Electroweak precision measurements and collider probes of the standard model with large extra dimensions

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The elementary particles of the standard model may reside in more than 3 + 1 dimensions. We study the consequences of large compactified dimensions on scattering and decay observables at high-energy colliders. Our analysis includes global fits to electroweak precision data, indirect tests at high-energy electron-positron colliders (CERN LEP2 and NLC), and direct probes of the Kaluza-Klein resonances at hadron colliders (Fermilab Tevatron and CERN LHC). The present limits depend sensitively on the Higgs sector, both the mass of the Higgs boson and how many dimensions it feels. If the Higgs boson is trapped on a (3 + 1)-dimensional wall with the fermions, large Higgs boson masses (up to 500 GeV) and relatively light Kaluza-Klein mass scales (less than 4 TeV) can provide a good fit to precision data. That is, a light Higgs boson is not necessary to fit the electroweak precision data, as it is in the standard model. If the Higgs boson propagates in higher dimensions, precision data prefer a light Higgs boson (less than 260 GeV) and a higher compactification scale (greater than 3.8 TeV). Future colliders can probe much larger scales. For example, a 1.5 TeV electron-positron linear collider can indirectly discover Kaluza-Klein excitations up to 31 TeV if 500 fb⁻¹ integrated luminosity is obtained.

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I. INTRODUCTION

The original motivation for adding a large compact dimension was to generate (3 + 1)-dimensional vector gauge fields from a purely gravitational action in higher dimensions (see [1,2] for a review). Describing nature completely by this mechanism is not viable. Not only is matter unexplainable in this approach, but the (3 + 1)-dimensional action inescapably contains a massless scalar particle that successfully competes with a spin-2 particle (graviton) to create a strong mix of scalar-tensor gravity unacceptable to modern experiment.

One conceptual cousin of the original Kaluza-Klein (KK) idea is string theory, or M theory, where strings and D-branes populate the higher dimensional space rather than just a spin-2 graviton (see [3,4] for reviews). A strong motivation for string theory is that it may be finite, and may thus provide a self-consistent description of quantum gravity. String theory also predicts troublesome scalar moduli particles that make it a challenge to identify the ground state of the theory. Solutions to this problem have been postulated, and progress has been made on other aspects of the theory, giving hope that it may be possible in time to write down a string theory description of nature.

Recently, it has been pointed out that there are more reasons to suggest extra dimensions than just having a self-consistent description of gravity [5–7]. The additional motivations include new directions to attack the hierarchy problem [5] and the cosmological constant problem [8,9], unifying the gravitational coupling with the gauge couplings [10–12], perturbative supersymmetry breaking in string theory [13,14], and low-scale compactifications of string theory [15,16,7,17,18]. An important breakthrough was the realization that the gravitational scale could be as low as the weak scale and still be phenomenologically viable [5,6]. Two or more large extra dimensions felt by gravity are required to make this possible. Another tantalizing realization is that gauge couplings may unify with a greatly reduced string scale if gauge fields feel one or more large extra dimensions [19–27]. A tentative picture is filling in for a viable scenario with TeV-scale extra dimensions, and especially TeV-scale string theory with a vastly reduced Planck scale, compactification scale, string scale, and unification scale.

In this article, we focus on the phenomenology of the gauge and matter sectors of theories with large extra dimensions. In particular, the Kaluza-Klein states of the gauge particles and matter particles can have important observable consequences at high-energy colliders. It is these consequences that we wish to study. We build on previous studies that assumed a similar framework and discussed relevant collider phenomenology [28–35].

In principle, gravitational radiation into extra dimensions and virtual graviton induced observables are correlated with observables generated by KK excitations of the gauge and matter fields. Many detailed studies on gravitational effects at high-energy colliders [36–58] and important astrophysical bounds [59–62] have appeared. However, to know the correlations between these effects and what we study here requires that we either specify the underlying theory, or assume that gravitational effects do not pollute the signals. We choose the latter path by assuming that gravity propagates in significantly more extra dimensions than the gauge and matter fields do. For example, gravity may propagate in ten dimensions, while gauge fields are confined to a pₜ brane (gauge bosons) or 3-brane (fermions). (pₜ) is defined to be the

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number of spatial dimensions that bulk gauge fields feel, and
$p_\perp$ is defined to be the number of remaining spatial dimen-
sions in which gravity propagates.) Also, we assume that the
higher-dimensional gravity scale and the gauge-unification
scale are comparable, as is expected in string theory. These
assumptions imply that gravitational radiation will not be as
significant as gauge KK excitations in collider phenomenol-
ogy. The exact strength of virtual graviton exchange effects
is not calculable, and so it is difficult to tell how probing they
are with respect to the gauge interactions pursued here. Es-
timates based on naive dimensional analysis suggest that the
virtual graviton exchange processes in some cases may be
comparable in the probing power of extra dimensions as KK
excitations of gauge bosons given the above assumptions.

