Gravity particles from Warped Extra Dimensions, a review.  
Part I - KK Graviton.

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Abstract

On face of the latest LHC experimental results on the direct searches for Beyond Standard Model physics we review the basic of Warped Extra Dimensions scenarios and the physics of the heavy gravity particles, their most unique signature.

In this first part we summarize the physics behind the hypothesis of a heavy spin-2 Beyond Standard Model particle on the context of WED and also to address some of the interesting phenomenology issues of model building hypothesis.

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Introduction

Among the five known elementary forces of nature four are described by Standard Model of elementary particles (electromagnetic, weak, strong and Higgs interactions). Gravitational interaction remains a odd personage - the (existence of a) quantum description of gravitational field is still a live and turbulent field of research.

The gravitational potential between two objects $M_1$ and $M_2$ at a given distance $r$ is.

$$V(r; M_1, M_2) = \frac{1}{M_{Pl}^2} \frac{M_1 M_2}{r}$$

The mass scale that weakens the gravitational interaction is known as Planck mass. To simplify notation on the course of calculations it is usual to define the reduced Planck mass as $M_{P} = M_{Pl}/\sqrt{8\pi}$ GeV. The value of reduced Planck mass is $M_{P} \sim 2.4 \times 10^{18}$ while the natural mass scale for the mediators of weak force is $M_{weak} \sim 100$ GeV.

On the course of last century it was realized that the existence of extra spatial dimensions (ED) – where gravitational interaction could propagate modify the gravity behaviour at short distances. The $n$–dimensional gravitational scale approaches the $n$–dimensional weak scale – justifying the weakness of gravity interactions with respect the remaining forces of nature we experience on our four dimensional effective world – the $M_{Pl}$ value would not be a fundamental scale of nature anymore.[1, 2, 3, 4, 5].
The original proposal for one Warped Extra Dimension (WED) was made on [6,7]. The focus of these works were to show a mechanism where the effective difference between $M_{Pl}$ and $M_{weak}$ is an outcome of the existence of one finite additional spatial dimension with non trivial metric. The five dimensional space defined between the two branes is referred as the bulk.

Quantum fluctuations of the five dimensional space-time metric are interpreted as particles, the fluctuations around the $3+1$ infinite part of metric correspond to the five dimensional Graviton field (spin-2) while the fluctuations around the finite size of the spatial extra dimension manifest itself as a spin-0 field, common called Radion [8].

The finiteness of the extra spatial dimension makes the four dimensional effective theory (our world) to contain excitation modes, acting as heavy resonances, known as Kaluza-Klein (kk) modes. The mass properties of the kk modes are directly calculable from the space-time geometry. The mass of the first kk-mode of Graviton field (spin-2) is expected to be on TeV range if the size and curvature of the extra dimension is chosen such that the Planck-TeV hierarchy is solved on the full theory.

On the original setup the SM matter is localized on the same brane of the Higgs doublet. The possibility of matter fields to be allowed to propagate along the extra dimension under the space-time setup [7] started to be investigated just some time after their publication. This hypothesis opened a big window for particle phenomenology: to promote the SM matter fields to be 5-dimensional introduces heavy kk-modes for all the matter fields on the effective 4 dimensional theory [9,10].

Indirect consequences of the above models were studied by many authors, in particular [11,12,13,14].

Deviations of the WED scenario, that holds similar phenomenology but not so strict structure for kk-particles couplings and masses can be done using solely effective field theory. One can adopt a more generic approach, based holographic principles [15,16].

One can also consider constructions were instead of a continuous dimension one considers a product of finite number of copies of space-time gauge with hierarchical gauge couplings, so called sites [17,18], this scenario is also known as "Dimensional Desconstruction" (or, by its old fashioned name "moose models" [19,20]). This naming is used because we recuperate the theory described by a continuous extra spacial dimension in the limit of infinite sites.

