about 10^{-3} eV, this embarrassing result undoubtedly casts a grievous shadow over the hierarchy problem.

Let me come back to the SM as an effective theory. Unknown new physics at the cut-off is parametrized in terms of non-renormalizable operators. To be most conservative, I will add only operators that preserve all local and global SM symmetries and satisfy the criterion of minimal flavor violation.[2] As shown in Table 1, typical limits on $\Lambda_{\rm LH}$, defined as the effective scale of the new dimension-six operators ($\mathcal{L} = \pm\Lambda_{\rm LH}^{-2}\mathcal{O}$), are $\Lambda_{\rm LH} > 5$–10 TeV.

I will refer to the "little" hierarchy (LH) problem as the tension between the constraints on $\Lambda_{\rm LH}$ – which imply that new-physics virtual effects could only emerge at energies larger than 5–10 TeV – and the no-fine-tuning condition – which requires the presence of new dynamics at the scale $\Lambda_{\rm SM}$, below the TeV. To some readers this could seem like a marginal problem, but I believe it provides a very useful guideline in the search for the correct theory beyond the SM. It is a typical LEP heritage, where any new-physics theory has to be confronted with very precise data.

The "little" hierarchy between $\Lambda_{\rm SM}$ and $\Lambda_{\rm LH}$ is already giving us some lessons on the hypothetical theory beyond the SM. First of all, new physics at $\Lambda_{\rm SM}$ is most likely weakly-interacting, or else effective operators would be generated with $\Lambda_{\rm LH} \simeq \Lambda_{\rm SM}$. Strong dynamics, if any, can only appear at scales larger than $\Lambda_{\rm LH}$. Moreover, even for weak dynamics, physics at $\Lambda_{\rm SM}$ should not induce (sizable) tree-level contributions to effective operators. Recently there has been growing interest in constructing new theories that describe the physics between $\Lambda_{\rm SM}$ and $\Lambda_{\rm LH}$, as I will report in Secs. 4, 5, and 6.

2. Supersymmetry

Supersymmetry[5] provides an elegant solution to the "big" hierarchy problem, by eliminating quadratic divergences with a symmetry principle. In softly-broken supersymmetry, the superpartner masses (generated by gauge-invariant terms) effectively play the rôle of $\Lambda_{\rm SM}$. Because of the absence of quadratic divergences, the validity of the theory can be extended to energy scales as large as $M_{\rm Pl}$, without encountering naturalness problems. This extension has several welcome consequences. It offers the possibility to link the SM to speculative ideas about quantum gravity at $M_{\rm Pl}$. It provides a viable setting for GUT's (with successful unification of gauge coupling constants), a framework for neutrino masses, and for suppressing proton decay. Finally, it allows an implementation of cosmological mechanisms, such as inflation or baryogenesis, which benefit from extrapolations of particle physics models to high energies.

Incidentally, the need for an extrapolation of the SM up to super-heavy scales has been challenged by recent research aiming at constructing new models whose validity cut-off is not very far from the electroweak scale. In these scenarios, at first sight, the advantages of the energy "desert" are lost. Nevertheless, theoreticians have suggested new ways of recovering some of the positive aspects. Just as an example, let me consider gauge coupling unification. In theories with extra dimensions, Kaluza–Klein excitations of SM particles can accelerate the running, possibly leading to a precocious gauge unification,[6] although predictivity is lost in the power running. Alternatively, one can consider an N-fold replication of the SM gauge group, where the unification scale becomes $10^{13/N}$ TeV.[7] Yet, one can abandon the usual tree-level expression for the weak mixing angle in GUT's: $\sin^2\theta_W = {\rm Tr}\ I_3^2/{\rm Tr}\ Q^2 = 3/8$ (where the trace is over any GUT irreducible representation). Taking an electroweak group $SU_3 \times SU_2 \times U_1$ broken to $SU_2 \times U_1$, in the limit $\tilde g_2, \tilde g_1 \gg \tilde g_3$ (where $\tilde g_i$ are the coupling constants of the high energy group), one finds $\sin^2\theta_W = 1/4$.[8] This is much closer to the experimental value $\sin^2\theta_W = 0.231$ than 3/8, and therefore little running is needed.

Aside from this connection to super-high energies, supersymmetry has several interesting and successful features in the low-energy domain. *(i)* Gauge-coupling unification leads to a prediction of a SM parameter: $\alpha_s(M_Z) = 0.124$ (for typical supersymmetric thresholds at 1 TeV) and 0.130 (for thresholds at 250 GeV). Given the intrinsic uncertainty of GUT thresholds, this is consistent with the PDG value $\alpha_s(M_Z) = 0.117 \pm 0.002$.[9] *(ii)* Electroweak symmetry is triggered radiatively by the top Yukawa interaction, with the broken group dynamically chosen (color SU_3 is preserved since squarks do not develop negative square masses, while weak SU_2 is broken). *(iii)* The Higgs boson is predicted to be lighter than about 130 GeV, compatibly with the indications from electroweak precision data. *(iv)* The "little" hierarchy is satisfied, if one assumes R-parity conservation. In this case, supersymmetric virtual effects are loop-suppressed and $\Lambda_{\rm LH} \sim 4\pi\Lambda_{\rm SM}$. *(v)* There is a natural candidate for cold dark matter, if R-parity is conserved.