In the following sections we define a five-dimensional
standard model (5DSM). Particularly important is the defini-
tion of the Higgs boson fields in this Lagrangian, since elec-
troweak symmetry breaking effects will correlate strongly
with some observables. We then compactify the extra dimen-
sions and work in an effective field theory that is the stan-
dard model plus additional non-renormalizable interactions
arising from integrating out KK excitations. We then do a
global fit to precision electroweak (EW) data and find limits
on the gauge compaction scale. Several comparisons of
precision electroweak data to the SM with extra dimensions
have been published recently [30–32,63]. Our contributions
in this direction are to construct a global fit to all relevant
data, and to present results in terms of operator coefficients
rather than just a fifth dimension compaction scale. One
result from the global fit demonstrates that a light Higgs
boson is not necessary, in contrast to the standard model fit
which requires it. We also study the possibility of finding the
first excited state at hadron colliders, and derive sensitivities
to the full KK tower at $e^+e^-$ colliders. In the last section we
conclude and summarize the results.

II. STANDARD MODEL IN EXTRA DIMENSIONS

We begin by considering only one extra spatial dimension
beyond the usual $3 + 1$ dimensions. Our first task is to state
which standard model particles reside in five spacetime di-
mensions and which reside only in four dimensions. In order
to obtain massive chiral fermions we assume that the fermi-
ons reside in the “twisted sector” of string theory, and so are
naturally confined to “walls” of an orbifold fixed point in
the higher dimensional space. It may be possible for the chil-
mal fermions to reside in a higher dimensional space and change
some of the results in this paper. The gauge fields are
non-chiral and so may reside with impunity in higher dimen-
sions, that is the fifth dimension, or the “bulk.” These as-
sumptions are essentially identical to those made in Refs.
[64,66,31].

A. Higgs sector

It is somewhat more difficult to decide what to do with
the Higgs fields. They are non-chiral fields as well, and with
no reference to a more fundamental theory it appears natural
to put them in the bulk with the gauge fields. To answer this
question more satisfactorily, it is necessary to discuss the
role of supersymmetry [64–66]. The more fundamental
theory is likely to contain space-time supersymmetry. In-
deed, one of the motivations for large extra dimensions is the
ability to obtain tree-level supersymmetry breaking at $R^{-1}$
from Scherk-Schwartz compactification of a TeV string
theory. The superpartners will then have masses near $R^{-1}$
and will have little effect on current collider phenomenol-
yogy as long as $R^{-1}$ is as low as several $200$ GeV.

As a consequence of supersymmetry, two Higgs doublets
are necessary in the spectrum, $H_u$ which gives mass to up-
type quarks and $H_d$ which gives mass to down-type quarks
and leptons. If one Higgs boson is on the wall and the other
Higgs boson is in the bulk, then successful gauge coupling
unification is possible with only the states of the minimal
supersymmetric standard model (MSSM) in the low-energy
spectrum [20]. Unification is also possible by putting both
Higgs fields in the bulk along with extra particles that may
be necessary for proton stability and other reasons
[19,20,22,67,68]. Alternatively, it may not be necessary to
require both Higgs fields to be zero-mode excitations under
orbifolding [69,20].

We therefore allow our Higgs sector to contain Higgs
field(s) in the bulk and Higgs field(s) on the wall [31]. We
define

$$\tan \phi = \frac{\langle \varphi_{\text{wall}} \rangle}{\langle \varphi_{\text{bulk}} \rangle},$$

where $\langle \varphi_{\text{wall}} \rangle$ is the vacuum expectation value of the Higgs
field on the wall, and $\langle \varphi_{\text{bulk}} \rangle$ is the vacuum expectation value of
the Higgs field in the bulk. In some theories $\varphi_{\text{wall}}$ can be
identified with either $H_u$ or $H_d$, and $\varphi_{\text{bulk}}$ can be identified
with the other Higgs field of the MSSM. In these cases,
$\tan \phi = \tan \beta$ or $\tan \phi = 1/\tan \beta$, where $\tan \beta = (H_u)/(H_d)$. However, the low-energy effective theory may more natu-
aturally be described in terms of a single Higgs boson
originating from non-chiral bulk field(s), in which case
$\tan \phi = 0$ ($\sin \phi = 0$). Furthermore, although supersymmetry
may be necessary for a viable string scenario, the most eco-
nomical model is the standard model with one Higgs field
either in the $p_1$-brane bulk or confined on the wall. There-
fore, the choices $\tan \phi = 0$ and $\tan \phi = \infty$ will be of particular
interest when we discuss the EW precision measurement pre-
dictions below.