On the last years the searches for heavy particles on 8 TeV LHC data had consistently advanced on both CMS and ATLAS. Several channels were already probed on di-jet [21], di-vector boson [22], di-leptons [23]. The di-top final state remains a challenge both from the difficulty of tag a boosted top and big $t\bar{t}$ QCD BKG, even in high mass region [24,25,26,27].

Since the celebrated ($\sim 125$ GeV) higgs discovery [28,29] the experimental searches for heavy resonances on LHC are being extended to channels with final states that contain higgs particles. The ultimate signature to define the possible signal as a spin even particle is to find its evidence also on di-higgs final state. The first results of this channel can be already found on [30].
On this first part of review we focus on the physics of heavy kk-graviton on the universe scenario of [7]. Our main purpose is to highlight similarities and differences on kk-graviton phenomenology provoked on consider bulk or RS1 scenario and to mention the experimental techniques that allows to differentiate the scenarios on particular LHC searches. The intention is to keep attached to the most simple models that still can be predictive and point the model dependent hypothesis of those, sharing intentions with the recently proposed Bridge model [31].

We present this document with the follow structure: a mini review of the motivations of for WED scenarios is found on section (I). Section (II) introduces the dynamics of Graviton field and Standard Model matter fields. On section (IV) we present the couplings between kk-graviton and matter fields. The following section (V) contains the results we collected. We finalize on section (VI), with some highlights of kk-graviton phenomenology resultant of the structure of kk-graviton couplings to Higgs sector.

I WED and the problem of hierarchy of scales

The proposition of a non factorizable geometry with one small spatial extra dimension was made on [7], the authors exploited a universe scenario with one compact Extra Dimension, the compactification scheme of this dimension allows us to be describe the ED as a line segment between two four dimensional branes, known as Planck and TeV brane (see figure I.1). The presence of a negative five dimensional cosmological constant ($\Lambda_{bulk}$) and tensions between the two branes (with the strength $\Lambda_{TeV/Planck}$) produces a exponential warp factor along the compactifyed Extra Dimension. If the Standard Model Higgs doublet is localized on TeV brane the exponential warp naturally address the hierarchy problem.

To state the above paragraph on mathematical terms we set off the most general solution form that solves the classical Einstein’s motion equations and holds 4D Poincaré invariance:

$$ds^2 = e^{-2\sigma(\phi)}\eta_{\mu\nu}dx^\mu dx^\nu + r_c^2d\phi^2 \equiv g_{MN}dx^M dx^N \quad (I.1)$$

Where $\mu\nu = 1, ..., 4$ are the 4D indexes. The full classical action reads:

$$S = S_{Gravity} + S_{TeV} + S_{Planck} + S_{Matter}, \quad (I.2)$$

Where $S_{Gravity}$ is the bulk gravitational action, $S_{Matter}$ is the action matter fields, we consider separately from $S_{Gravity}$ the pieces of gravitational action confined on the branes, we denoted them by $S_{TeV/Planck}$. The gravitational part of the action can be written as:

$$S_{i=TeV/Planck} = -\int d^4x\sqrt{g(\phi = 0, \pi)}\Lambda_{i=TeV/Planck} \quad (I.3a)$$

$$S_{Gravity} = \int d^4x\int_{-\pi}^{\pi}d\phi\sqrt{g}(-\Lambda_{bulk} + 2M_5^3R) \quad (I.3b)$$
Figure I.1: Scheme of dimensions on RS theory. The Gravity (Planck) and Weak (TeV) branes are the 4 dimensional boundaries of the extra dimension $\phi$ compactified on a interval $(0, \pi)$. The figure also illustrate the metric behaviour along the extra dimension.

Where the $\Lambda$’s are vacuum energy densities on the branes/bulk, $R$ is the Ricci tensor, $M_5$ is the 5D Planck mass and $g_{XY}$ is the 5D metric. If we denote the trace of the space-time metric as $g = -g^X_X$ the unity of volume reads $(-\int d^5x \sqrt{g})$.