These positive aspects of low-energy supersymmetry are well known, but let me spend a few words on the negative aspects. The first obvious drawback of low-energy supersymmetry is that no superpartner has been observed. This is in contradiction with the requirement of less than 10% fine-tuning, which predicted a discovery at LEP2.[11] In conventional models of low-energy supersymmetry, the most strin-

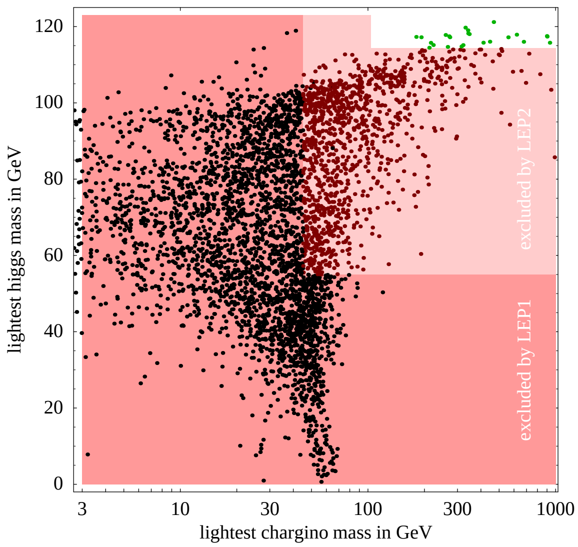

Figure 1. Scatter plot of the chargino and Higgs mass obtained by sampling the parameter space of minimal supergravity. The regions excluded by LEP1 and LEP2 searches are shown. This figure was produced by A. Strumia, updating earlier results.[10]

gent direct bounds come from chargino ($m_{\chi^+} > 103.5$ GeV),[12] charged slepton ($m_{\tilde{e}_R} > 99.4$ GeV),[13] and Higgs searches. The SM Higgs mass bound of 114.4 GeV[14] applies only in the limit of a heavy pseudoscalar, while the general bound is 91.0 GeV[15] (the final combined LEP analysis on the supersymmetric Higgs has not appeared yet). These limits imply a fine-tuning of the order of the per cent. This is illustrated in Fig. 1, which shows the values of Higgs and chargino masses obtained by a random sample of supersymmetric models with universal boundary conditions at the GUT scale and correct electroweak breaking.[10] Less than 1% of the points have simultaneously $m_{\chi^+} > 103$ GeV and $m_H > 114$ GeV (for pseudoscalar Higgs masses close to m_H, the bound becomes weaker but the allowed region of Fig. 1 remains essentially the same).

This is a very disappointing result of supersymmetry, leading to three possible explanations. *(i)* Supersymmetry is not realized at the weak scale. *(ii)* Low-energy supersymmetry is correct, but parameters are fine-tuned at the level of the per cent (and LHC will discover the superpartners). In this case, you may be relieved to know that there is one example in nature of an accidental correlation of this order of magnitude. The apparent sizes of the Sun and the Moon in the sky are roughly equal. If you applied the naturalness criterion to conclude that their actual sizes are roughly equal too, you would be quite wrong. Their radius and distance to the Earth are "fine-tuned" to give, on average, $(\theta_{\rm Sun} - \theta_{\rm Moon})/\theta_{\rm Sun} \sim 3\%$, where $\theta_{\rm Sun,Moon}$ are the solar and lunar angular sizes. *(ii)* Low-energy supersymmetry exists, but its realization is different from what we had imagined. Let me expand on this.

The origin of the fine-tuning is linked to the success of supersymmetry to trigger the electroweak breaking. Indeed, from the Higgs mechanism we know that $M_Z^2 = -g^2\mu_H^2/\lambda$, where μ_H^2 and λ are the coefficients of the quadratic and quartic terms in the Higgs potential. In supersymmetry, $\lambda \sim g^2$, and therefore $M_Z^2 \sim -\mu_H^2$. The effective μ_H^2 has a tree-level direct contribution from the soft terms and a one-loop negative contribution roughly proportional to the stop mass (which, in turn, is usually determined by the contribution proportional to the gluino mass). In the conventional supergravity approach, the one-loop contribution has a large logarithm and it overcomes the tree-level effect. Therefore $|\mu_H|$ is of the order of the supersymmetric masses, which are expected to lie close to M_Z. One may think that the situation improves in gauge-mediated models,[16] since the logarithm can be made small with sufficiently light messengers. However, this is not the case, because gauge mediation gives a boundary condition for the stop mass at the messenger scale, which is significantly larger than the tree-level contribution to μ_H^2. Actually, in gauge mediation, the fine-tuning is usually more acute, because the right-handed selectron is a factor of about 2 lighter than the Higgs mass parameter, and because of the limit on the visible decay of the lightest neutralino.[10] The situation could improve in models where the stop and gluino masses are not large (with respect to the other soft masses) or in models with a low supersymmetry-breaking scale,[17] or in models where μ_H^2 is truly a one-loop factor smaller than all supersymmetric masses (see for instance ref.[18]).

