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LIMITS ON THE SIZE OF EXTRA DIMENSIONS[1]

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ABSTRACT

We give a brief summary of present bounds on the size of possible extra-dimensions from collider experiments.

[1] Based on talks given by I.A. at Pascos99 [1] Lake Tahoe, California, 10-16 December 1999 and by K.B. at Rencontres de Moriond [2] on Electroweak Interactions and Unified Theories, Les Arcs, France, March 11-18, 2000.

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1 Introduction

In how many dimensions do we live? Could they be more than the four we are aware of? If so, why don't we see the other dimensions? Is there a way to detect them?. While the possibility of extra-dimensions has been considered by physicists for long time, a compelling reason for their existence has arisen with string theory. It seems that a quantum theory of gravity requires that we live in more than four dimensions, probably in ten or eleven dimensions. The remaining (space-like) six or seven dimensions are hidden to us: observed particles do not propagate in them. The theory does not tell us yet why four and only four have been accessible to us. However, it predicts that this is only a low-energy effect: with increasing energy, particles which propagate in a higher dimensional space could be produced. What is the value of the needed high energy scale? could it be just close by, at reach of near future experiments?

Another scale which appears in our attempts to answer the previous questions is related to the extended nature of fundamental objects. It is the scale at which internal degrees of freedom are excited. In string theory this scale M_s is related to the string tension and sets the mass of the first heavy oscillation mode. The point-like behavior of known particles as observed at present colliders allows to conclude that M_s has to be higher than a few hundred GeV. However to answer the question of what energies should be reached before starting to probe this substructure of the "fundamental particles", more precise determination of experimental lower bounds on M_s and understanding the assumptions behind them is needed.

It is the aim of this talk as to provide a short summary of the present status of limits on these scales of new physics: extra-dimensions and string-like sub-structure of matter.

2 Hiding Extra-Dimensions

There is a simple and elegant way to hide the extra-dimensions: compactification. It is simple because it relies on an elementary observation. Suppose that the extra-dimensions form, at each point of our four-dimensional space, a D-dimensional torus of volume $(2\pi)^D R_1 R_2 \cdots R_D$. The $(4+D)$-dimensional Poincare invariance is replaced by a four-dimensional one times the symmetry group of the D-dimensional space which contains translations along the D extra directions. The $(4+D)$-dimensional momentum satisfies the mass-shell condition $P^2_{(4+D)} = p_0^2 - p_1^2 - p_2^2 - p_3^2 - \sum_i p_i^2 = m_0^2$ and looks from the four-dimensional point of view as a (squared) mass $M_{KK}^2 = p_0^2 - p_1^2 - p_2^2 - p_3^2 = m_0^2 + \sum_i p_i^2$. Assuming periodicity of the wave functions along each compact direction, one has $p_i = n_i/R_i$ which leads to:

$$M_{KK}^2 \equiv M_{\vec{n}}^2 = m_0^2 + \frac{n_1^2}{R_1^2} + \frac{n_2^2}{R_2^2} + \cdots + \frac{n_D^2}{R_D^2}\,, \tag{1}$$

with m_0 the four-dimensional mass and $n_i 0$ non-negative integers. The states with $\sum_i n_i \neq 0$ are called Kaluza-Klein (KK) states. It is clear that getting aware of the ith extra-dimension would require experiments that probe at least an energy of the order of $min(1/R_i)$ with sizable couplings of the KK states to four-dimensional matter.

Let us discuss further some properties of the KK states that will be useful for us below. We parametrise the "internal" D-dimensional box by $y_i \in [-\pi R_i, \pi R_i]$, $i = 1, \cdots, D$ while the four-dimensional Minkowski spacetime is spanned by the coordinates x^μ, $\mu = 0, \cdots 3$. It is useful to choose for the KK wave functions the basis:

$$\Phi^\alpha_{\vec{n},\vec{e}}(x^\mu, y_i) = \Phi^\alpha(x^\mu) \prod_i \left[(1-e_i)\cos(\frac{n_i y_i}{R_i}) + e_i \sin(\frac{n_i y_i}{R_i})\right]\,, \tag{2}$$

where the vector $\vec{n} = (n_1, n_2, \cdots, n_D)$ gives the energy of the state following Eq. 1 while $\vec{e} = (e_1, \cdots, e_D)$ with $e_i = 0$ or 1 correspond to a choice of cosine or sine dependence in the coordinate y_i, respectively. The index α refers to other quantum numbers of Φ.

