Light Composite Higgs
The Third Way to Electroweak Symmetry Breaking

1. Theoretical Models

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The biggest gap in our understanding of particle physics is the mystery of electroweak symmetry breaking.

We know that SU(2)xU(1) is a symmetry of the couplings of γ, W, Z to quarks and leptons. These couplings are measured to high precision and follow the SU(2)xU(1) relations.

In order to give mass to the quarks and leptons, and to W and Z, the SU(2)xU(1) symmetry must be spontaneously broken.

In the Standard Model, this breaking is accomplished by the fundamental Higgs field $\varphi(x)$.

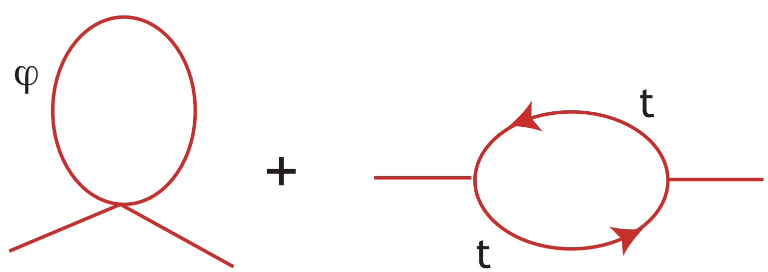
Of course, this field does not explain anything. The Higgs potential is

$$V(|\varphi|) = \mu^2|\varphi|^2 + \lambda|\varphi|^4$$

Spontaneous symmetry breaking occurs because

$$\mu^2 < 0$$

The Standard Model gives no physical explanation of the sign of μ^2. In fact, μ^2 receives divergent radiative corrections with both signs.



$$\mu^2 = \mu^2_{\rm bare} + \frac{\lambda}{8\pi^2}\Lambda^2 - \frac{3y_t^2}{8\pi^2}\Lambda^2 + \cdots$$

We would like to find a theory in which spontaneous symmetry breaking arises naturally from physical principles.

Two well-known examples are technicolor and supersymmetry.

Technicolor is the idea that electroweak symmetry breaking arises from fermion pairing and condensation due to a new strong interaction at TeV energies.

Technicolor has severe phenomenological difficulties:

- awkwardness of coupling techni-fermions to quarks and leptons to generate their masses

- tendency of this coupling to generate new flavor-changing weak interactions

- techni-rho - Z mixing, which generates large shifts of the Z mass and precision electroweak observables

- difficulty of obtaining a large enough top quark mass without making the top itself a technifermion

- for "topcolor-assisted technicolor", prediction of a Higgs boson with mass $m_h \sim 2m_t$

Supersymmetry is often motivated in its own right as a beautiful extention of space-time symmetry.

However, it also provides a solution to the problem of electroweak symmetry breaking:

- SUSY gives a raison d’être for a fundamental scalar field.

- SUSY protects the Higgs mass from large radiative corrections.

- The coupling of the Higgs field to the top squarks generates

$$\mu^2 < 0$$

However, SUSY has been searched for intensively in the first two years of the LHC.   At the moment, all results are negative.

Personally, I am still a believer in supersymmetry.

The models excluded are models in which the squarks and gluino have roughly the same mass scale, with a large gap to the lightest supersymmetric particle.   This is a plausible expectation for the spectrum, but not the only one.

To naturally explain the electroweak symmetry breaking scale, only the top squarks and Higgsinos have to be very light.

A crucial question for 2012 is:

Do there exist top squarks of mass 500 GeV ? These might be directly produced from gluons

$$gg \to \widetilde{t}\widetilde{t}^* \to b\bar{b}\ell^+\ell^- + \mathrm{MET}$$

even if the gluino is above 1 TeV.   In 2012 we should have enough data to discover or exclude this scenario.

Still, it is logical to ask:

If technicolor and supersymmetry are being excluded, what is next ?

If we find a light Higgs and no SUSY, are we stuck with the Standard Model ?

In these lectures, I will describe three related models of electroweak symmetry breaking with light but composite Higgs bosons. These theories are known as

- Little Higgs
- Gauge-Higgs Unification
- Randall-Sundrum / warped extra dimensions

These models are similar in their general characteristics and in their signatures.