The second drawback of low-energy supersymmetry is the lack of predictivity in the supersymmetry-breaking sector, which translates into the "supersymmetric flavor problem" (i.e. the large flavor violations present with generic soft terms). We are now aware of various schemes, alternative to the conventional supergravity scenario, which are more predictive and address the flavor problem: gauge mediation,[16] anomaly mediation,[19] gaugino mediation.[20] However, at present, we cannot say that one scheme is preferable to the others. For instance, in operator-based schemes (like supergravity) the Higgs mixing mass μ can be properly accounted for,[21] while in other schemes where the soft terms are calculable, the Higgs mass parameters μ and $B\mu$ typically appear at an undesired order in perturbation theory. Luckily, the different schemes usually have quite distinct features both in the mass spectrum and in the nature of the lightest supersymmetric particles, leading to distinct signals at colliders. Therefore, in the case of a discovery, the question of the correct scheme of supersymmetry-breaking terms can be settled experimentally.

Recently many new models have been proposed, where supersymmetry at the weak scale is implemented in unconventional ways. These proposals make use of new ingredients, coming from extra dimensions. Before discussing them, let me first introduce extra dimensions.

3. Extra Dimensions

The hierarchy problem can motivate the existence of experimentally-accessible extra dimensions.[22,23] The scenarios are by now well known: our three-dimensional space is embedded in a larger D-dimensional space-time. While SM fields are confined on the three-dimensional brane, gravity is described by the geometry of the full space. The crucial assumption is that the fundamental Planck scale of the D-dimensional theory M_D is of the order of the weak scale, avoiding any hierarchy. Any distance scale shorter than TeV^{-1} (associated with Newton's constant, with GUT's, right-handed neutrinos, etc.) has to be explained by some geometrical property.

For flat and compact extra dimensions, the observed weakness of gravity is translated into a large value of the compactification radius R, since the ordinary Planck mass is given by[22] $M_{\rm Pl} \sim R^{\delta/2} M_D^{1+\delta/2}$, where $D = 4 + \delta$. The interpretation is that we observe a weak gravitational force not because gravity is intrinsically weak ($M_D \sim$ TeV), but because of the small overlap between our world and the graviton wave function, which is spread in a very large volume.

For spaces with non-factorizable metrics, the large hierarchy $M_{\rm Pl}/M_W$ can be reproduced with much smaller values of R, exploiting the strong functional dependence in the warp factor. In this case, the interpretation is that, because of the non-trivial geometry, the zero-mode graviton wave function is peaked on a brane far from ours and only its exponential tail overlaps with our world, leading to the observed weakness of gravity. The hierarchy is the result of the familiar gravitational redshift in general relativity. A photon that climbs out of a gravitational potential sees its measured energy reduced. It is known that for time-independent metrics with $g_{0j} = 0$, the product $E\sqrt{|g_{00}|}$ is conserved, rather than the energy E. Therefore, a photon emitted at a distance r from a spherically symmetric gravitational source is observed at infinity to be redshifted by an amount $(E_{obs} - E_{em})/E_{em} = \sqrt{|g_{00}|} - 1 \simeq -G_N M/r$, since the Schwarzschild metrics generated by a mass M has $g_{00} = 1 - 2G_N M/r$. Similarly, in the Randall–Sundrum set-up, masses in the hidden (Planck) brane are blueshifted by a warp factor, when measured on the observable (TeV) brane, as a result of the non-trivial gravitational field.

Although the scenarios with a quantum-gravity scale at the TeV have various theoretical and cosmological difficulties, they represent a wonderful possibility for new-physics searches at future colliders. At the LHC, graviton emission is observed in jet plus missing-energy events[24,25] (for flat dimensions) or as resonances in Drell–Yan[26] (for warped dimensions). Tree-level graviton exchange leads to an effective dimension-8 operator $T_{\mu\nu}T^{\mu\nu}$, which predicts correlated signals in diphoton and dilepton final states.[24,27] Graviton loops give rise to dimension-6 operators, with a flavor-universal axial-vector contact interaction playing a special rôle.[28] Because of the lower dimensionality of the induced operators, loop effects are generally more important than tree-level effects, unless the short-distance behavior of gravity is modified at an energy scale significantly smaller than M_D. The radion, a scalar field contained in the extra-dimensional graviton, can mix with the Higgs boson, giving rise to interesting variations in the Higgs searches.[29]

The theoretical description of the experimental signals described above relies upon linearized gravity, whose validity ends as the relevant energy of the process ($\sqrt{s}$) approaches the fundamental gravity scale

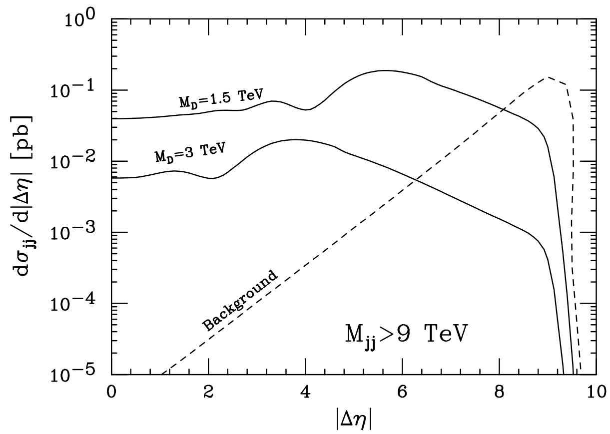

Figure 2. The di-jet differential cross section $d\sigma_{jj}/d|\Delta\eta|$ versus the jet rapidity separation $\Delta\eta$ at the LHC from transplanckian gravitational scattering for $\delta = 6$, $M_D = 1.5$ TeV and 3 TeV. Cuts are chosen such that the di-jet invariant mass $M_{jj} > 9$ TeV, and both jets have $|\eta| < 5$ and $p_T > 100$ GeV. The dashed line is the expected background from QCD.