The simplest example of the models we will be using for getting experimental bounds are obtained by gauging the Z_2 parity: $y_i \to -y_i \mod 2\pi R_i$. This leads to compactification on segments of size πR_i. In general, the consistency of this "orbifold" projection implies that the Z_2 space parity should be associated with a Z_2 action on the internal quantum numbers α of Φ. As a result one has the following properties:

- Only states invariant under this Z_2 are kept while the others are projected out. There are two classes of states left in the theory: those for which $\Phi^{(even)}(x^\mu)$ is even under Z_2 action and $e_i = 0$ and those for which $\Phi^{(odd)}(x^\mu)$ is odd and $e_i = 1$. It is important to notice that the latter are not present as light four-dimensional states i.e. they have $\sum_i n_i \neq 0$ and thus always correspond to higher KK states.

- At the boundaries $y_i = 0, \pi R$ fixed by the Z_2 action, new states $\Phi^{(loc)}(x^\mu)$, have to be included. These "twisted" states are localized at the fixed points. They can not propagate in the extra-dimension and thus have no KK excitations.

- The odd bulk states $\Phi^{(odd)}(x^\mu)$ ($e_i = 1$) have a wave function which vanishes (the $\sin(\frac{n_i y_i}{R_i})$ in Eq. 1) at the boundaries. Their coupling to localized states involves a derivative along y_i. For example three boson interactions of the form $\partial_i \phi^{(odd))} \phi^{(loc)} \phi^{(loc)}$ can be non-vanishing.

- The even states, in contrast, can have non-derivative couplings to localized states. The gauge couplings for instance are given by:

$$g_n = \sqrt{2} \delta^{-|\frac{\vec{n}}{R}|^2/M_s^2} g \tag{3}$$

where $\delta > 1$ is a model dependent number ($\delta = 1/2$ in the case of Z_2). The $\sqrt{2}$ comes from the relative normalization of $\cos(\frac{n_i y_i}{R_i})$ wave function with respect to the zero mode while the exponential damping is a result of tree-level string computations that we do not present here.

Use of compactification is an elegant way to hide extra-dimensions because some of the quantum numbers and interactions of the elementary particles could be accounted to by the topological and geometrical properties of the internal space. For instance chirality, number of families in the standard model, gauge and supersymmetry breaking as well as as some selection rules in the interactions of light states could be reproduced through judicious choice of more complicated internal spaces.

3 Theoretical constraints

The basic requirement on the theoretical side is that there exist theories that allow the correct magnitude for the strength of the gauge and gravitational couplings for given compactification and string scales just above the present experimental energies. In the simplest string models, the four-dimensional Planck mass can be expressed as:

$$M_{pl}^2 \equiv f_{pl} \frac{(M_s^D V_D)}{g_s^p} M_s^2 \,, \tag{4}$$

where V_D is the D-dimensional internal volume felt by gravitational interactions, g_s the string coupling and p an integer. The four-dimensional gauge coupling can be written as

$$\frac{1}{g_{YM}^2} \equiv f_{YM} \frac{(M_s^d V_d)}{g_s^q} \,, \tag{5}$$

where V_d is the d-dimensional internal volume felt by gauge interactions, and the coefficients f_{pl}, f_{YM} have been computed for known classical string vacua. In the lowest order approximation, they are moduli-independent $\mathcal{O}(1)$ constants.

In the past, weakly coupled heterotic strings were providing the most promising framework for phenomenological applications. In this case, the standard model was considered as descending from the ten-dimensional E_8 gauge symmetry, and we have $V_d = V_D$, $D = d = 6$ and $p = q = 2$. Taking the ratio of the two equations, one finds $\frac{M_s^2}{M_{pl}^2} = \frac{f_{YM}}{f_{pl}} g_{YM}^2 \sim g_{YM}^2$. Requiring $g_{YM} \sim \mathcal{O}(1)$, it was concluded that both the string scale M_s and the compactification scale $R^{-1} \equiv V_6^{-1/6}$ had to lie just below the Planck scale, at energies $\sim 10^{18}$GeV far out of reach of any near future experiment [3, 4].

The situation changed during recent years [5] when it was discovered that string theory provides classical solutions (vacua) where gauge degrees of freedom live on subspaces i.e. $d < D$ along with the possibility of $p \neq q$. For instance, while $D = 6$ and $p = 2$, $(d, q) = (d, 1)$ in type I and $(d, q) = (2, 0)$ in type II or weakly coupled heterotic strings with small instantons. In these cases, it is an easy exercise to check that both the string and compactification scales can be made arbitrarily low [4, 6].

