DECONSTRUCTED HIGHER DIMENSIONAL GAUGE THEORIES

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Abstract

The idea of deconstructing higher dimensional gauge theories is reviewed. Its several potential physical applications are also discussed.

Higher dimensional gauge theories offer interesting new tools to understand the roots of the Standard Model. Among other things, gauge symmetries can be broken by compactification process on singular manifolds, e.g. on orbifolds. Moreover compactification on orbifolds is a simple mechanism to generate chirality in four dimensions. Gauge theories on orbifolds may offer new solutions to the hierarchy problem, with the electroweak scale calculated in term of the compactification radius R. Another important virtue of higher dimensional theories is the possibility of localizing wave functions in extra dimensions. This can explain the hierarchy of various physical parameters, e.g. fermion masses, as a result of a small overlap of wave functions localized at different positions in extra dimensions. Also, this may be a tool to solve the doublet-triplet splitting problem in Grand Unified Theories. However, gauge theories in more than four dimensions are non-renormalizable and some quantum problems cannot be addressed in an unambiguous way.

It has recently been demonstrated that the physics of higher dimensional gauge theories can be reproduced in certain four dimensional gauge theories with enlarged gauge symmetry. For example, the correspondence exists between five dimensional gauge theories with the gauge group G and four dimensional gauge theories with the gauge group G replicated N times, $G \times G \times \cdots \times G$. The four dimensional theory is referred to as 'latticized' or 'deconstructed' and can be viewed as a renormalizable completion of the higher dimensional theory. A more general view on deconstruction is that, inspired by higher dimensional gauge theories, one arrives at a class of purely 4d renormalizable gauge theories that offer (and often generalize) similar benefits to those of higher dimensional gauge theories with no need of extra dimensions at all.

Let me first recall the correspondence between the four-dimensional theory constructed in [1,2] and the gauge interactions in 4+1 dimensions. A gauge theory, e.g. with $SU(n)$

gauge symmetry, in 5 dim has three free parameters: the cut-off scale M_s, the dimensionful coupling constant

$$g_0 = 1/\sqrt{M} \equiv g_5/\sqrt{M_s} \tag{1}$$

and the compacitificatin radius R (we assume that $M_s \gg 1/R$). The physics seen in four dimensions is obtained by integrating out the fifth dimension. The simplest compactifications are on a circle S_1 or on an orbifold S_1/Z_2. In the latter case the four components of the vector potential $A_\mu(x_\mu, x_5)$, $\mu = 0,1,2,3$, are even under Z_2 and the fifth component $A_5(x_\mu, x_5)$ is odd (here and in the following gauge group indices are suppressed). Fourier expanding in the discrete fifth component of momentum we define the four-dimensional degrees of freedom (Kaluza-Klein modes): vectors $A^n_\mu(x_\mu)$, $n = 0,1,\ldots$, and scalars $A^n_5(x_\mu)$, $n = 1,\ldots$. In the compactification on orbifold there is no zero mode of A_5 and, actually, $A_5(x_\mu, x_5)$ can be gauged away. If we compactify on a circle, the zero mode of A_5 remains in the physical spectrum (as a scalar in the adjoint representation of the $SU(n)$ group) and the massive vector KK modes are doubled. For the orbifold compactification, in the axial gauge $A_5 \equiv 0$, the effective Lagrangian after integrating over x_5 contains the following well known terms:

zero-mode kinetic term invariant under $SU(n)$ gauge transformations in four dimensions,

$$-\frac{1}{4}F^{(0)}_{\mu\nu}F^{(0)\mu\nu} \tag{2}$$

were

$$F^{(0)}_{\mu\nu} = \partial_\mu A^{(0)}_\nu - \partial_\nu A^{(0)}_\mu + \tilde{g} f A^{(0)}_\mu A^{(0)}_\nu \tag{3}$$

and $\tilde{g} = g_5/\sqrt{M_s R}$, with the coupling g_5 defined by equation (1);

KK kinetic and mass terms

$$\sum_{n=1}^{N}(\partial_\mu A^n_\nu - \partial_\nu A^n_\mu)^2 + \sum_{n=1}^{N}\left(\frac{n\pi}{R}\right)^2 A^n_\mu A^{n\mu} \tag{4}$$

and triple couplings and quartic couplings of the zero mode and K-K modes.