Lecture 1 will describe the theoretical framework.
Lecture 2 will discuss LHC signatures and searches.

We might begin with general requirements for a light Higgs boson.

If we accept that SU(2)xU(1) is an exact gauge symmetry of Nature, it is not possible to construct a model without the Higgs mechanism. There is always a resonance in the I = J = 0 partial wave in WW scattering; this is the Higgs boson.

At the moment, the LHC Higgs exclusions leave two possibilities:

1. The Higgs boson is light: $114 < m_h < 130$, with slightly looser bounds in exceptional scenarios

2. The Higgs boson is heavy, $m_h > 600$, and is likely to be a broad resonance

In these lectures, we will follow the path in which the Higgs boson is light. In this case, whether the Higgs is fundamental or composite, there is an effective field theory, valid at the TeV scale and below, in which the Higgs is described by an elementary scalar field.

In this context, we can describe the requirements for a theory of electroweak symmetry breaking in two complementary ways:

1. Large radiative corrections to μ^2 must cancel

We have seen that, in the Standard Model, μ^2 receives divergent radiative corrections from diagrams with t and W, Z as well as from the Higgs self-coupling. Within the effective field theory, these divergences must be cancelled by additional diagrams.

In SUSY, the cancelling diagrams involve $\widehat{t}$ and $\widetilde{\chi}$.

In non-SUSY theories, the cancelling diagrams will include new particles with the same statistics as the Standard Model states. Thus, we require some W', Z', T to be part of these models.

This already raises an issue:

In the Standard Model, all masses are of the form

$$m \sim \lambda v \quad \text{where} \qquad v = 246 \ \ \mathrm{GeV}$$

and λ is a perturbative coupling. This limits masses to be below about 500 GeV.

If we want W’, Z’, T to be heavier, the main part of their masses cannot come from electroweak symmetry breaking.

So these cannot be simple sequential W, Z or 4th generation T. We need a different principle for the mass generation.

2. A physical principle should imply $\mu^2 = 0$ in the leading order of some expansion.

There are two interesting principles that can accomplish this.

The underlying theory by have a spontaneously broken symmetry for which the Higgs doublet fields in φ are Goldstone bosons.

The underlying theory may be higher-dimensional, with gauge bosons A_M^a such that some extra components A_5^a provide the Higgs double fields φ .

I will now describe some models that make these ideas more concrete.

A very nice review of this subject is:

G. Bhattacharyya, arXiv: 0910.5095 , Rept. Prog. Phys. 74, 026201 (2011).

First, explore the idea of Higgs as a Goldstone boson. This gives the Little Higgs models of Arkani-Hamed, Cohen, Katz, Nelson ...

To avoid complex group theory, I will consider a model with global symmetry SU(3)xSU(3) -- like QCD with massless u,d,s. There are 8 Goldstone bosons. As in QCD, we can describe these as parameters of an SU(3) unitary matrix

$$V = e^{2i\Pi^a t^a/f} \qquad\qquad 2i\Pi^a t^a = \begin{pmatrix} \Phi & H \\ -H^\dagger & \phi \end{pmatrix}$$

Note that H is an SU(2) doublet.

We will want $f \sim 1 \text{ TeV} \ , \ M_\rho \sim 10 \text{ TeV}$

All fields in V must be massless if the SU(3) symmetries

$$V \to \Lambda_R V \Lambda_L^\dagger$$

are respected.

To generate the top quark mass, this structure must couple to top. We need to put top quark into the representations

$$\chi_L = \begin{pmatrix} u \\ b \\ U \end{pmatrix}_L \qquad U_R \qquad u_R$$

with an extra singlet quark. The effective Lagrangian is

$$\mathcal{L} = -\lambda_1 f \begin{pmatrix} 0 & 0 & \overline{u}_R \end{pmatrix} V \chi_L - \lambda_2 f \overline{U}_R U_L$$

The first term has the symmetry $V \to V \Lambda_L^\dagger \qquad \chi \to \Lambda_L \chi$

The second term has the symmetry $V \to \Lambda_R V$

Either symmetry suffices to insure that H is exactly massless. Thus, to build a Higgs potential, we need to involve both interaction terms.