M_D, and we enter the region where experiments are directly sensitive to the underlying quantum-gravity theory. However, it is interesting that there exists another kinematical regime where the theoretical description is tractable. This is the transplanckian region $\sqrt{s} \gg M_D$, where a semi-classical description is appropriate. This can be understood by noticing that, in the transplanckian region, the classical length associated with a gravitational scattering – the Schwarzschild radius $R_S \sim (G_D\sqrt{s})^{1/(\delta+1)}$, where $G_D \sim M_D^{-(\delta+2)}$ is the D-dimensional Newton constant – becomes larger than the Planck length M_D^{-1} associated to quantum-gravity effects.

When the collision impact parameter b is larger than R_S, the non-linearity of the gravitational field is negligible, and the gravitational scattering, although non-perturbative, can be theoretically calculated. This leads to a prediction at hadron colliders for two-jet final states at large invariant mass and large rapidity separation with characteristic distributions[30] (an example is shown in Fig. 2). The peculiarity in the rapidity distribution comes from a diffractive pattern associated with a new scale in the gravitational potential $b_c \sim (G_D s)^{1/\delta}$, which exists in any dimensions different from 4 (i.e. $\delta \neq 0$) – the familiar Coulomb-like potential is actually very special. These features can be used to confirm the gravitational nature of possible discoveries of new physics.

As the impact parameter b approaches the Schwarzschild radius R_S, the gravitational field becomes strong, and analytic calculations cannot be done. However, there are solid indications[31] that, when $b < R_S$, black holes are formed.[32] This opens the exciting possibility that (unstable) black holes can be observed at future colliders, with spectacular experimental signals.

I should remark that given the present bounds on M_D from direct searches, and the "limited" energy available at the LHC, the window for observing transplanckian phenomena is rather narrow and, more importantly, the separation between classical and quantum scale may not be sufficient to avoid pollution from quantum-gravity effects. Although it is of course important to pursue these searches at the LHC, transplanckian phenomena can offer a good motivation to consider future hadron colliders, such as the VLHC, operating at $\sqrt{s} \sim 100$–200 TeV.

As we have considered extra-dimensional gravity, it is possible that also SM particles live in (part of) the extra-dimensional compactified space.[33] Collider experiments can look for their Kaluza–Klein (KK) excitations. If only gauge bosons live in five dimensions, while fermions are conventional, a combination of present direct and indirect searches yield a 95% CL lower limit on the compactification scale M_C (the mass of the first excited mode) of 6.8 TeV.[34] Searches at the LHC have a sensitivity reach of up to 13–15 TeV. Much weaker bounds ($M_C > 0.3$ TeV) apply to the case of universal extra dimensions,[35] where all SM fields live in five dimensions. This is because, after compactification, the momentum conservation along the fifth dimension corresponds to a conserved KK number. This plays a rôle analogous to R-parity conservation in supersymmetry. Indeed KK particles can only be pair-produced and cannot contribute to SM processes by virtual tree-level exchange. Moreover, the lightest KK mode, corresponding to the hypercharge gauge boson, is stable and is a possible candidate for cold dark matter.[36]

4. Supersymmetry Breaking and Extra Dimensions

At the end of Sec. 2, I anticipated that extra dimensions bring new elements in the construction of low-energy supersymmetric models. Let me briefly discuss these new ingredients, which are related to the issues of supersymmetry and gauge-symmetry breaking, orbifold projection, and the AdS/CFT correspondence.

Consider a field Φ, which lives in a five-dimensional space with the fifth dimension y compactified on a circle of radius R. Because of periodicity, the field can be expanded in discrete Fourier modes in y, which correspond to the familiar KK tower of four-dimensional fields with masses $m_n = n/R$, for integers n. Let us now impose the non-trivial boundary condition that the field Φ picks up a phase as it goes around the circle:

$$\Phi(x, y + 2\pi R) = e^{2\pi i Q_\Phi} \Phi(x, y). \qquad (2)$$

The KK expansion consistent with condition (2) is

$$\Phi(x, y) = e^{iQ_\Phi y/R} \sum_{n=-\infty}^{+\infty} e^{iny/R} \Phi_n(x). \qquad (3)$$

The mass of the n-th mode (i.e. its momentum along the fifth coordinate) is shifted by the boundary condition to the value $m_n = (n+Q_\Phi)/R$. Since the mass spectrum of different fields can be shifted in different ways, symmetries of the five-dimensional theory can be hidden if we look at its four-dimensional truncated version, obtained by keeping only light modes. If the phase Q_Φ corresponds to an R-charge (i.e. if particles inside the same supermultiplet have different boundary conditions), then supersymmetry is broken in the four-dimensional effective low-energy theory. This is the Scherk–Schwarz symmetry-breaking mechanism through boundary conditions in the extra dimensions.[37] Notice that the mechanism is intrinsically non-local, since it involves the global structure of the compactified space. Therefore, if we look at small distances (smaller than R), any symmetry-breaking effect should disappear. This is a very useful property. For instance, in order to stabilize the weak scale it is essential to maintain the cancellation of quadratic divergences in the ultraviolet, and therefore we want to recover supersymmetry at short distances.