The structure of couplings reflects conservation of the fifth component of momentum. The truncation in the number of KK modes in understood, $n < N$ where $N \approx M_s R$, so that $M_N < M_s$. The effective Lagrangian has four-dimensional $SU(n)$ gauge invariance, with massive KK modes transforming linearly under the adjoint representation of $SU(n)$, but the full gauge invariance of the five-dimensional Lagrangian is lost because of the truncation $n < N$. The theory is manifestly non-renormalizable.

The renormalizable four-dimensional theory that, in its infrared region, generates the interactions described by the truncated Lagrangian is as follows [1,2]. Let us consider the gauge structure

$$SU(n)_0 \times SU(n)_1 \times \cdots \times SU(n)_N \equiv SU^{N+1}(n) \tag{5}$$

where the vector potentials are $A^a_{j\nu}$ and the dimensionless gauge coupling constants are equal for all of the $SU(n)$ symmetries, $g_{4j} \equiv g_4$. We suppose, in addition, that there is a set of scalar fields $\Phi_j (j = 1,\ldots,N)$ (elementary or effective) which transform as $(\bar{n}, n)$ under $SU(n)_{j-1}$ and $SU(N)_j$ groups. We shall call $\Phi_j s$ the link-Higgs fields. Their transformation properties define the links to be nearest neighbors. The Lagrangian reads

$$\mathcal{L} = -\frac{1}{4}\sum_{i=0}^{N} F^a_{i\mu\nu}F^{ai\mu\nu} + \sum_{i=1}^{N}(D_\mu\Phi)^\dagger_i D^\mu\Phi_i + V(\Phi_i) \tag{6}$$

where $D_\mu = \partial_\mu + ig_4\sum_{i=0}^{n} A^a_{i\mu}T^a_i$ (T^a_i are generators of the i th $SU(n)_i$ gauge symmetry), and the potential has full chiral symmetry $SU^N(n)\times SU^N(n)$ and is symmetric under interchange $\Phi_j \to \Phi_i$.
If the dynamics of the scalar sector (described here by the potential $V(\Phi_j)$; the scalars Φ_j can be replaced by technicolour-like condensates [1]) is such that the diagonal components of the scalars Φ_j acquire a vacuum expectation value v, the chiral symmetry of the link-Higgs sector is spontaneously broken $SU^N(n)\times SU^N(n)\to SU^N(n)$ leaving N Nambu-Goldstone bosons, each transforming under the adjoint representation of $SU(n)$. The gauge symmetry $SU^{N+1}(n)$ of the full Lagrangian is also spontaneously broken to the diagonal $SU(n)$ and n vector bosons acquire masses eating up the Nambu-Goldstone bosons. If we want to have a four-dimensional model corresponding to a compactification on a circle, we introduce one additional link-Higgs field Φ_0 transforming as $(n_0, \bar{n}_N)$. The chiral symmetry of the Higgs sector is now $SU^{N+1}(n)\times SU^{N+1}(n)$. Its breaking to $SU^{N+1}(n)$ leaves $N+1$ Nambu-Goldstone bosons and N of them, as before, are eaten up by longitudinal components of N vector bosons. Thus, one massless scalar remains in the spectrum [3], in exact correspondence to the zero mode of A_5 in the five-dimensional theory. The massless scalar transforms under the adjoint representation of the diagonal, unbroken, $SU(n)$ symmetry.