The 5D Einstein gravity equations in terms of the Ricci tensor and scalar ($R_{MN}$ and $R \equiv R^M_M$) are:

$$\sqrt{g} \left( R_{MN} - \frac{1}{2} g_{MN} R \right) = T_{MN}$$

See for example [32] and references within. The indexes $M,N = 1,\ldots,5$ are the 5D indexes, we use the sign convention $(-,+,+,+,+)$. The most general solution form to solve the classical Einstein motion equations maintaining 4D Poincareé invariance is:

$$ds^2 = e^{-2\sigma(\phi)} \eta_{\mu\nu} dx^\mu dx^\nu + r^2 d\phi^2$$

(I.4)

Where $\mu\nu = 1,\ldots,4$ are the 4D indexes. Using (I.4) the Ricci tensor and scalar are given by:

$$R_{\mu\nu} = e^{-2k(\phi)} (4\sigma(\phi)^2 + \sigma(\phi)''') \eta_{\mu\nu}$$

(I.5a)

$$R_{55} = -4\sigma(\phi)^2 - 4\sigma(\phi)''$$

(I.5b)

The primes denotes derivations with respect to the $\phi$ variable. If we neglect the matter back-reaction on the Einstein gravity equations one finds the function $\sigma(\phi)$ is linear on
the extra dimension, this function acts as an exponential warp factor along the compact extra dimension. Taking $\Lambda_{\text{bulk}} = \Lambda_{\text{Planck}} = -\Lambda_{\text{TeV}} = \Lambda$ we arrive at [7]:

$$\sigma(\phi) = r_c |\phi| \sqrt{-\frac{\Lambda}{24M_5^2}} \equiv r_c |\phi| k$$

(I.6)

Literature refers to the factor $k \equiv \sqrt{-\frac{\Lambda}{24M_5^2}}$ as the curvature factor, this form strongly depends on the relation between the negative cosmological constant and the brane tensions.

Integrating out the extra dimension we find the four dimensional effective gravity action, and consequently the four dimensional Planck mass:

$$M_{\text{Pl}}^2 = M_5^3 \frac{1}{k} (1 - e^{-2\pi kr_c})$$

(I.7)

Usually the curvature factor $k$ is assumed to be of the order of 5D Planck scale $M_5$, like this any large value of $kr_c > 1$ the equation (I.7) does not produce strong hierarchy between the mass constants of the theory ($k$, $M_5$ and $M_{\text{Pl}}$).

Both $k$ and $r_c$ are free parameters of the theory, probes on the range of validity of (I.1) constrain the value of the compactification radius $r_c$. To do not deal with extremely big numbers we use the combinations $\tilde{k} \equiv \frac{k}{M_{\text{Pl}}}$ and $kr_c$ as basic parameters of the theory.

On next section we discuss the relation of $r_c k$ value and the hierarchy of Weak and Gravitational mass scales.

I.1 The higgs mechanism

The higgs sector is added to the theory just like one its SM construction: as a $SU(2)_L \times U(1)_Y$ scalar doublet, we also admit the existence of a quartic potential. The Higgs mass is assumed to be the same 5D mass considered on Gravitational action (I.3b) $M_5$.

If we confine the higgs doublet ($H$) to live on the TeV brane its 5D action is:

$$S_H = - \int d^5 x \sqrt{-g} \left[ \partial_\mu H^\dagger \partial^\mu H - \lambda(|H|^2 + v_0^2)^2 \right] \delta(\phi r_c - \pi r_c),$$

(I.8)

when we integrate out the extra dimension ($\phi$) we arrive to

$$S_H = - \int d^4 x \left[ e^{-2\pi kr_c} \partial_\mu H^\dagger \partial^\mu H - e^{-4\pi kr_c} \lambda(|H|^2 + v_0^2)^2 \right].$$

(I.9)

We canonize the kinetic term redefining $H \rightarrow e^{-\pi kr_c} H$, the 4 dimensional higgs action reads on equation (I.10):

$$S_H = - \int d^4 x \left[ \partial_\mu H^\dagger \partial^\mu H - \lambda(|H|^2 + e^{-2\pi kr_c} v_0^2)^2 \right].$$

(I.10)

We define the four dimensional vacuum expectation value of the Higgs doublet as:

$$v \equiv e^{-\pi kr_c} v_0$$

(I.11)
If we assume the 5D Higgs vev \( v_0 \) of the order of the 5D fundamental mass scale \( M_5 \) the separation between Planck (gravity) and EW (higgs mass) scale is assured when \( kr_c \sim 11 \).