The compactification of the fifth dimension into a circle was obtained by identifying the point y with $y+2\pi R$. We can further restrict the space by identifying y with $-y$: the circle collapses into a segment. The KK decomposition (take Eq. (3) with $Q_\Phi = 0$ for simplicity) can be written as $\Phi = \Phi^{(+)} + i\Phi^{(-)}$, where

$$\Phi^{(+)}(x, y) = \sum_{n=0}^{\infty} [\Phi_n(x) + \Phi_{-n}(x)] \cos(ny/R)$$

$$\Phi^{(-)}(x, y) = \sum_{n=1}^{\infty} [\Phi_n(x) - \Phi_{-n}(x)] \sin(ny/R). \qquad (4)$$

The fields $\Phi^{(+)}$ and $\Phi^{(-)}$ are even and odd, respectively, under the symmetry $y \to -y$. However, only the even tower has a massless zero mode ($n = 0$). This projection into an orbifold (essentially a manifold with boundaries) is very useful in model building. In particular,[38] starting from vector-like fermion representations in extra dimensions, one can retain only some chiral components in the truncated four-dimensional theory, where all massive modes have been integrated out, and construct realistic models for weak interactions. Combinations of Scherk–Schwarz boundary conditions and orbifold projections give rise to many interesting possibilities for reducing the symmetry and the particle content of extra-dimensional theories.

The next ingredient I want to discuss is the AdS/CFT correspondence,[39] which is the conjecture that properties of conformal field theories in D dimensions are related to those of $(D+1)$-dimensional theories with gravity in anti-de Sitter space. It has also been suggested that the non-compact version of the RS model (RS2)[40] is dual to a four-dimensional strongly-coupled conformal theory with gravity.[41] One can take advantage of these speculations to extract some physical information on the structure of the RS1 set-up with two branes in a slice of AdS_5, discussed in Sec. 3.[42]

The line element of the RS set-up is given by

$$ds^2 = \exp\left(-\frac{2y}{L}\right) dx^\mu dx_\mu + dy^2, \qquad (5)$$

where y is the extra coordinate and L is the AdS radius. From Eq. (5) we observe that a scale transformation $x \to \xi x$ can be reabsorbed by a shift of the coordinate y. This is suggestive of the holographic interpretation of the fifth coordinate as the renormalization scale of the four-dimensional theory. The hidden (Planck) brane corresponds to the ultraviolet cut-off of the conformal theory and to the addition of gravity. Indeed, gravity decouples from the four-dimensional theory as the Planck brane is moved to infinity. The visible (TeV) brane is the infrared cut-off, where the observable fields live. It corresponds to a spontaneous breaking of the conformal symmetry. Local gauge symmetries in the bulk of the five-dimensional theory are matched to global symmetries of the conformal theory. This connection between higher-dimensional and quasi-conformal theories could be very useful for gaining new insights on some properties of the strongly-coupled regime, which can be accessed by perturbative calculations in the dual theory.

The ingredients presented above have been used by many authors to construct new implementations of supersymmetry at the weak scale. Although I find many of these attempts quite interesting, I do not intend to give a review of these proposals here. Instead I will take only two examples, to show the use of the higher-dimensional ingredients.

The first example[43] is based on the extension of the supersymmetric model into five dimensions, one of which is compactified on a circle with two orbifold projections, $S^1/(Z_2 \times Z_2)$. In five dimensions, there are two supersymmetries. Each boundary of the orbifold breaks one supersymmetry and therefore the effective low-energy theory is non-supersymmetric. However, supersymmetry is recovered at distances smaller than the compactification radius R, leaving only a *finite* contribution to the Higgs mass, and predicting[43] $m_H = 127 \pm 10$ GeV. It is remarkable that one is able to compute a SM parameter (most new-physics theories add new parameters instead of predicting the values of SM quantities!). Unfortunately a quadratic sensitivity to the ultraviolet cut-off is actually introduced by a one-loop contribution to the Fayet–Iliopoulos term.[44] Also the contribution to the ρ parameter is too large, unless we tune the coefficient of an ultraviolet-sensitive operator. Finally unitarity is violated at a scale of 1.7 TeV, where an ultraviolet completion is needed. Variations of this model have been proposed,[45] in which some of these problems can be alleviated, in particular by raising the cut-off into the 5–10 TeV region.

In this model we can clearly appreciate the attractive features of Scherk–Schwarz supersymmetry breaking and orbifolding. Note also that, from the point of view of collider searches, these models are completely distinct from the conventional low-energy supersymmetric ones. Indeed, here there is only one Higgs doublet and two superpartners for each SM particle (the remnant of supersymmetry in five dimensions). The lightest supersymmetric particle is a stable (or metastable, if small R-parity violations are included) stop with a mass of about 210 GeV.