The physics of our renormalizable four-dimensional theory below the scale v of spontaneous gauge symmetry breaking can be easily studied in detail in the nonlinear σ-model approximation. For the orbifold compactification, upon substituting

$$\Phi_j \to v\ \exp(i\phi^a_j T^a/2v) \tag{7}$$

the Φ_j kinetic terms lead to a mass matrix for the gauge fields

$$\sum_{i=1}^{N}\frac{1}{2}g_4^2v^2(A^a_{(i-1)\mu} - A^a_{i\mu})^2 \tag{8}$$

This mass matrix has the structure of a nearest-neighbor coupled oscillator Hamiltonian. We can diagonalize the mass matrix to find the eigenvalues:

$$M_n = \sqrt{2}g_4 v\sin\left[\frac{\gamma_n}{2}\right] \qquad \gamma_n = \frac{n\pi}{N+1} \quad n = 0,1,\ldots,N. \tag{9}$$

Thus we see that for small n this system has a KK tower of masses given by

$$M_n \approx \frac{g_4 v\pi n}{\sqrt{2(N+1)}} \qquad n \ll N \tag{10}$$

and $n=0$ corresponds to the zero mode.
To match on to the spectrum of the KK modes, we require

$$\frac{g_4 v}{\sqrt{2}(N+1)} = \frac{1}{R}. \tag{11}$$

Hence, the system with $SU(n)^{N+1}$ and N Φ_i is provides a description of the KK modes of the 5d theory by generating the same mass spectrum.
Next, it is crucial to examine the interactions from the model. The gauge fields A^j_μ can be expressed as linear combinations of the mass eigenstates $\tilde{A}^n_\mu$ as

$$A^j_\mu = \sum_{n=0}^{N} a_{jn} \tilde{A}^n_\mu. \tag{12}$$

The a_{nj} form a normalized eigenvector ($\vec{a}_n$) associated with the nth $n \neq 0$ eigenvalue and has the following components:

$$a_{nj} = \sqrt{\frac{2}{N+1}} \cos\left(\frac{2j+1}{2}\gamma_n\right) \qquad j = 0, 1, \ldots, N. \tag{13}$$

The eigenvector for the zero mode, $n = 0$ is always $\vec{a}_0 = (1/\sqrt{N+1})(1, 1, \ldots, 1)$. We can now rewrite the Lagrangian equation (6) in the mass eigenstates of the vector bosons ($\tilde{A}^n_\mu$) and derive the interactions between them. They are identical to those of the 5d theory, with the identification

$$\tilde{g} = \frac{g_5}{\sqrt{M_s R}} = \frac{g_4}{\sqrt{(N+1)}}. \tag{14}$$

In the 4d theory, there are three relevant parameters, namely, the gauge coupling constant g_4, the total number of $SU(n)$ groups $N+1$ and the vacuum expectation values v of the Higgs fields v determined by the potential $V(\Phi_i)$. The mapping between them and the parameters of the 5d theory are $N+1 = M_s R, g_4 = g_5 = \sqrt{M_s/M}$ and $v = \sqrt{M_s M}$.

We conclude that the physics of the 4d theory below the scale v of spontaneous gauge symmetry breaking is the same as that of 5d theory with cut-off. That correspondence does exist for supersymmetric gauge theories, too [4]. Suppose we raise the scalar fields in the previous model (the one corresponding to compactification on a circle) to 4 dim, $\mathcal{N}$=1 chiral multiplets and we introduce vector superfields $V_j, j = 1 \ldots N$ associated to the gauge group $SU(n)_j$:

$$\Phi_i(\varphi_i, \psi_i) \qquad \textit{and} \qquad V_j(A^a_j, \lambda^a_j) \tag{15}$$