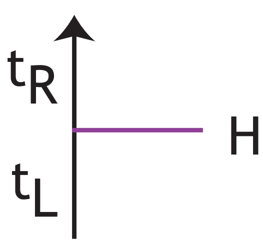
Transform to the top quark mass eigenstates:

$$t_L = u_L \qquad t_R = \frac{\lambda_2 u_R - \lambda_1 U_R}{\sqrt{\lambda_1^2 + \lambda_2^2}}$$

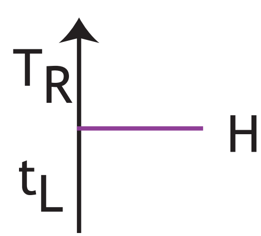
$$T_L = U_L \qquad T_R = \frac{\lambda_1 u_R + \lambda_2 U_R}{\sqrt{\lambda_1^2 + \lambda_2^2}}$$

$$m_T = \sqrt{\lambda_1^2 + \lambda_2^2} f$$

then the H vertices are:



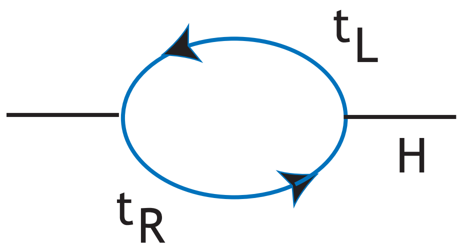
$$= -i\lambda_t$$



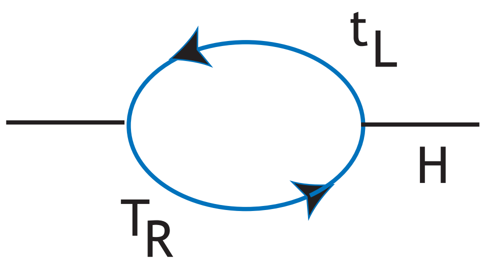
$$= -i\lambda_T$$

$$\lambda_t = \frac{\lambda_1 \lambda_2}{\sqrt{\lambda_1^2 + \lambda_2^2}} \qquad \lambda_T = \frac{\lambda_1^2}{\sqrt{\lambda_1^2 + \lambda_2^2}}$$

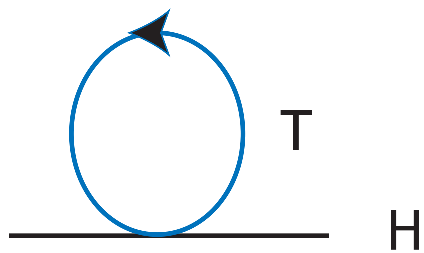
Using this structure, compute corrections to the H mass:



$$= -6\lambda_t^2 \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2}$$



$$= -6\lambda_T^2 \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2 - m_T^2}$$



$$= +6\frac{\lambda_T}{f} \int \frac{d^4k}{(2\pi)^4} \frac{m_T}{k^2 - m_T^2}$$

By the relations on the previous page, the quadratic divergences of these diagrams cancel. What is left over is:

$$m_H^2 = -3\frac{\lambda_1^2\lambda_2^2 f^2}{8\pi^2} \log\frac{M^2}{m_T^2} = -3\frac{\lambda_T^2 m_T^2}{8\pi^2} \log\frac{M^2}{m_T^2}$$

If the top quark is heavy, this set of contributions can dominate the H mass and produce EWSB.

Realistic Little Higgs models have more structure and some complications:

To find an unbroken SU(2)xU(1) symmetry after the SU(3)xSU(3) symmetry breaking, the original model should have had two copies of SU(2)xU(1) and two sets of SU(2)xU(1) gauge bosons. The second set of bosons become massive at the scale

$$m(W') \sim gf \sim \mathrm{TeV}$$

These are the massive bosons that we need to cancel the W,Z radiative correction to m_H^2 .

To obtain $\langle H \rangle \ll f \sim \mathrm{TeV}$, the theory must contain a Higgs quartic coupling of order 1, or at least not suppressed by λ_t^4. This puts group-theoretic restrictions on the symmetry breaking, leading to more complex patterns such as $SU(5) \to SO(5)$. These restrictions have recently been analyzed by Hook and Wacker.

If T and the new W and Z bosons are as light as 500 GeV, they create large perturbations of precision electroweak observables. A way to avoid this is to introduce a T-parity symmetry carried by the lightest new states (Cheng and Low). Then these states are produced only in pairs.