Lowering the string scale, one increases the strength of higher (non-renormalizable) operators leading to the possibility of inducing exotic processes at experimentally excluded rates. Although an explicit string realization of the scenario is necessary in order to have a satisfactory solution, at the effective field theory level many discrete or global symmetries can be displayed that forbid these operators.

4 Experimental constraints

4.1 The scenario:

In order to pursue further, we need to provide the quantum numbers and couplings of the relevant light states. In the scenario we consider:

- Gravitons [3] which describe fluctuations of the metric propagate in the whole 10- or 11-dimensional space.

- In all generality, the gauge bosons propagate on a $(3 + d)$-brane, with $d = 0, ..., 6$. However, as we have seen in the previous section, a freedom of choice for the values of the string and compactification scales requires that gravity and gauge degrees of

[3] Along with gravitons, string models predict the presence of other very weakly coupled states as gravitinos, dilatons, moduli, Ramond-Ramond fields....These might alter the bounds obtained in Section 4.3.

freedom live in spaces with different dimensionalities. This means that $d_{max} = 5$ or 6 for 10- or 11-dimensional theories, respectively. The value of d represents the number of dimensions felt by KK excitations of gauge bosons.

- The matter fermions, quarks and leptons, are localized on a 3-brane and have no KK excitations. Their coupling to KK modes of gauge bosons are given in Eq. 3. This is the main assumption in our analysis and limits derived in the next subsection depend on it. In a more general study it could be relaxed by assuming that only part of the fermions are localized. However, if all states are propagating in the bulk, then the KK excitations are stable and a discussion of the cosmology will be necessary in order to explain why they have not been seen as isotopes.

The possible localization of the Higgs scalar, as well as the possible existence of supersymmetric partners do not lead to important modifications for most of the obtained bounds.

4.2 Extra-dimensions along the world brane: KK excitations of gauge bosons

To simplify the discussion, let us first consider the case $d = 1$ where some of the gauge fields arise from a 4-brane. Since the couplings of the corresponding gauge groups are reduced by the size of the large dimension $R_{\|} M_s$ compared to the others, if $SU(3)$ has KK modes all three group factors must have. Otherwise it is difficult to reconcile the suppression of the strong coupling at the string scale with the observed reverse situation. As a result, there are 5 distinct cases that we denote (l, l, l), (t, l, l), (t, l, t), (t, t, l) and (t, t, t), where the three positions in the brackets correspond to the 3 gauge group factors of the standard model $SU(3)_c \times SU(2)_w \times U(1)_Y$ and those with l feel the extra-dimension, while those with t (transverse) do not.

The experimental signatures of extra-dimensions are of two types:

- Observation of resonances due to KK excitations. This needs a collider energy $\sqrt{s} \gtrsim 1/R_{\|}$ at LHC. The discovery limits in the case of one extra-dimension are given in table 1.

- Virtual exchange of the KK excitations which lead to measurable deviations in cross-sections compared to the standard model prediction. The exchange of KK states gives

Table 1: Limits on $R_\parallel^{-1}$ in TeV at present and future colliders. The luminosity is given in fb^{-1}.

Collider	Luminosity	Gluons	$W^\pm$	$\gamma + Z$
Discovery of Resonances				
LHC	100	5	6	6
Observation of Deviation				
LEP 200	4×200	-	-	1.9
TevatronI	0.11	-	-	0.9
TevatronII	2	-	-	1.2
TevatronII	20	4	-	1.3
LHC	10	15	8.2	6.7
LHC	100	20	14	12
NLC500	75	-	-	8
NLC1000	200	-	-	13

rise to an effective operator:

$$\bar{\psi}_1 \psi_2 \bar{\psi}_3 \psi_4 \sum_{|\vec{n}|} \frac{g^2(|\vec{n}|)}{m_0^2 + \frac{|\vec{n}|^2}{R_\parallel^2}} . \tag{6}$$

The usual approximation of taking $g^2(|\vec{n}|)$ independent of $|\vec{n}|$ fails for more than one dimension because the sum $\sum_{n_i} \frac{1}{n_1^2+n_2^2+\ldots}$ becomes divergent. This divergence is regularized by the exponential damping of Eq. 3. For $d > 1$ the result depends then on both parameters $R_\parallel$ and M_s. Example of analysis for $d = 2$ can be found in Ref. [7]. The simpler case of $d = 1$ has been studied in detail. Possible reaches of colliders experiments [8, 7] are summarized in table 1.