After the diagonal breaking we can identify 4 dim $\mathcal{N} = 2$ massless vector

$$\begin{pmatrix} & A^\mu & \\ \lambda & & \psi \\ & \varphi & \end{pmatrix} \begin{matrix} \text{2 components} \\ \text{4 components} \\ \text{2 components} \end{matrix} \tag{16}$$

and $N-1$ 4 dim $N = 2$ massive vectors

$$\begin{pmatrix} & A^\mu & \\ \lambda & & \psi \\ & \varphi & \end{pmatrix} \begin{matrix} \text{3 components} \\ \text{4 components} \\ \text{1 components} \end{matrix} \tag{17}$$

in exact correspondence to the zero-mode and K-K modes of 5 dim vector multiplet

$$\begin{pmatrix} A^\mu \\ \psi \\ \phi \end{pmatrix} \begin{matrix} \text{3 components} \\ \text{4 components} \\ \text{1 components} \end{matrix} \tag{18}$$

Thus, the 4d $\mathcal{N}$ =1 $SU(n)^N$ gauge theory is broken at the scale v to 4d $\mathcal{N}$ =2 $SU(n)$ diagonal gauge theory. Similar correspondence exists after including hypermultiplets and, as it will be discussed later on, for chiral theories compactified on S_1/Z_2 orbifold. As another extension of those ideas, it was also shown that the 5d and 4d theories show the same non-perturbative structure [5].

One should stress that the correspondence between the 5d and 4d theories holds only below the scale v of the breaking of the gauge symmetry of the latter to the diagonal subgroup. Generically, the 4d theory contains above the scale v more active degrees of freedom with the masses $m \gtrsim 0(v)$ which do not have any extra dimensional interpretation. For instance, a supersymmetric 'aliphatic' model [4] which mimices in its IR a supersymmetric gauge theory compactified on S_1/Z_2 has chiral anomalies in the first and the N th gauge group caused by the missing of the link superfield $\phi(\bar{n}_N, n_1)$. They can be canceled by adding superfields $\phi(\bar{n})$ and $\phi(n_N)$ in the adjoint representations of the respective gauge groups. They can be given masses $0(v)$ by the superpotential $W = \frac{1}{v^{N-2}}\phi(\bar{n}_1)\phi(n_1, \bar{n}_2)\ldots\phi(n_{N-1}, \bar{n}_N)\phi(n_N)$. (Another possibility to cancel those anomalies is of course, to use the Green-Schwarz mechanism.)

One should also mention that the nature of the correspondence between the extra space dimension and the 'group product space' is, in general, quite subtle and the precise lattice interpretation can be used only in the simplest cases (e.g. for pure gauge theories compactified on a circle). The two classes of the theories should rather be compared at the level of effective theories in four dimensions: the 5d theory- after integrating out the 5 th dimension, and the 4d theory - after breaking the full gauge symmetry to the diagonal subgroup.

We can summarize this discussion as follows: 'Deconstruction' is a prescription for a gauge invariant and renormalizable regularization of $d > 4$ dimensional gauge theory by reduction to 4d. For instance, the power-law like running of the gauge couplings in five dimensions [6,7] can be discussed in the renormalizable setting. In the 4d theory, the running is logarithmic but with the β function changing every time we pass the threshold M_n [2]. In the limit of large N one recovers power-law-like running. However, a more general attitude to deconstructing higher dimensional gauge theories is to consider a class of 4d theories that implement many seemingly extra dimensional ideas, even with small N and including also models with no obvious extra dimensional correspondence. Thus, we can talk about model building in 4d inspired by extra dimensions with extra dimensional benefits. But at the end extra dimensions are not there.

In fact, theories with replicated gauge groups and bi-fundamental matter have already been discussed in the literature in various contexts. The results of ref. [1,2] encourage to their further exploration and suggest they may be a model building tool complementary to higher dimensional gauge theories. Two particularly interesting examples of such theories discussed earlier and related to the recent developments are $SU(5) \times SU(5)$ GUT models [8], and the quiver theories [9] widely discussed in the context of the search for CFT's and of A_dS_5/CFT correspondence.