The second example[46] I want to present exploits the possibility of coexistence of supersymmetric and non-supersymmetric sectors in extra dimensions.[47] Let us consider a supersymmetric RS1 set-up where the Higgs sector lies on the TeV brane and the other SM multiplets live in the bulk. A source of supersymmetry breaking is localized on the Planck brane. The bulk fields are directly coupled to the supersymmetry-breaking sector and superpartners (squarks, sleptons, and gauginos) acquire masses of the order of $M_{\rm Pl}$: supersymmetry is badly broken. However the Higgs sector can only feel supersymmetry-breaking effects through loops involving brane-to-brane mediation of bulk particles. Since this corresponds to a non-local interaction, the contribution is finite. Moreover, because of the gravitational redshift discussed in Sec. 3, the Higgs sector feels that the supersymmetry-breaking source is a warp factor smaller than $M_{\rm Pl}$. Using the AdS/CFT correspondence, we can interpret the theory in four dimensions. The SM fields are elementary, while the Higgs sector corresponds to bound states of a quasi-conformal field theory. The electroweak scale is generated at the one-loop level, and therefore $\langle H \rangle \sim (1/4\pi) L^{-1}$, where L (assumed to be TeV^{-1}) is the size of the bound states. The loop factor is very welcome to explain the little hierarchy discussed in Sec. 1. The collider phenomenology is again quite distinct from usual low-energy supersymmetry. Only the Higgs has supersymmetric partners (higgsinos), while a tower of CFT bound states will appear at energies of the order of L^{-1}. This interpretation allows a connection of this set-up with models of composite Higgs bosons.[48] Unfortunately, precision electroweak data set a lower bound on L^{-1} of the order of 9 TeV, implying a certain degree of fine-tuning in the generation of the weak scale. I hope this example has clarified how AdS/CFT can be used in model building to extract information in theories with strong interaction and composite Higgs bosons.

In conclusion, extra dimensions have brought in the game new ingredients for low-energy supersymmetric models. Supersymmetry could reveal itself at colliders in a variety of ways: not only with missing energy, as in supergravity models, but also accompanied by leptons, or photons, or stable charged particles (as in gauge mediation[16]) or with nearly-degenerate $\tilde{W}^{\pm,0}$ (as in anomaly mediation[49]), or with new characteristic signals (stable stop, partially supersymmetric spectrum, etc.).

5. Electroweak Breaking and Extra Dimensions

The ingredients that I have discussed in the previous section can also be used to break gauge symmetries, and it is therefore a logical possibility to use extra dimensions to break the electroweak group.

The hierarchy problem is solved by a symmetry that forbids the Higgs mass term. Supersymmetry (in conjunction with chiral symmetry) is an example. Another example could be gauge symmetry $A_\mu \to A_\mu + \partial_\mu \Lambda$ that forbids a mass term $m^2 A_\mu^2$, but the Higgs boson is not a gauge particle. However, this is true in four dimensions, but not necessarily in more dimensions. A gauge boson in $D = 4 + \delta$ dimensions can be decomposed as $A_M = (A_\mu, A_j)$, $j = 1, \ldots, \delta$, where A_j are scalar fields under the Lorentz group. The extra-dimensional components could be interpreted as Higgs bosons.[50,51] This provides a Higgs–gauge unification completely analogous to the old version of the KK unification of electromagnetism and gravity, in which photon and graviton are different components of the same field – the five-dimensional metrics.

Since, unlike the gauge boson, the Higgs does not transform under the gauge group as the adjoint representation, one has to rely on the orbifolding discussed in Sec. 4 to project out unwanted states. More difficult is to generate the Higgs Yukawa and quartic couplings without reintroducing quadratic divergences. This problem is typically harder than the similiar one encountered in Little Higgs theories (discussed in Sec. 6), since the structure of the Higgs interactions is tightly constrained by the embedding in of extra dimensions. Nevertheless, there has been progress towards the construction of realistic models.[52]

Following a different approach, one can try to break the electroweak group by appropriate boundary conditions of the gauge bosons in the extra dimensions, and eliminate the Higgs altogether. Of course, in this case, one expects KK excitations for the W and Z bosons with masses too low to satisfy present collider bounds.[53] By warping the extra dimension, one can lift the masses of the excited KK modes with respect to the zero modes (to be identified with ordinary gauge bosons) and make these particles phenomenologically acceptable,[54,55] opening a path for the construction of realistic models for electroweak breaking with no Higgs (neither fundamental nor composite). Notice that these proposals differ substantially from the case of a Higgsless SM where the gauge symmetry is non-linearly realized. In the latter case unitarity in scattering processes involving longitudinally-polarized W bosons is violated at an energy scale $G_F^{-1/2}$, below the TeV. On the contrary, violation of unitarity in higher-dimensional theories is postponed to energy scales in the 10 TeV region.[56] This is sufficient to allow these theories to pass the test of the little hierarchy (see Sec. 1). Effectively, the tower of gauge boson KK modes (partly) plays the rôle of the Higgs boson in modifying the cross section high-energy behavior.

At present the main obstacle for Higgsless extra-dimensional theories seems to come from electroweak precision data. Contributions to the ρ parameter can be made small by imposing a gauge custodial symmetry in the bulk (which has the holographic interpretation of a global custodial symmetry in the four-dimensional theory, see Sec. 4). However, contributions to the oblique parameter S[57] (or ϵ_3[58]) appear to be positive and unacceptably large.[59]

6. Little Higgs

We have seen that supersymmetry and gauge symmetry are two examples of principles that can be used to forbid a mass term for the Higgs boson and solve the hierarchy model. A third example is a non-linearly realized global symmetry, where the Higgs has to be identified with a Goldstone boson.