The effects of exchange of virtual KK modes are also constrained by high precision data [9, 10], such as the fit of the measured values of M_W, Γ_{ll} and Γ_{had}. If the Higgs is assumed to be a bulk state like the gauge bosons, then one finds $R^{-1} \gtrsim 3.5$ TeV. Inclusion of Q_W measurement, which does not give a good agreement with the standard model itself, raises the bound to $R^{-1} \gtrsim 3.9$ TeV [10]. The presence of a localized Higgs allows tree-level mixing of different KK, and makes the bounds model dependent [10].

There are some ways to distinguish the corresponding signals from other possible origin of new physics, such as models with new gauge bosons. In the case of observation of resonances, one expects three resonances in the (l,l,l) case and two in the (t,l,l) and (t,l,t) cases, located practically at the same mass value. This property is not shared by most of other new gauge boson models. Moreover, the heights and widths of the resonances are directly related to those of standard model gauge bosons in the corresponding channels. In the case of virtual effects, these are not reproduced by a tail of Bright-Wigner shape and a deep is expected just before the resonance of the photon+Z, due to the interference between the two. However, good statistics will be necessary.

4.3 Extra-dimensions transverse to the brane world: KK excitations of gravitons

The localization of (infinitely massive) branes in the $(D-d)$ dimensions breaks translation invariance along these directions. Thus, the corresponding momenta are not conserved: particles, as gravitons, could be absorbed or emitted from the brane into the $(D-d)$ dimensions. Non observation of the effects of such processes allow us to get bounds on the size of these transverse extra dimensions. In order to simplify the analysis, it is usually assumed that among the $D-d$ dimensions n have very large common radius $R_\perp \gg M_s^{-1}$, while the remaining $D-d-n$ have sizes of the order of the string length.

During a collision of center of mass energy $\sqrt{s}$, there are $(\sqrt{s}R_\perp)^n$ KK excitations of gravitons with mass $m_{KK\perp} < \sqrt{s} < M_s$, which can be emitted. Each of these states looks from the four-dimensional point of view as a massive, quasi-stable, extremely weakly coupled (s/M_{pl}^2 suppressed) particle that escapes from the detector. The total effect is a missing-energy cross section roughly of order:

$$\frac{(\sqrt{s}R_\perp)^n}{M_{pl}^2} \sim \frac{1}{s}(\frac{\sqrt{s}}{M_s})^{n+2} \qquad (7)$$

Explicit computation of these effects leads to the bounds given in table 2 [11]. The results require some remarks:

- The amplitude for emission of each of the KK gravitons is taken to be well approximated by the tree-level coupling of the massless graviton as derived from General Relativity. Eq. 3 suggests that this is likely to be a good approximation for $R_\perp M_s \gg 1$.

- The cross-section depends on the size $R_\perp$ of the transverse dimensions and allows to derive bounds on this *physical* scale. As it can be seen from Eq. 3, transforming these bounds to limits on M_s there is an ambiguity on different factors involved, such as the string coupling. This is sometimes absorbed in the so called "fundamental quantum gravity scale $M_{(4+n)}$". Generically $M_{(4+n)}$ is bigger than M_s, and in some cases, as in type II strings or in heterotic strings with small instantons, it can be many orders of magnitude higher than M_s, so it does not correspond to a scale where some physical phenomena open up.

- There is a particular energy and angular distribution of the produced gravitons that arises from the distribution in mass of KK states given in Eq. 1. It might be a smoking gun for the extra-dimensional nature of such observable signal.

- For given value of M_s, the cross section for graviton emission decreases with the number of large transverse dimensions. The effects are more likely to be observed for the lowest values of M_s and n.

- Finally, while the obtained bounds for $R_\perp^{-1}$ are smaller than those that could be checked in table-top experiments probing macroscopic gravity at small distances, one should keep in mind that larger radii are allowed if one relaxes the assumption of isotropy, by taking for instance two large dimensions with different radii.

In table 2, we have also included astrophysical and cosmological bounds. Astrophysical bounds [13, 14] arise from the requirement that the radiation of gravitons should not carry on too much of the gravitational binding energy released during core collapse of supernovae. In fact, the measurements of Kamiokande and IMB for SN1987A suggest that the main channel is neutrino fluxes.