B. 5DSM Lagrangian and renormalized parameters

Our starting framework is equivalent to Refs. [64,66,31],
where we assume the vector bosons and one Higgs field ($\varphi_h$)
reside in the 5D bulk, and the fermions and another Higgs
field ($\varphi_w$) reside on the 4D wall or boundary of the $S^1/Z_2$
orbifold. In five dimensions, the kinetic terms of the La-
grangian are simply

$$L_5 = \int d^4x \left[ -\frac{1}{4} F_{AB}^2 + |D_A \varphi_h|^2 ight. $$
$$+ \left. (i \bar{\psi} \gamma^\mu D_\mu \psi + |D_\mu \varphi_w|^2) \delta(y) + \cdots \right],$$

where $\hat{g}$ is the 5D gauge coupling in the covariant derivative,
$\varphi_b$ is the Higgs boson in the bulk, and $\varphi_w$ is the Higgs boson on the wall. The $\delta(y)$ function indicates that the fermions and $\varphi_w$ fields are localized at $y=0$, the location of the 3-brane wall.

Compactifying the fifth dimension on a $S^1/Z_2$ line segment, one finds

$$L_4 = \int d^4x \left[ \sum_{n=0}^{\infty} \left( -\frac{1}{4} F_\mu^n \right)^2 + \frac{1}{2} \left( \frac{n^2}{R^2} + 2 g^2 |\varphi_b|^2 \right) V_\mu^n V_\mu^n \right]$$

$$+ g^2 |\varphi_w|^2 \left( V_\mu^{(0)} + \sqrt{2} \sum_{n=1}^{\infty} V_\mu^n \right)^2 + i \bar{\psi} \sigma^\mu \left[ \partial_\mu + ig V_\mu^{(0)} \right] \psi + \cdots ,$$

where $g = \tilde{g}/\sqrt{\pi R}$ is the four dimensional gauge coupling. In the non-Abelian case, one should replace $V_\mu^{(n)}$ with $\lambda^a \gamma_\mu^{(n) a}$, where $\lambda^a$ are the group generators, to obtain the appropriate expressions. From this Lagrangian interactions in the theory are specified. The KK states have an additional $\sqrt{2}$ strength in their interactions, which may appear odd at first sight. This factor arises from rescaling the gauge kinetic terms to be canonically normalized for all $n$. However, we emphasize that it is a model dependent consideration. For two or more extra dimensions the sum diverges. However, a more accurate application of the fundamental theory indicates that $g$ depends on $n$, and in general given by [29,34]

$$g_n \sim g \exp \left( \frac{-n \cdot \tilde{n}}{R^2 M_s^2} \right),$$

where $M_s$ is the string scale. This behavior is in qualitative agreement with string scattering amplitudes at high energy which tend toward zero. The exponential suppression then cures the problem of divergent summations of KK states. However, the precise coefficients and form of Eq. (6) is model dependent.

Also, there are many other model dependent considerations that will yield different couplings of KK gauge bosons to different fermions. For example, in Ref. [70] it was pointed out that this situation arises if fermions are stuck to different points of a thick wall. In this case, the KK phenomenology could be qualitatively different than what is presented here.

In an effort to be as model independent as possible, we present all our “indirect” search results in terms of a parameter $V$ which is defined to be

$$V = 2 \sum_n \left( \frac{g_n^2}{g^2} \right) \frac{m^2_W}{n^2 M_s^2} .$$

Although not entirely model independent, this is the quantity that can account for variations of $g$ for different $n$ in the summation of the correct effective theory and the regularization of the KK sum. Often, for concreteness, we will translate a limit of $V$ that we find into a limit on $M_s$ by assuming one extra dimension and that $g_n^2 = g$ for all $n$. We must also keep in mind that other subtleties of the full theory may contribute to collider phenomenology in addition to what we have discussed here [71–73].

### III. Precision Measurements

In the standard model, all physical observables can be predicted in terms of a small set of input observables. Equivalently, the standard model contains several parameters in the Lagrangian which can be fit to by comparing calculations within the model to measurements. There are more ob-
servables than there are parameters, and so the fit is over-constrained. A global $\chi^2$ analysis to precision electroweak data can determine if a particular model, such as the SM, is a consistent description of nature.