It is possible that the lightest T-parity odd vector boson could be a stable dark matter or MET particle. However, Hill and Hill argued that - generically - this boson decays to $\gamma\gamma$,WW, ZZ .

The other method to construct a theory with a Higgs mass that is naturally zero at leading order is to consider a universe with extra space dimensions. These dimensions are not visible, so they must be very small. I will begin by discussing a single new dimension x^5 with zero curvature and periodic boundary conditions around a circle of circumference $2\pi R$.

There are other possibilities; I will discuss one later.

Let first collect some basic formulae for fields in an extra dimension.

Consider the Klein-Gordon equation in this space:

$$(\partial_M \partial^M + m^2)\phi = 0 \qquad\qquad M = 0, 1, 2, 3,\ 5$$

To study this theory, Fourier-analyze in the 5th coordinate:

$$\phi(x, x^5) = \sum_n \phi_n(x) e^{ik_n x} \qquad\qquad k_n = n/R$$

Then ϕ_n obeys

$$(\partial_\mu \partial^\mu + k_n^2 + m^2)\phi_n = 0$$

so this is an ordinary Klein-Gordon field with a mass

$$M = [(n/R)^2 + m^2]^{1/2}$$

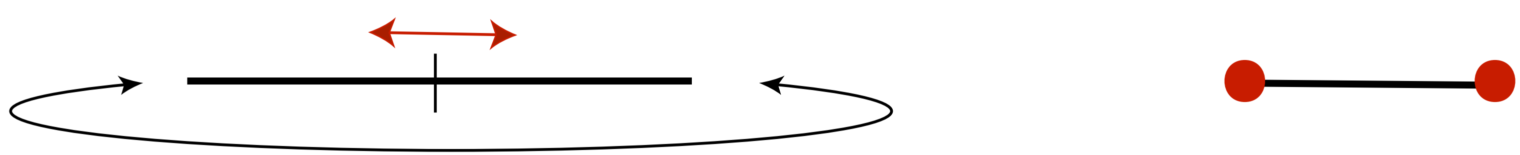
If the field was originally massless, the Fourier modes will be massive fields with the simple mass relation

$$M = |n|/R$$

This is called a Kaluza-Klein reduction, leading to a Kaluza-Klein tower (KK tower) of massive particles.

To lower the symmetry and to build in more structure, it is often convenient to perform an orbifold reduction of the space. That is, we identify a discrete symmetry P and restrict the theory to states invariant under P .

For example, consider a periodically connected interval, and impose the symmetry $P: \quad x^5 \to -x^5$. Effectively, we now have an interval $0 < x^5 < \pi R$.



The boundaries of the interval are fixed points of the symmetry operation. Wavefunctions must be even under reflection across these points.

The fixed points are actually 3+1-dimensional planes in the 4+1 dimensional space. In principle, there can be particles bound to these points. String theory gives illustrations of this.

An interesting case is provided by an SU(3) gauge theory. Consider states invariant under

$$P: \quad x^5 \to -x^5 \quad \text{and} \quad \xi \to \begin{pmatrix} 1 & & \\ & 1 & \\ & & -1 \end{pmatrix} \xi$$

On the multiplet of gauge field, the action is

$$\begin{pmatrix} A_\mu(x^5) & H_\mu(x^5) \\ H_\mu^\dagger(x^5) & a_\mu(x^5) \end{pmatrix} \to \begin{pmatrix} +A_\mu(-x^5) & -H_\mu(-x^5) \\ -H_\mu^\dagger(-x^5) & +a_\mu(-x^5) \end{pmatrix}$$

and

$$\begin{pmatrix} A_5(x^5) & H_5(x^5) \\ H_5^\dagger(x^5) & a_5(x^5) \end{pmatrix} \to \begin{pmatrix} -A_5(-x^5) & +H_5(-x^5) \\ +H_5^\dagger(-x^5) & -a_5(-x^5) \end{pmatrix}$$

Components with +1 have wavefunctions $\cos(nx^5/R)$, and
n= 0 is allowed.

Components with -1 have wavefunctions $\sin(nx^5/R)$, and
n = 0 is forbidden.