The Planck-EW hierarchy reduction by the exponential space-time warp factor, represented by equation (I.11), is the most celebrated model feature of Warped Extra Dimensional scenarios.

Other constructions admit the higgs doublet could be a composite state of the heavy \( kk \) fermions, bounded by an additional strong force \([33, 34, 35, 36]\), or even by excitations of QCD \([37, 38]\).

## II Gravity particles

We refer as "gravity particles" the particles resultant of the quantum fluctuations around the classical metric solution (I.6). The fluctuation modes can be decomposed into a 4D tensor \( \otimes \) 4D vector \( \otimes \) scalar components:

\[
\delta g_{MN}(x, \phi) = \begin{pmatrix}
    h_{\mu\nu}(x, \phi) & h_{\mu,5}(x, \phi) \\
    h_{\mu,5}(x, \phi) & h_{55}(x, \phi)
\end{pmatrix}
\]  

(II.1)

The motion’s relations for the independent fields of (II.1) form a over constrained system of equations. The axial gauge, defined by fix \( h_{\mu,5} = 0 \), decouples the dynamics of tensor and scalar perturbations. The perturbed space-time metric on this gauge have the form:

\[
ds^2 = e^{-2(\sigma + \sigma_{\phi})}\left(\eta_{\mu\nu} + h_{\mu\nu}(x, \phi)\right) dx^\mu dx^\nu + (2h_{55}(x, \phi) - r_c)^2 d\phi^2
\]

(II.2)

The tensor fluctuations \( (h_{\mu\nu}) \) correspond to Graviton modes and the scalar fluctuation \( (h_{55}) \) correspond to Radion mode \([39, 40, 41]\).

We find the dynamics and interactions of the gravity particles substituting the form (II.2) on the full action (I.2):

\[
S_{\text{Gravity}} = S_{\text{Graviton}} + S_{\text{Radion}}
\]

(II.3)

The five dimensional motion’s equations are solved performing a Fourier expansion on the fields five dimensional wave function that separates the \( \phi \) and \( x_\mu \) behaviours as equation (II.4).

\[
X(x^\mu, \phi) = \sum_{n=0}^{\infty} X^{(n)}(x^\mu) f_X^{(n)}(\phi)
\]

(II.4)

Each four dimensional expansion mode is also known as the \( n \)-th Kaluza-Klein (KK) mode. The 5th dimensional part of kk-modes wave functions \( f_X^{(n)}(\phi) \) are commonly called profiles. The profiles \( f_X^{(n)} \) obey a differential Sturm-Liouville problem defined by the equations of motion, spin of the field and boundary conditions, the profiles for all the known
fields are described by a combination of exponential and Bessel functions. A didactic and complete guide to the behaviour of matter fields on bulk WED scenario can be found on [42]. It is out of scope of this note to write down and solve the 5D equations of motion to derive in details each matter field 5D profile, here we just highlight the results necessary for further discussion.

We denote the Graviton four dimensional wave functions (and profiles) as $h^{(n)}_{\mu\nu}(x_{\mu})$ (and $f^{(n)}_{h}(\phi)$). The zero mode Graviton correspond to the traditional graviton field -mediator of gravitational force. The equations of motion defines the kk-graviton profiles as:

$$f^{(n)}_{h}(\phi) = N_n e^{2k_{\phi}r_c} \left[ J_2 \left( \frac{m_{Gr}^{(n)}}{k e^{-k_{\phi}r_c}} \right) + Y_2 \left( \frac{m_{Gr}^{(n)}}{k e^{-k_{\phi}r_c}} \right) \right] , \quad (II.5)$$

where $J_2$ ($Y_2$) are Bessel (Neumann) functions of second rank, $N$ is a normalization constant and $m_{Gr}^{(n)}$ are the four dimensional mass for each mode.