In the following we do a global analysis of EW precision measurement data using the higher dimensional standard model (HDSM). In the limit that the extra compactified dimensions’ radii tend to zero, we will recover the standard model global fit results. It has been often stated that the SM fits the EW precision data very well; however, this is only true if we assume that the Higgs boson is light. In the 5DSM there are two more parameters in the theory beyond the usual SM parameters that will impact precision measurement predictions [30–33]. These parameters are $\tan\beta$ [the ratio of wall-Higgs vacuum expectation value (VEV) to bulk-Higgs VEV] and $R^{-1}=M_c$ (the compactification scale). We shall see below that strong correlations exist between allowed values of $\tan\beta$, $M_c$, and $m_H$ once we require that the 5DSM be consistent with all measurements.

### A. Global fits with physical observables

The procedure for carrying out a global fit is the same for the HDSM as it is for the SM: (1) Construct the full bare Lagrangian of the theory, $\mathcal{L}(g_0, m_0, \psi_0, \ldots$).

(2) Split the bare parameters and bare fields into renormalized quantities and counterterms, $\mathcal{L}(g + \delta g, m + \delta m, \psi + \delta\phi, \psi/2, \ldots)$.

(3) Decide on a renormalization scheme [modified minimal subtraction (MS)], on shell, etc. that sets the values of the counterterms (e.g., set to a loop correction at a particular scale). For tree-level calculations, it is most convenient to set the counterterms to zero.

(4) Calculate all observables using the renormalized Lagrangian. From the previous steps the result will be finite and depend only on renormalized couplings $O_i(g, m, \ldots)$.

(5) Perform a constrained global fit to see if there is a set of renormalized couplings $g, m, \ldots$ that allows $O_i(g, m, \ldots) = O_{\text{exp}}^{\text{fit}}$ to within experimental uncertainty.

In some cases a model can be completely ruled out by the above procedure, whereas in other cases like the SM and the 5DSM, the model can work for a limited range of parameter choices for the as-yet unknown parameters.

There are many observables that we wish to compare predictions with experiment. Above, we specified the Lagrangian and renormalized parameters that enable us to carry out this program. In this section we write down, analytically, the calculated observables at leading order in an expansion of $m_w^2/M_c^2$. We expand in $m_w^2/M_c^2$ (i.e., $V$) since we know that we recover the good SM fit to data as $M_c \gg m_w$. The physical vector boson masses are, to leading order in $m_w^2/M_c^2$,

$$m_w^2 = \frac{1}{2} g^2 V^2 \left[1 - s_\phi^4 V\right],$$

and

$$m_Z^2 = \frac{1}{2} \frac{g^2}{\cos^2 \theta} V^2 \left[1 - s_\phi^4 V/\cos^2 \theta\right].$$

where we define

$$V = 2 \sum_{n=1}^{\infty} \frac{1}{n^2} = \frac{\pi^2}{3} \frac{m_w^2}{M_c^2},$$

and $g = g_1/g_0$. $g_1 = g/\tan \theta$, and $V^2 = (\phi_{\text{wall}})^2 + (\phi_{\text{bulk}})^2$. The last equality in Eq. (10) is valid only if $p_1 = 4$ (one extra spatial dimension). It is also convenient to define a charge-current and neutral-current interaction coupling with the lightest $W$ and $Z$ mass eigenvalues,

$$g_w W^\mu + \text{H.c.} \rightleftharpoons \frac{\sqrt{2}}{\sqrt{2}} \left[1 - s_\phi^2 V\right]$$

$$g_Z \cos \theta Z^\mu + \text{H.c.} \rightleftharpoons \frac{g Z}{\cos \theta} \left[1 - s_\phi^2 V/\cos^2 \theta\right].$$

We can now express more easily other observables in terms of the physical vector boson masses $m_w$ and $m_Z$ and the definitions $g_w$ and $g_Z$ provided above. For example,

$$G_\ell (\mu \text{ decay}) = \frac{\sqrt{2} g_w^2}{8 m_w^2} \left[1 + V\right],$$

$$\Gamma(Z \rightarrow \ell\ell) = \frac{N_c m_Z}{12 \pi} \left|\frac{g Z}{2 \cos \theta}\right|^2 \left|V + a_f^2\right|^2,$$

$$A_f = \frac{2 \nu_f a_f}{\sqrt{2} + a_f^2},$$

$$Q_w = -\frac{4}{m_Z^2} \left[\frac{g^2 (1 - s_\phi^2 V/\cos^2 \theta)^2}{\cos^2 \theta} + \frac{g^2 V}{\cos^2 \theta}\right],$$