The resulting theory has states at n= 0 corresponding to an SU(2)xU(1) multiplet of gauge bosons and a complex SU(2) doublet scalar field.

This is Gauge-Higgs Unification (Hall-Nomura/Hosotani)

The Kaluza-Klein towers contain both gauge bosons and extra scalars.   The extra scalars provide the longitudinal components of the massive gauge bosons.

For the zero mass Dirac equation in 4 dimensions:

$$i\gamma \cdot \partial \Psi = 0$$

we can introduce the representation

$$\gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \overline{\sigma}^\mu & 0 \end{pmatrix} \qquad \sigma = (1, \vec{\sigma}) \qquad \overline{\sigma} = (1, -\vec{\sigma}) \qquad \Psi = \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix}$$

and notice that the equation splits into two Weyl equations, e.g.

$$i\overline{\sigma} \cdot \partial \psi_L = 0$$

In 5 dimensions, this is not possible. We must add one more Dirac matrix, and this does not respect the splitting.

$$\gamma^5 = -i\Gamma = -i \begin{pmatrix} -1 & 0 \\ 0 & 1 \end{pmatrix}$$

So the basic fermion representation in 5 dimensions is a 4-component Dirac fermion.

Now perform the KK reduction of this Dirac fermion

$$\Psi = \sum_n \hat{\Psi}_n(x) e^{ik_n x^5}$$

$\hat{\Psi}_n$ obeys $(i\gamma^\mu \partial_\mu - \gamma^5 k_n)\hat{\Psi}_n = 0$

Let $\Psi_n = \frac{1+\gamma^5}{2}\hat{\Psi}_n$; multiply this equation by $\frac{1-\gamma^5}{2}$ (not a projector in this notation). We find

$$(i\gamma^\mu \partial_\mu - k_n)\Psi_n = 0$$

That is, $\hat{\Psi}_n$ obeys the 4-d Dirac equation with mass

$$k_n = n/R$$

It seems that, for n=0, we get massless left- and right-handed Weyl fermions. However, we must be more careful about the boundary conditions.

Since $\gamma^5 \partial_5$ is odd under P and mixes ψ and ψ_R, we must assign these fields opposite parity:

$$P\psi_L = +\psi_L \quad P\psi_R = -\psi_R$$

$$\Psi = \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix}$$

Then these fields obey different boundary conditions at the fixed points. One is compatible with $k^5 = 0$; the other is not.

$$\psi_L(x^5) = +\psi_L(-x^5) \quad \psi_R(x^5) = -\psi_R(-x^5)$$

Thus, each 5-d fermion field yields one 4-d Weyl fermion, plus a KK tower of 4-d massive Dirac fermions.

With a little more work, it can be shown that the same conclusion holds if we give the 5-d fermion a nonzero 5-d mass term.

The 5-d fermion thus leads to a massless chiral fermion in 4-d, plus a KK tower of massive fermions.

These massive fermions have the same definite SU(2)xU(1) quantum numbers as the chiral fermion; thus, they can obtain mass consistent with electroweak symmetry.

The first KK excitation of t_R is a massive singlet quark T.

The KK excitations of $(t, b)_L$ are massive vectorlike SU(2) doublet quarks.

In this setup, symmetry breaking can be driven by the Hosotani-Toms mechanism. I will explain this in the context of an Abelian gauge theory.

Consider a constant gauge field value $A_5 \neq 0$
Ordinarily, we think, a constant gauge field can be gauged away.

However, in the periodic space, this gauge transformation changes the boundary condition

$$\Psi(x+2\pi R) = \Psi(x) \quad \longrightarrow \quad \Psi(x+2\pi R) = e^{i2\pi g R A_5}\Psi(x)$$

The vacuum energy can depend on the boundary condition, and so integrating out the field Ψ generates a potential for A_5 .

In fact, $A_5 = 0$ is a local maximum. To understand this note that

anti-periodic b.c.s for Ψ give $\mathrm{tr}[e^{-\beta H}]$

periodic b.c.s for Ψ give $\mathrm{tr}[(-1)^F e^{-\beta H}]$

Now we have all of the ingredients needed to make a realistic model.

As in the Little Higgs case, we need another effect that generates a large enough 4-Higgs coupling. This requires that SU(2)xU(1) be embedded appropriately in a larger non-Abelian group.