We shall now discuss briefly a few potentially interesting possibilities offered by theories with replicated gauge groups, inspired by the idea of deconstruction. The Grand Unified Theory based on the $SU(5) \times SU(5)$ group broken spontaneously to the diagonal Standard Model group has been revived in ref. [10], as an solution to the doublet-triplet splitting problem and as an explanation of the hierarchy of the fermion masses. A discrete symmetry in a four-dimensional $SU(5)$ model commutes with $SU(5)$, modulo a standard model gauge transformations. Such a discrete symmetry leaves the doublet mass term $H\tilde{H}$ in the superpotential invariant if and only if the triplet $Q\tilde{Q}$ terms is invariant. So such a discrete symmetry cannot solve the doublet-triplet splitting problem. However, starting with the $SU(5) \times SU(5)$ group it is easy to find unbroken discrete symmetries with the desired properties. The same symmetry can explain the hierarchy of fermion masses. Other solutions to the doublet-triplet splitting problem, also in the framework of theories with replicated GUT groups, have been proposed in ref.[11].

Theories inspired by the deconstruction idea often show interesting structure at the level of quantum corrections and give new insight to the hierarchy problem. In the bottom-up approach to the hierarchy problem one may be tempted to take the following attitude: rather than solving perturbatively the hierarchy problem to a very high scale, postpone 'sensitive' ($\wedge^2$) UV dependence of the electroweak symmetry breaking mechanism, using perturbative physics, to a scale $\wedge \gtrsim 0(10\text{TeV})$. In other words, the question is this: can one solve the hierarchy problem if a non-perturbative cut-off to the Standard Model is at $\wedge \gtrsim 0(10\text{TeV})$ (and neither at $\wedge \sim 0(1\text{ TeV})$ nor at $\wedge \sim M_{PL}$). In practice, one then looks for theories with δm_H^2 finite (or at most logarithmically divergent) at one-loop level, but not necessarily in higher orders of perturbation theory.

One class of models inspired by deconstruction and discussed in the above context are non-supersymmetric models with the Higgs boson appearing as a pseudo- Goldstone boson [12]. This mechanism has its correspondence in higher dimensional gauge theories with a Higgs boson as a gauge boson in extra dimensions. Models based on deconstruction have the virtue of a large collection of symmetries protecting the Higgs mass.

The basic idea underling that approach can be illustrated by the following simple (and totally unrealistic) example. The scalar sector of the discussed earlier $SU(n)^N$ gauge theory with N link-Higgs fields in bi-fundamental representations $\phi(\bar{n}_i, n_{i+1})$ (cyclic case) has $SU(n)^N \times SU(n)^N$ 'chiral symmetry (invariance under independent 'left' and 'right' rotations: $\phi_i \to U_i \phi_i V_i^+$). This symmetry is broken spontaneously by the vevs $\phi_i = v1$ to the diagonal subgroup $SU(n)^N$. The $(N-1)$ Goldstone bosons are ' eaten up' by the gauge fields when $SU(n)^N$ gauge symmetry is broken to the diagonal subgroup $SU(n)$ and one Goldstone boson remains in the physical spectrum. However, the original chiral symmetry of the scalar sector is not only broken spontaneously by $\phi_i = v1$ but it is also broken explicitly by the $SU(n)^N$ gauge interactions. Thus, the Goldstone boson that remains in the physical spectrum is actually a pseudo-Goldstone boson, with non-zero radiative corrections to its mass. But, the gauge interactions which break explicitly the chiral $SU(n)^N \times SU(n)^N$ symmetry of the scalar sector, at the same time, protect the corrections to the Goldstone boson mass against any divergences! Indeed, the only gauge invariant operator that contributes to the Goldstone boson mass is $[Tr\phi_1\phi_2 \ldots \phi_N]^2$ and for $N > 2$ it is dim>4 operator; so if generated radiatively, it must have a finite coefficient.