$$\times a_e \left[\nu_e (2Z + N) + \nu_d (2N + Z)\right].$$

$$R = \frac{\sigma_{NC}^v - \sigma_{NC}^v}{\sigma_{CC}^v - \sigma_{CC}^v} = \left[\frac{g^2 Z}{(1 - \sin^2 \theta)m_Z^2} + \frac{g^2 V}{(1 - \sin^2 \theta)m_z^2}\right]
\times \left[\frac{g^2_w}{m_w^2} + \frac{g^2 Z}{m_Z^2}\right]^{-1} \left(1 - \frac{1}{2} \sin^2 \theta\right),$$

$$\sin^2 \theta_{\text{eff}} = x + \frac{x(1-x)}{1-2x} \left[\frac{c_\phi}{x} - \frac{s_\phi}{1-x}\right],$$

$$m_w^2 = m_Z^2 (1-x) \left[1 + \frac{1}{1-2x} \left(c_\phi (1-x) - s_\phi\right)\right],$$

where $Q_w$ is a measure of atomic parity violation, $x$ is the solution to the equation

$$x(1-x) = \frac{\pi \alpha}{\sqrt{2} G_f m_Z^2}.$$
In this fit we have held fixed $m_Z = 91.1867$ GeV [79], $G_F = 1.16637 \times 10^{-5}$ GeV$^{-2}$ [75], $\alpha_s(m_Z) = 0.119$ [79], $m_t = 173.8$ GeV [76], and $\alpha_{\text{QED}}(m_Z) = 1/128.933$ [80].

We assume that one physical Higgs scalar boson is present in the spectrum which interacts with the fermions and gauge bosons like a SM Higgs boson. The other physical Higgs degrees of freedom either do not exist or have interactions decoupled from the zero modes of the gauge bosons and the fermions. This is analogous to the MSSM, where one Higgs boson acts like a SM Higgs boson and the rest decouple, being irrelevant for precision measurement analyses.

Our procedure, then, is to choose a Higgs boson mass and vary $V$ to see how the predictions change for the observables. We wish to minimize the $\chi^2$ function defined as

$$\chi^2 = \sum_i \frac{[C_i^{\text{theory}}(g, m_h, V, \ldots)-C_i^{\text{exp}}]^2}{(\Delta C_i^{\text{exp}})^2}.$$  (29)

We also define $\Delta \chi^2 = \chi^2 - \chi^2_{\text{min}}$.

In Fig. 1 we plot $\chi^2$ with $\tan \phi = 0$ and for differing choices of $m_h$. In the SM, the 95% C.L. upper bound on the Higgs boson mass is 260 GeV [79]. In this plot the $\chi^2$ value for $V = 0$ (decoupled extra dimensions) and $m_h = 260$ GeV is $\chi^2 = 22.1$. We then allow $V$ to vary from zero and $m_h$ to vary, and define the allowed region of parameter space to be that which has $\chi^2 < 22.1$. From Fig. 1 we can see that the light Higgs boson is favored for $V = 0$, just as the well-known SM results indicate. Furthermore, as we increase the Higgs boson mass the best fit value of $V$ drifts more and more into the $V < 0$ region. Within the context of the 5DSM, negative values of $V$ are not physical. Increasing the value of $V$, or, equivalently in the 5DSM, lowering the compactification scale $M_c$, we see that the electroweak precision data fit only gets worse for any value of $m_h$. The largest value of $V$ with $\chi^2 < 22.1$ is $V = 0.0015$. Therefore, the limit on $V$ in this theory is $V < 0.0015$ which is equivalent to $M_c > 3.8$ TeV in the 5DSM.

For $\tan \phi = \infty$, meaning the only Higgs boson(s) associated with electroweak symmetry breaking (EWSB) on the wall, we find the opposite behavior. In Fig. 2 we plot the $\chi^2$ for various choices of $m_h$, with $\tan \phi = \infty$ and with $V$ varying. In this case, the fit remains good as $V$ increases and $m_h$ increases. (Note that the $V = 0$ slice is equivalent to the $V = 0$ slice of Fig. 2.) A similar relaxing of the SM Higgs boson mass limit from precision data has been demonstrated in other contexts [63,82,83]. Now, all the minima of the $\chi^2$ fits are in the $V > 0$ physical region. For $m_h > 260$ GeV, a non-zero value of $V$ is required to be present in the theory in order to provide an acceptable fit to the data. As demonstrated in Fig. 2, the Higgs mass could be heavy and as high as 500 GeV and still have $\chi^2 < 22.1$ as long as $V = 0.0016$. That is, KK excitations of gauge bosons must substantially
In this case, the low-mass Higgs boson is not needed if \( V \) is sufficiently high. If the Higgs mass gets above 500 GeV, then there is no longer a choice of \( V \) for which \( \chi^2 < 22.1 \). This value is compatible with precision electroweak data with \( \tan \beta = 1 \). For \( V > 0 \) the value of \( \sin^2 \theta_W \) decreases and \( m_W \) increases, precisely what lowering the Higgs boson mass would do to the predictions. In this case, the low-mass Higgs boson is not needed if \( V \) is sufficiently high.