The particle spectrum is similar to that of Little Higgs models, at least for the lowest levels. A geometrical symmetry

$$x^5 \to -x^5$$

can play the role of a parity that requires the lowest KK states to be pair-produced.

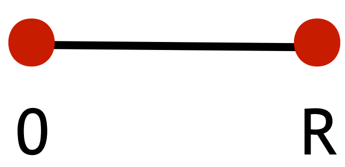
In this example, the higher dimensional space is flat. However, there is another approach to an extra-dimensional explanation of the weak interaction scale that makes essential use of strong spacial curvature. This is the model of Randall and Sundrum.

In this model, we again have the 5th dimension as an interval with two boundaries. But now we will include gravity as a more serious element. We take the boundaries to have associated energy that is significant on the scale of $m_{\rm Pl}$, and we take the interior of the space to have energy that also scales with $m_{\rm Pl}$.

Now look for a solution with 4-d Lorentz invariance, but with nontrivial dependence on the 5th dimension:

$$ds^2 = e^{-f(y)}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2$$

$$y = x^5$$



Randall and Sundrum, and Gogberashvili, found the following solution:

$$\Lambda_0 = -\Lambda_R = C m_{\mathrm{Pl}}^4 \qquad \Lambda = -C m_{\mathrm{Pl}}^4 \cdot k$$

where $k = b\, m_{\mathrm{Pl}}$ and the metric is

$$ds^2 = e^{-2ky}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2$$

The metric is warped, with excursions in x representing smaller distances at larger y.

Now put matter fields at the fixed points (branes) at the ends of the space. The fixed point Lagrangian is 4-dimensional,

$$\int d^4x(\mathcal{L}_* - \sqrt{g}\Lambda_*)$$

where * = 0,R, and it includes the metric induced from the curvature of the 5-d space.

It is interesting to exploring postulating a Higgs field at y = R:

$$\int d^4x \mathcal{L} = \int d^4x \sqrt{g}[g^{\mu\nu}\partial_\mu \varphi^\dagger \partial_\nu \varphi - \lambda(|\varphi|^2 - V^2)^2]$$

Now, $g_{\mu\nu} = e^{-2ky}\eta_{\mu\nu}$ $\qquad \sqrt{g} = e^{-4ky}$

so

$$\int d^4x \mathcal{L} = \int d^4x [e^{-2kR}\partial^\mu \varphi^\dagger \partial_\mu \varphi - e^{-4kR}\lambda(|\varphi|^2 - V^2)^2]$$

We should rescale φ so that the kinetic term is in the standard normalization. Then

$$\int d^4x \mathcal{L} = \int d^4x [\partial^\mu \varphi^\dagger \partial_\mu \varphi - \lambda(|\varphi|^2 - e^{-2kR}V^2)^2]$$

The scale-invariant part of the Langrangian is not affected, but the potential is now minimized at $v = e^{-kR}V$

for $V \sim m_{\mathrm{Pl}}$ but $kR \sim 30$, $v \sim 100\ \mathrm{GeV}$.

In the warped geometry, coordinate excursions cover exponentially smaller distances as we move to larger y. This brings us closer to the Planck scale. An observer on the brane, though, sees the energy scale decrease exponentially with y. In this view, the Planck scale -- which also gives the cutoff for quantum field fluctuations -- decreases exponentially with y. For y = R, the highest-frequency quantum fluctuations are at the weak scale. The stability of the weak scale then follows from the 5-d spacetime geometry.

In the following, I will refer to the brane at y = R as the weak brane. The opposite endpoint, at y = 0, is the Planck brane.

In our examples with flat extra dimensions, each field in the bulk generated a KK tower of excitations that would appear as new particles in collider experiments. In the Randall-Sundrum geometry, the gravitational field in the bulk will give KK excitations. If we put matter fields into the bulk, these will also give towers of KK states.