Consider a scalar field Φ, transforming under an abelian global symmetry as $\Phi \to e^{ia}\Phi$, and assume that the symmetry is spontaneously broken by $\langle\Phi\rangle = f$. I can parametrize the complex field Φ in terms of two real fields ρ and θ as $\Phi = (\rho + f)e^{i\theta/f}/\sqrt{2}$. The symmetry transformation properties are $\rho \to \rho$ and $\theta \to \theta + a$. If I identify θ (the Goldstone boson) with the Higgs, the symmetry forbids a mass term $m^2\theta^2$.

The non-trivial part of the problem is to generate the gauge, Yukawa, and self-interaction of the Higgs (which are non-derivative couplings and therefore forbidden by the symmetry) without reintroducing quadratic divergences. Some (not entirely successful) attempts to construct realistic models for the Higgs as a pseudo-Goldstone boson were made in the past.[60]

Recently, a much less ambitious program was proposed, which is known under the name of "Little Higgs".[61] The proposal is to solve only the little hierarchy problem; in other words, one searches for a description valid only up to 10 TeV or so. Recall, from Sec. 1, that the big hierarchy problem is formulated in terms of one-loop corrections to the Higgs mass:

$$\delta m_H^2 \sim \frac{G_F}{\pi^2} m_{\rm SM}^2 \Lambda_{\rm SM}^2. \tag{6}$$

If we require this correction not to exceed the masses of the SM particles $m_{\rm SM}$, we obtain

$$\Lambda_{\rm SM} < \frac{\pi}{\sqrt{G_F}} \sim \text{TeV}. \tag{7}$$

Suppose that at the scale $\Lambda_{\rm SM}$ ($<$ TeV) new physics cancels the one-loop power divergences. Then we are left with one-loop logarithmic divergences and two-loop power divergences of the form

$$\delta m_H^2 \sim \frac{G_F^2}{\pi^4} m_{\rm SM}^4 \Lambda^2. \tag{8}$$

The same naturalness argument now implies

$$\Lambda < \frac{\pi^2}{G_F m_{\rm SM}} \sim 10 \text{ TeV}, \tag{9}$$

which is an energy scale of the order of $\Lambda_{\rm LH}$ (defined in Sec. 1). Therefore, cancelling only one-loop divergences is just sufficient to postpone the big hierarchy problem beyond the scale of the little hierarchy problem, thus being consistent with the phenomenological constraints.

In order to achieve the cancellation, at least at one-loop order, we can use the mechanism of "collective breaking". There are two (or more) approximate global symmetries under which the Higgs is a Goldstone boson. Therefore, each symmetry protects the Higgs mass and one requires that no single term in the Lagrangian simultaneously break all these symmetries. More than one term is necessary to generate the Higgs mass, and therefore the corresponding Feynman diagram has more than one loop.

The collective breaking can be achieved with gauge-group replication and, in this sense, the Little Higgs can be interpreted as a model with discrete extra dimensions through deconstruction.[62] The model-building recipe is the following. One starts from a non-linear sigma model of the Goldstone bosons in the coset G/H. The group G has a weakly-gauged subgroup $G_1 \times \ldots \times G_n$, where $n \geq 2$. Each of the gauge groups G_j preserves a different non-linear global symmetry under which the Higgs transforms as a Goldstone boson. The SM group is a subgroup of $G_1 \times \ldots \times G_n$ which breaks all the global symmetries. Divergent contributions to the Higgs mass necessarily involve the gauge couplings from all the G_j groups and are therefore generated only at the n-th loop:

$$\delta m_H^2 \sim \frac{g_1^2}{(4\pi)^2} \cdots \frac{g_n^2}{(4\pi)^2} \Lambda^2. \tag{10}$$

Instead of using gauge-group replication, one can obtain collective breaking by replicating the field content. Consider the example[63] of an $SU(N)$ gauge group with two scalar fields $\Phi_{1,2}$ in the fundamental representation. Take a scalar potential of the form $V(\Phi_1^\dagger\Phi_1, \Phi_2^\dagger\Phi_2)$ such that both $\Phi_{1,2}$ acquire vacuum expectation values, spontaneously breaking the $SU(N)$ gauge symmetry. In the limit in which you turn off the gauge coupling to Φ_1, the theory has a local $SU(N)$ under which only Φ_2 transforms, and a global $SU(N)$ under which only Φ_1 transforms. Both symmetries are spontaneously broken and the spectrum contains a Goldstone boson. In the limit in which you turn off the gauge coupling to Φ_2, the situation is analogous after replacing the indices 1 and 2, and therefore we still find a massless Goldstone boson. This means that power-divergent contributions can only appear in diagrams with gauge couplings to *both* Φ_1 and Φ_2, i.e. at two loops:

$$\delta m_H^2 \sim \frac{g^4}{(4\pi)^4} \Lambda^2. \tag{11}$$

Along these lines, for both gauge-group and field replications, many models have been constructed. They are quite interesting although, admittedly, rather elaborate.[64,63,65] Effectively, the common operative mechanism that leads to the one-loop cancellation of power divergences is the introduction, at the scale $f \sim \Lambda_{\rm SM}$, of new degrees of freedom. One-loop diagrams from each SM particle are compensated by one-loop diagrams involving a new particle, much in the same fashion as supersymmetry cancels quadratic divergences. The peculiarity, however, is that in Little-Higgs theories the cancellation occurs between particles with the same spin. For instance the top-quark contribution to the Higgs mass is cancelled by a loop of a new vector-like charge-2/3 fermion with a dimension-five effective coupling with the Higgs. The W and Z contribution is cancelled by new electroweak triplet and singlet gauge bosons with negative couplings to the Higgs. The content of scalar particles is more model-dependent and, in some cases, new electroweak triplet scalar fields are necessary for the cancellation. All these new degrees of freedom represent the signature of Little-Higgs theories, which can be searched for in future colliders, in particular the LHC.[66]