We next ask what the 95% C.L. range is for \( V \) given a fixed \( m_h \). This question differs slightly from the previous question, in that we are no longer asking how good of a fit a particular value of \( m_h \) is, but rather what deviations of \( V \) would be tolerated if \( m_h \) were given to us from another source (direct experiment, “by God,” etc.). For this we must analyze the \( \Delta \chi^2 \) distribution, which is defined to be \( \chi^2 - \chi^2_{\text{min}} \), where \( \chi^2_{\text{min}} \) is the lowest value of \( \chi^2 \) for a fixed \( m_h \) but variable \( V \). Then, the 95% range of \( V \) is determined by requiring \( \Delta \chi^2 < (1.96)^2 \). The case where a negative \( V \) provides the best fit must be handled by following the Feldman-Cousins prescription [81]. In Table I we show these ranges of \( V \) for a given \( m_h \) and \( \sin^2 \phi \). The blank spaces in the table mean that there is no range of \( V \) allowed in the physical region, and the parentheses mean that there is no choice of \( V \) for that particular \( m_h \) and \( \sin^2 \phi \) which gives \( \chi^2 < 22.1 \). Therefore, any non-blank entry without parentheses means that the corresponding Higgs boson mass is not ruled out for that given choice of \( \sin^2 \phi \) and at least one value of \( V \).

### IV. KALUZA-KLEIN EXCITATIONS AT HIGH-ENERGY COLLIDERS

The results of the previous section were obtained by comparing precision measurements of electroweak observables to the theoretical predictions of the HDSM. These results were mainly derived from how the zero-mode vector bosons interact with the KK excitations and how the KK excitations of the \( W \) and \( Z \) directly affect observables with characteristic energy below \( m_Z \) (\( \mu \) decay, \( \nu N \) scattering, etc.). In this section, we estimate the sensitivity of KK excitations to \( e^+ e^- \) scattering at high-energy colliders above \( m_Z \). This will involve operators induced by higher modes of the \( W/Z/\gamma \) gauge bosons and also on-shell production of KK excitations of the SM gauge bosons.

### Table I. 95% C.L. \([\Delta \chi^2 < (1.96)^2]\) allowed ranges of \( V/10^{-3} \) for different values of \( \sin^2 \phi \) and \( m_h \) in the global \( \Delta \chi^2 \) distribution to precision electroweak data. Blank spaces in the table mean that the fit to the data is too poor to quote a bound in the physical region of \( V > 0 \). Parentheses mean that there is no value of \( V \) such that \( \chi^2 < 22.1 \), implying that the corresponding Higgs boson mass is not allowed from the global \( \chi^2 \).

<table>
<thead>
<tr>
<th>( m_h ) [GeV]</th>
<th>( \sin^2 \phi = 0 )</th>
<th>( \sin^2 \phi = 0.25 )</th>
<th>( \sin^2 \phi = 0.50 )</th>
<th>( \sin^2 \phi = 0.75 )</th>
<th>( \sin^2 \phi = 1.0 )</th>
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<td>&lt;1.56</td>
<td>&lt;2.25</td>
<td>&lt;1.42</td>
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<td>1.21</td>
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<td>0.53</td>
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<td>0.20–1.45</td>
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<td>2.08</td>
<td>0.33–2.30</td>
<td>0.34–1.63</td>
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<td>0.44–1.77</td>
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<tr>
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<td>0.54–1.87</td>
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<td>(0.63–2.75)</td>
<td>0.69–2.03</td>
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<td>(0.85–2.97)</td>
<td>0.88–2.22</td>
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<td>(1.29–3.41)</td>
<td>(1.25–2.61)</td>
<td></td>
<td></td>
</tr>
<tr>
<td>900</td>
<td>(1.54)</td>
<td>(1.39–3.51)</td>
<td>(1.34–2.70)</td>
<td></td>
<td></td>
</tr>
<tr>
<td>1000</td>
<td>(1.50)</td>
<td>(1.49–3.61)</td>
<td>(1.45–2.79)</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>
As $M_c$ gets larger the excited modes of $\gamma^{(n)}$ and $Z^{(n)}$ obtain heavier and heavier mass and have minimal impact on the overall scattering amplitude. The amplitude contributions from the excited modes can be analyzed effectively by integrating out the heavy modes and constructing operators which take into account all their effects. For example, integrating out the higher modes of the photon yields an operator of the form