To find the KK reduction, we need to find the eigenstates of the wave equation in the warped geometry. This is easiest for a massless scalar field.

$$\int d^5x\mathcal{L} = \int d^5x\sqrt{G}\,\frac{1}{2}G^{MN}\partial_M\phi\partial_N\phi$$
$$= \int d^5x e^{-4ky}[e^{2ky}\partial^\mu\phi\partial_\mu\phi - (\partial_y\phi)^2]$$

Look for a solution of the variational equation of the form

$$\phi_n(x,y) = e^{-ip_n\cdot x}f_n(y)$$

The equation for $f_n(y)$ is

$$-\partial_y e^{-4ky}\partial_y f_n(y) = e^{-2ky}p_n^2 f_n(y)$$

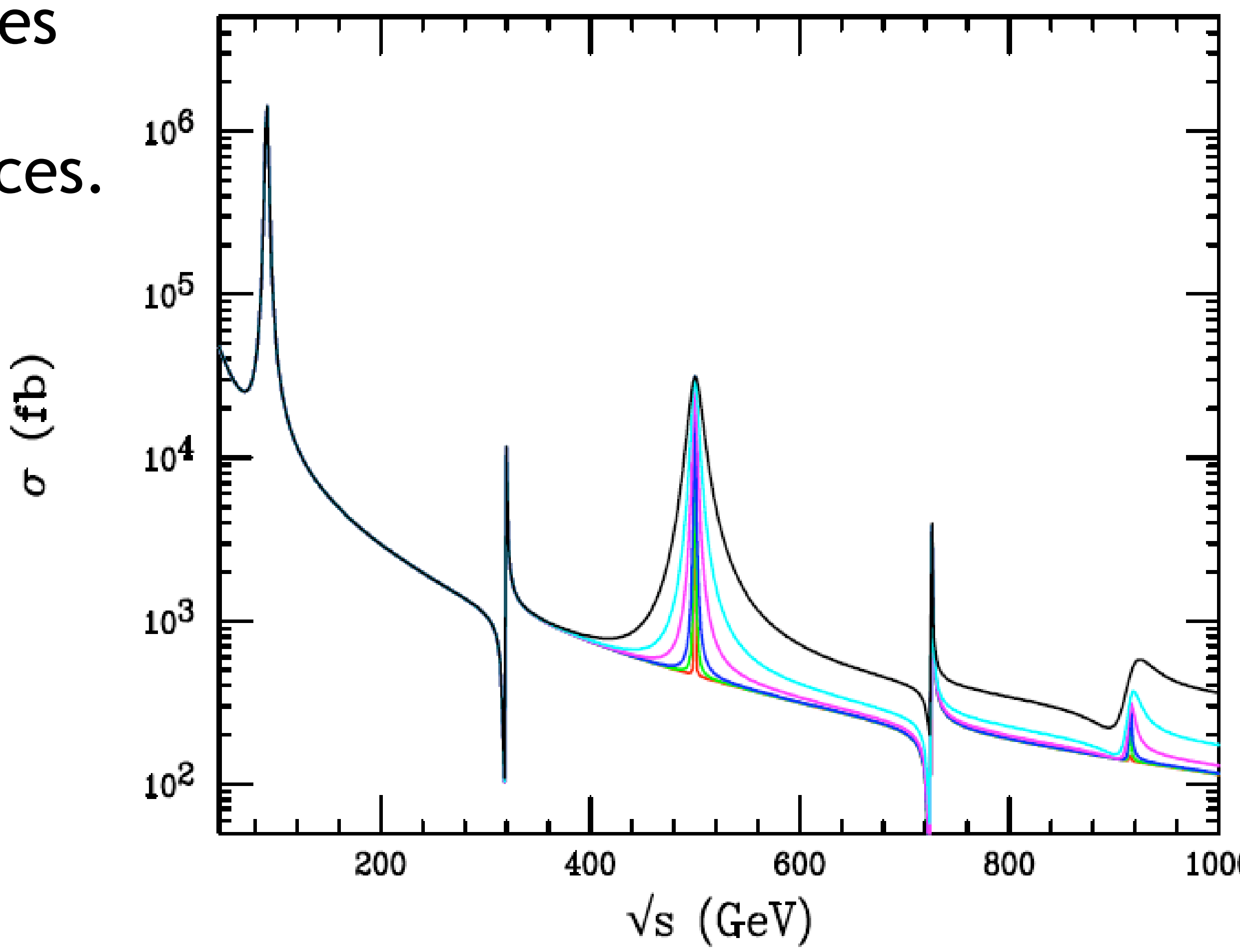
The solutions of this equation are Bessel functions

$$f_n(y) = e^{2ky}J_2(\frac{m_n}{k}e^{ky})$$

where $m_n^2 = p_n^2$ is the mass of the excitation. The m_n are determined by imposing a boundary condition at y = R. This fixes $x_n = (m_n/k)e^{kR}$. That is $m_n \sim ke^{-kR}$,
the KK masses are at the weak scale!