Tevatron searches for new gauge bosons and LEP precision data (in particular $\Delta\rho$ contributions from the new top-like fermion, from mixing with the new gauge bosons and, possibly, from the scalar triplet) provide significant constraints on Little-Higgs models. In the minimal model of Arkani-Hamed *et al.*,[64] the symmetry-breaking scale f can be, at best, as low as 5 TeV.[67] This leads to a lower bound on the mass of the new top-like quark $m_{t'} > 2\sqrt{2}(m_t/v)f = 14\ \mathrm{TeV}(f/5\ \mathrm{TeV})$. Comparing with the top contribution to the Higgs mass, we find that this implies a fine-tuning of at least one part in a thousand. Variations of the minimal model can significantly reduce the amount of fine-tuning needed.

The validity of the Little-Higgs theories extends up to a region of about 10 TeV, where we expect an embedding into a more fundamental region. The construction of such an ultraviolet completion is still an open and compelling question.[68]

7. Conclusions

During the last few years we have significantly enriched our basket of theories for the electroweak scale, substantially deepened our understanding of them, and constructed many previously-undiscovered variations. Nevertheless, none of the known models is fully satisfying or totally free from fine-tuning. The excessive fine-tuning may well be just a fortuitous accident, or it is a sign that some theoretical ingredient is still missing (or, possibly, that we are on a completely wrong track!)

A great distinction exists between theories with and without a “desert”. I define “desert” the case in which some new dynamics modify the ultraviolet behavior of the Higgs mass parameter below the TeV, but then physics can be extrapolated up to a very large energy scale without inconsistencies. Conventional low-energy supersymmetry is the best known example. As previously discussed, the desert scenario has some indisputable advantages.

- It allows a connection between SM physics and quantum gravity or other speculative theories at very short distances, such as GUT’s or strings.
- It offers the possibility of predicting some SM parameters (α_S, m_b/m_τ) through GUT relations. In particular, gauge-coupling unification in low-energy supersymmetry is certainly the most remarkable information on new physics we have.
- It naturally embeds the existence of a new mass scale Λ, as suggested by the evidence for neutrino masses and for a new dimension-5 operator $(1/\Lambda)\ell\ell HH$ to be included in the SM Lagrangian.
- It allows a natural suppression of fast proton decay and unwanted flavor violations.
- It provides a set-up for interesting cosmological theories (inflation, baryogenesis, etc.), which require an extrapolation of particle physics to very small distances.

However, for each of these points, alternative (and more or less attractive) explanations have been proposed in scenarios with a low cut-off. In my presentation I did not have the time to enter into the various solutions, but at present it is fair to conclude that the verdict for desert versus non-desert is still pending.

The interest in constructing theories that extend the SM at the TeV, but are valid only up to the 10 TeV region, has both experimental and theoretical roots. The experimental LEP data have convinced us that the scale at which new-physics virtual effects appear (Λ_{LH}) cannot be the same as the scale at which the ultraviolet behavior of the Higgs mass is modified (Λ_{SM}). This has forced us to abandon old versions of electroweak-breaking theories with strong dynamics. From a theoretical perspective, extra dimensions have brought new tools into the game and there has been a great urge to apply them at the weak scale. However, extra-dimensional theories are non-renormalizable, and unitarity is violated at an energy of typically few times the mass of the first KK mode. This implies that the cut-off is not far from the electroweak scale and it opens the urgent question of finding acceptable ultraviolet completions. The tools from extra dimensions have also given new hope for constructing theories of composite Higgs bosons and to better understand their properties.

Of course, in both desert and non-desert scenarios, many fundamental questions are postponed to physics at the cut-off, be it $M_{\rm Pl}$, 10 TeV, or 1 TeV. Although theoretically this is equally unsatisfactory, the value of the cut-off makes a big difference from a phenomenological point of view. When I discussed the little hierarchy problem, I restricted myself to operators that satisfy the same symmetries as the SM Lagrangian. Had I introduced operators that arbitrarily violate flavor, CP, lepton or baryon number, I would have obtained much more stringent bounds on the corresponding mass scale. These bounds pose severe constraints on the dynamics of an ultraviolet completion at 10 TeV, while they are irrelevant for physics at $M_{\rm Pl}$.

Finally, I want to point out that the question of desert versus non-desert has important implications for the strategy of future collider planning. In a scenario like conventional low-energy supersymmetry, a multi-TeV linear collider seems the most appropriate next step after the LHC, because precise studies of new-particle masses and properties are going to be of the utmost importance. On the other hand, in non-desert scenarios, the new physics to be discovered at the LHC is just the first shell of a more complicated structure. Moving towards the highest possible energy then becomes a priority, motivating research for hadron colliders in the 100–200 TeV region, like the VLHC.

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