$$O_{\gamma^{(n)}} = \sum_n \left( \frac{-2g^2\sin^2\theta W Q_f}{n^2 M_c^2} \right) [\bar{e} \gamma^\mu e][\bar{f} \gamma^\nu f]$$

$$= -g^2\sin^2\theta W Q_f \frac{V}{m_W^2} [\bar{e} \gamma^\mu e][\bar{f} \gamma^\nu f].$$

(31)

Similar operators arise from integrating out $W^{(n)}$ and $Z^{(n)}$,

$$O_{W^{(n)}} = -\frac{g^2 V}{2m_W^2} [\bar{e} \gamma^\mu P_L e][\bar{f} \gamma^\nu P_L f].$$

(32)

$$O_{Z^{(n)}} = -\frac{g^2 V}{4\cos^2\theta m_W^2} [\bar{e} \gamma^\mu (\gamma_\nu a_\tau) e][\bar{f} \gamma^\nu (\gamma_\nu - a_\tau) f].$$

(33)

Limits can be set on $V$ from the effect of these operators on the total rates and polarization asymmetries of $e^+e^- \to f\bar{f}$ for all accessible fermions (see [85] for a discussion of the observables). In Fig. 3 we plot the search reach of $V$ versus integrated luminosity for $\sqrt{s} = 195$ GeV. To construct this plot we have included initial state radiation with a $10^\circ$ polar angle cut on the photons. The $b$ and $c$ quark tagging efficiencies are taken to be 35% and 20% respectively. We also assume that the KK states decay only into SM particles.

The conclusions may be weakened if these KK excitations were to decay some fraction of the time into superpartners. With over 1 fb$^{-1}$ one can either detect or rule out $V > 2.4 \times 10^{-3}$ (or $M_c < 3.1$ TeV for the 5DSM).

The same analysis can be applied at the Next Linear Collider (NLC). However, here we are able to add the top quark to the list of final states. Furthermore, we can include polarization asymmetries at the NLC, and we can utilize observables associated with $\tau$ lepton polarization. At the NLC we assume that $b, t$ and $c$ quark identification efficiencies are 60% and the efficiency for measuring tau polarization is 50% [84,85].

The sensitivity to $V$ at the NLC is substantial. In Fig. 4 we plot the search reach for $V$ versus integrated luminosity for $\sqrt{s} = 500$ GeV, 1000 GeV and 1500 GeV $e^+e^-$ colliders. With more than 500 fb$^{-1}$ the reach is $V > 12.5 \times 10^{-5}$, $V > 4.0 \times 10^{-5}$, and $V > 2.2 \times 10^{-5}$ for the three ascending center-of-mass energies. In the 5DSM, we can translate these reaches in $V$ into reaches of $M_c$ and find 13 TeV, 23 TeV and 31 TeV respectively. These are significantly greater than for a typical grand-unified-theory- (GUT) inspired $Z'$ [85] due to (1) the larger couplings (i.e., $\sqrt{2}$ enhancement), (2) the tower contribution, and (3) both $Z^{(n)}$ and $\gamma^{(n)}$ contribute.

B. Direct searches at hadron colliders

One can also look for direct production of the KK states at hadron colliders. We do not consider the capabilities of indirect, off-shell contributions of the KK excitations to hadron collider observables, since $Z'$ studies indicate that resonant production searches are more probing. The neutral current mode of producing final state lepton pairs in $Z^{(n)}/\gamma^{(n)}$ mediated Drell-Yan processes is the most useful mode to search for evidence of extra dimensions at the Tevatron and LHC.

The scattering processes $f\bar{f} \to l^+l^-$ through intermediate KK excitations $Z^{(n)}$ and $\gamma^{(n)}$ leads to peaks in the invariant mass spectrum at high energies. The couplings of $Z^{(n)}$ and $\gamma^{(n)}$ to fermions are the same as their corresponding zero-mode couplings, except for an overall enhancement of $\sqrt{2}$. In our analysis we estimate the sensitivity to $M_c$ from only the...
first excited state. Incorporating other excited states into the search would yield a slightly higher sensitivity than what we present here.