Of course, identifying the electroweak Higgs boson with that pseudo-Goldstone boson is totally unrealistic. The latter is in the adjoint representation of the gauge group and, moreover, the generated mass is positive. However, more realistic models have been constructed [13] along similar lines, at the expense of introducing more interactions and more symmetries. The Higgs boson mass is no longer fully protected from divergences but one-loop quadratic divergences are absent.

Using higher dimensional supersymmetric gauge theories, there has also been proposed another approach to the "weak" solution to the hierarchy problem in which the Higgs boson mass is calculated in terms of the compactification radius, $\delta m^2 \sim f(R, \wedge)$, with weak dependence on the cut-off scale $\wedge$. In ref. [14] this is achieved by breaking supersymmetry by the Scherk-Schwarz mechanism. The deconstructed theories offer an analogous possibility of calculating the electroweak scale in terms of the deconstruction scale v, in a renormalizable theory [15]. In ref. [15], the $\mathcal{N}$ =1 supersymmetry of the $SU(2)^N$ gauge theory is broken in a hard way, by removing some of the degrees of freedom within supermultiplets. In consequences, the zero mode spectrum is not supersymmetric. Nevertheless, the one-loop contribution to the Higgs boson mass coming from the large top quark Yukawa coupling remains finite and triggers the electroweak symmetry breaking.

Finally, it is worth mentioning that the structure of deconstructed gauge theories resembles the $U(n)^N$ gauge theories with bi-fundamental matter constructed in open string compactifications [9]. So-called quiver theories are constructed by studying a stack of D3 branes on orbifolds of R^6 transverse to the D-branes. Equivalently, they can be obtained by orbifolding $\mathcal{N}$ = 4 supersymmetric $U(Nn)$ gauge theory in 4d by a discrete Z_N group embedded in the $SU(4)$ R-symmetry group of the $\mathcal{N}$ =4 theory. This way one can construct $U(n)^N$ gauge theories with $\mathcal{N}$=2,1 or 0 supersymmetries. The scalar potential of those theories has flat directions. One of them is $\phi_i = v_i 1$ which breaks $U(n)^N \rightarrow U(n)_{\text{diagonal}}$ The stringy interpretation of the diagonal breaking is that branes are moved away from the fixed points of Z_N [16].

The idea of deconstruction provides some phenomenological motivation for quiver theories. The $\mathcal{N}$ =2,1 theories are superconformal so one can contemplate the possibility of embedding realistic low energy gauge theories (Standard Model?) into superconformal (finite) theories with replicated gauge symmetries.

The non-supersymmetric quiver theories are also interesting as they show several interesting properties [17]. In such theories there is equal number of fermions and boson but they are in different gauge group representations. Nevertheless one finds [17] no one-loop quadratic divergences in the effective potential ($STrM^2 = 0$ identically) for any pattern of gauge symmetry breaking (if there are no scalars in the adjoint representation of one of the $(U(n)'s)$. Moreover , in the deconstruction phase, with $\phi_i = v_i 1$ and $U(n)^N \rightarrow U(n)_{\text{diagonal}}$, one can identify "custodial" supersymmetry: all terms in the Lagrangian up to terms quadratic in the heavy modes match the structure of $\mathcal{N}$= 1 supersymmetric theory (the zero mode sector has $\mathcal{N}$= 4 supersymmetry). Therefore it follows that any nth level spectrum is boson-fermion degenerate and the universal vevs remain flat direction at one loop. Indeed, $STrM^{2q} = 0$ and one-loop effective potential is zero. Also, one-loop corrections to the zero mode masses vanish.

Thus, this is an example of a theory with supersymmetry partially restored in the low energy phase, with gauge symmetry broken to its diagonal subgroup. It remains to be seen if this fact may have any interesting implications for the hierarchy problem.

In summary, the deconstruction idea shows that theories with replicated gauge groups broken down the diagonal subgroup are a useful alternative to extra dimensions. Moreover, they offer an interesting link to open string theories.

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