The best cosmological bound [15] is obtained from requiring that decay of bulk gravitons to photons do not generate a spike in the energy spectrum of the photon background measured by the COMPTEL instrument. The bulk gravitons are themselves expected to be produced just before nucleosynthesis due to thermal radiation from the brane. The limits assume that the temperature was at most 1 MeV as nucleosynthesis begins, and become stronger if the temperature is increased.

Table 2: Limits on $R_\perp$ in mm from missing-energy processes.

Experiment	$R_\perp(n=2)$	$R_\perp(n=4)$	$R_\perp(n=6)$
Collider bounds			
LEP 2	4.8×10^{-1}	1.9×10^{-8}	6.8×10^{-11}
Tevatron	5.5×10^{-1}	1.4×10^{-8}	4.1×10^{-11}
LHC	4.5×10^{-3}	5.6×10^{-10}	2.7×10^{-12}
NLC	1.2×10^{-2}	1.2×10^{-9}	6.5×10^{-12}
Present non-collider bounds			
SN1987A	3×10^{-4}	1×10^{-8}	6×10^{-10}
COMPTEL	5×10^{-5}	-	-

4.4 Dimension-Eight Operators and Limits on The String Scale:

At low energies, the interaction of light (string) states is described by an effective field theory. Non-renormalizable dimension-six operators are due to the exchange of KK excitations of gauge bosons between localized states. If these are absent, then there are deviations to the standard model expectations from dimension-eight operators. There are two generic sources for such operators: exchange of virtual KK excitations of bulk fields (gravitons,...) and form factors due to the extended nature of strings.

The exchange of virtual KK excitations of bulk gravitons is described in the effective field theory by an amplitude involving the sum $\frac{1}{M_p^2}\sum_n \frac{1}{s-\frac{\vec{n}^2}{R_\perp^2}}$. For $n>1$, this sum diverges and one cannot compute it in field theory but it only in a fundamental (string) theory. In analogy with the case of exchange of gauge bosons, one expects the string scale to act as a cut-off with a result:

$$A g_s^2 \frac{T_{\mu\nu}T^{\mu\nu} - \frac{1}{1+d_\perp} T_\mu^\mu T_\nu^\nu}{M_s^4} \,. \tag{8}$$

The approximation $A = \log\frac{M_s^2}{s}$ for $d_\perp = 2$ and $A = \frac{2}{d_\perp - 2}$ for $d_\perp > 2$ is usually used for quantitative discussions. There are some reasons which might invalidate this approximation for particular cases. In fact, the result is very much model dependent: in type I string models it reflects the ultraviolet behavior of open string one-loop diagrams which are model (compactification) dependent.

In order to understand better this issue, it is important to remind that in type I string models, gravitons and other bulk particles correspond to excitations of closed strings. Their tree-level exchange is described by a cylinder joining the initial $|Bin>$ and final $|Bout>$ closed strings lying on the brane. This cylinder can be be seen on the other hand as an open string with one of its end-points describing the closed (loop) string $|Bin>$, while the other end draws $|Bout>$. In other words, the cylinder can be seen as an annulus which is a one-loop diagram of open strings with boundaries $|Bin>$ and $|Bout>$. Note that the validity of this duality, which only assumes the branes to be Dirichelet branes of string theory, seems to require the presence of other weakly interacting closed strings besides gravitons.

More important is that when the gauge degrees of freedom arise from Dirichelet branes, it is expected that the dominant source of dimension-eight operators is not the exchange of KK states but instead the effects of massive open string oscillators [8, 12]. These give rise to contributions to tree-level scattering that behave as $g_s s/M_s^4$. Thus, they are enhanced by a string-loop factor g_s^{-1} compared to the field theory estimate based on KK graviton exchanges. Although the precise value of g_s requires a detail analysis of threshold corrections, a rough estimate can be obtained by taking $g_s \simeq \alpha \sim 1/25$, implying an enhancement by one order of magnitude.

What is the simplest thing one could do in practice?. There are some processes for which there is only one allowed dimension-eight operator; an example is $f\bar{f} \to \gamma\gamma$. The coefficient of this operator can then be computed in terms of g_s and M_s. As a result, in the only framework where computation of such operators is possible to carry out, one cannot rely on the effects of exchange of KK graviton excitations in order to derive bounds on extra-dimensions or the string scale. Instead, one can use the dimension-eight operator arising from stringy form-factors.

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