For m = 0, there is a constant solution to the equation on the previous slide. This gives the lowest state in the tower. For a Yang-Mills theory in the bulk, this state gives the effective Yang-Mills gauge boson.

The KK tower of states then appears as a sequence of resonances. For example, for photon, Z, and graviton states in $e^+e^- \to \mu^+\mu^-$



Davoudiasl, Hewett, Rizzo

There is an additional interesting feature for fermions in the Randall-Sundrum space. Consider the Dirac action, including a mass term:

$$\int d^5x\sqrt{G}\ \overline{\Psi}(i\Gamma^M D_M - m)\Psi$$

Γ^M is the curved-space Dirac matrix, and D_M includes the spin connection. We would like to apply orbifold boundary conditions and find chiral fermions, just as we did in the flat space example. To do this, write the variational equation and look for solutions with $\gamma^\mu\partial_\mu\Psi = 0$

Let $\hat{\Psi}(x) = e^{-ky\cdot(3/2)}\Psi$ then

$$[\Gamma\partial_y - \frac{1}{2}k\Gamma - m]\hat{\Psi} = 0 \qquad \Gamma = \begin{pmatrix} -1 & \\ & 1 \end{pmatrix}$$

The second term comes from the spin connection. The orbifold boundary conditions allow a solution only for $\Gamma = -1$. This is

$$\hat{\psi} = \psi(x)e^{(\frac{1}{2}-c)ky} \qquad c = \frac{m}{k}$$

For model-building, we can think of c = m/k as a parameter that we can adjust for each 5-d fermion multiplet. Setting $c > 1/2$ moves the wavefunction of the fermion toward the Planck brane. Setting $c < 1/2$ moves the wavefunction of the fermion toward the weak brane.

Since the Higgs field is localized at the weak brane, multiplets with larger c have small masses. Also, flavor-changing effects due the Higgs and to KK resonance exchange are highly suppressed.

However, there is a problem with the top quark. To suppress corrections to $\Gamma(Z^0 \to b\bar{b})$, the multiplet $Q_t = (t, b)$ must be assigned $c > 1/2$. However, then we must set $c < 1/2$ for the $\bar{t}$. Then KK resonances should create form factor effects in the $\bar{t}$ gauge couplings.

Through some theoretical insight, the physics of warped 5-d space can be connected to physics of bound states in 4-d space.

In 1997, Maldacena proposed that 4-d N=4 super-Yang-Mills theory and Type IIB string theory on the 10-d space $AdS_5 \times S^5$ were equivalent theories. Much evidence has accumulated since then that operator matrix elements in N=4 super-Y-M can be computed by considering these as operators on the boundary of anti-de Sitter space.

An essential part of this construction is the interpretation of shifts in the extra dimension of AdS (anti-de Sitter space) as scale transformations in the the Yang-Mills theory.

We have a similar structure in a general Randall-Sundrum model. The geometry of the bulk is a space of constant negative curvature, that is, AdS. Up to the effects of the boundary conditions, the model is scale-invariant, with a shift in y corresponding to a scale factor $e^{-k\Delta y}$ in the energies of KK modes.

So we can think of Randall-Sundrum theory in an alternative way as providing a tool to construct working models of a strongly-coupled gauge theory. In this correpondence, we reinterpret the KK resonances as bound states of the strong interactions.

For a general choice of fields, the resulting theory will probably not be a local field theory when reinterpreted in 4-d. That is the price for having a model of strong interaction in which we can do explicit computations.

The three types of models described in this lecture all have superficially similar particle content. We can search for the particles predicted by these models in a unified way.

After we find these particles and measure their properties, we can ask whether the correct model is literally one of these three or a more complex model in this general class.

That is not the question for today. First, we want to find the first evidence of new particles connected to electroweak symmetry breaking. I will discuss that search in the next lecture.