The search strategy is based on leptonic final states in the narrow width approximation. For the $Z^{(1)}$ first excited state with high mass, the search is for a narrow high-energy dilepton invariant mass excess. For $W^{(1)}$ at high mass, the search is for a high-energy lepton plus large missing energy. For a given luminosity the cuts are chosen such that no standard model background events are expected, and a signal is declared if there are more than 10 events presents given the same cuts. The strategy is summarized in Ref. [85].

We can now present search capabilities for 1 extra dimension, the 5DSM, at the Fermilab Tevatron and CERN Large Hadron Collider (LHC). At the Tevatron with $\sqrt{s}=2$ TeV and integrated luminosity of $2 \text{ fb}^{-1}$ (20 fb$^{-1}$) we find that $M_c$ up to 1.20 TeV (1.36 TeV) can be probed in the combined neutral channels mediated by $\gamma^{(1)}$ and $Z^{(1)}$. At the LHC, in this same channel with $\sqrt{s}=14$ TeV and 100 fb$^{-1}$ integrated luminosity, we find a reach of up to 5.9 TeV. Searches for the $W^{(1)}$ mode at the Tevatron allow discovery at 1.11 TeV and 1.34 TeV with 2 fb$^{-1}$ and 20 fb$^{-1}$ respectively. The $W^{(1)}$ can be discovered at the LHC with 100 fb$^{-1}$ if its mass is less than 6.35 TeV.

One could also look for enhancements in the dijet production at high invariant mass from KK excitations of the gluons. The procedure we employ here is similar to that used to constrain resonant production of squarks in theories with light gluinos [86]. We have extrapolated the CDF and D0 data to higher energies and luminosity, and find a reach capability of $M_c \leq 700$ GeV, which is somewhat lower than the reach capability from $Z^{(1)}$/$\gamma^{(1)}$ and $W^{(1)}$ induced processes. At the LHC we estimate the reach of the first gluon excited level to be below 5 TeV, although the precise number depends sensitively on the dijet energy resolution.

The search reach increases with more extra dimensions because there are more copies of the first KK excitation gauge bosons. For example, with one extra dimension presented above, there is only one copy of $Z^{(1)}$, $\gamma^{(1)}$ and $W^{(1)}$. With $d$ extra dimensions there are $d$ copies of these bosons, making a higher production rate of final state leptons for the same $M_c$. Also, with only one extra dimension the next KK excitation level is at $2M_c$, where only one copy of the gauge bosons resides, whereas with more than one extra dimension the next KK level is lower, $\sqrt{2M_c}$, where there may be many copies of the SM gauge bosons. Therefore, as the number of extra dimensions increases, it appears to become more important to consider the higher KK levels to get an accurate estimate of the maximum search reach. However, as discussed in Sec. II C, the naive effective field theory description of KK excitations cannot be correct, especially for more than one extra dimension, and the couplings of the higher KK states must necessarily be suppressed in a model dependent way. For this reason, we have focussed only on the first excited state.

### V. CONCLUSIONS

In conclusion, the standard model originating from more than four extra dimensions is just as good of a description of nature as the 4D standard model. The difference is in the allowed physical parameters that have not yet been detected. For example, in the ordinary 4DSM, the Higgs boson mass must be less than about 260 GeV in order for the precision electroweak data to match the data well. This is not the case in the 5DSM, where much larger masses (up to 500 GeV) for the Higgs boson are allowed as long as the Higgs boson is confined to the wall and KK excitations of the gauge bosons are rather light.

Table II contains a summary of many of the results presented in the text. All results have been translated into bounds or sensitivity on the compactification scale in the 5DSM, where one can see that current and future colliders will be able to probe well into the TeV region. This is especially relevant for the solution to the hierarchy problem, which we view as the strongest motivations for this scenario. If low-scale compactification theories do have some relevance for electroweak symmetry breaking and the hierarchy problem, it is then at the TeV scale that we expect evidence for them to show up. This is directly analogous to expectations for finding supersymmetry, since low-scale supersymmetry also solves the hierarchy problem. The scale of $M_c$ can then be thought of in the same way as the scale of superpartner masses and the ratio $m_{\chi}^2/M_c^2$ is one measure of how fine-tuned the electroweak potential is. It is for these reasons that we are optimistic that low-scale, sub-Planckian compactifications are more likely at lower scales near $m_w$ than at higher, inaccessible scales.

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[42] T.G. Rizzo, talk given at Physics at RUN II: Workshop on Top