The first kk mode is TeV localized, its 4D effective mass is:

$$m_{Gr} \equiv m_{Gr}^{(1)} = x_1 \frac{\bar{k}}{k} e^{-k_{\phi}r_c} M_{Pl} \sim \text{TeV}, \quad (II.6)$$

where $x_1 = 3.83$ is the first zero of the Bessel function $J_1$. If we want to use $m_{Gr}$ and $\bar{k}$ as the free parameter of theory we must note the relation:

$$k_{r_c} = \frac{1}{\pi} \ln \left( x_1 \frac{\bar{M}_{Pl}}{M_{Gr}} \right) \quad (II.7)$$

The value of $k_{r_c}$ parameter is dominated by $\bar{M}_{Pl}$. The dependence on the parameters $M_{Gr}$ and $\bar{k}$ is very mild. Following the relation of equation (II.7) the $k_{r_c}$ value varies between 10 and 12 if $k$ vary between 0.01 and 1, and $M_{Gr}$ vary between 100 GeV and 1.5 TeV. Remaining a phenomenologically acceptable value to reproduce the mass hierarchy of equation (I.11) on a wide range of conditions.

The physics of the scalar fluctuations (Radion) requires an external mechanism to stabilize the distance $r_c$ between branes this mechanism was first derived on [43], the references [8, 44] among others exploit its collider phenomenology. On this part of the review we treat the 4D graviton spectra.

### III The Standard Model fields

For definitiveness we will refer to WED case with matter fields allowed to propagate on the extra dimension as bulk scenario, for the case that considers all the SM fields to be confined on TeV brane as RS1 scenario.

On the simplest bulk WED construction only the SM gauge symmetry $SU(2)_L \times U(1)_Y$ is enlarged to the bulk. The 5D action for matter fields on the bulk model reads [10]:
\[ S_{\text{Matter}} = -\int d^5x \sqrt{-g} \left( \frac{1}{4} F_{MN} F^{MN} + i \bar{\Psi} \gamma^M D_M \Psi + i m_f \bar{\Psi} \Psi \right) \]  

(III.1)

Where \( F_{MN} = \partial_M A_N - \partial_N A_M + ig_5 [A_M, A_N] \) represents the SM gauge fields field strengths and \( \Psi \) the SM fermion field \( i \), where the fermionic basis is already diagonalized. The bulk fermion fields admits a five dimensional mass term \( m_f \). Usually literature calculates the masses of the kk spectrum ignoring the back-reaction of bulk matter on space-time metric\(^1\).

As pointed on the last chapter the masses of the kk modes of each bulk fields are geometrically generated, and therefore directly calculable from space-time geometry, the masses for the zero modes however are free of constrains and can always be taken as massless, to correspond to the SM fields before Electroweak Symmetry Breaking.

### III.1 Bulk fermions

The five dimensional mass term \( m_f \) results on one additional degree of freedom to define fermionic profiles. We find the fermionic kk modes to be vectorlike, each kk mode have both left and right handed chirality.

The boundaries conditions assures only one of the chiralities survive for the massless zero mode. To construct a massive four dimensional field we need to introduce two bulk fermions – to fix notation we designates the bulk fermion field where left (right) handed zero mode survive as left (right) handed fermion. The zero mode fermion profiles to a right/left handed fermion reads:

\[
f_{L/R}^{(0)}(\phi) = \sqrt{e^{(1 \mp 2 c_{f_{L/R}})k} - 1} e^{- (1 + 2 c_{f_{L/R}}) k^\phi r_c} \]  

(III.2)

Where the \( c_{f_{L,R}} \) are the arbitrary constants from the separation of variables used for each bulk fermion. To simplify the analysis of the theory free parameters literature takes for the light fermions \( c_{f_L} = -c_{f_R} \equiv c_f \), with this convention \( c_f < 1/2 \) fixes the zero mode fermion profile to be TeV localized and \( c_f > 1/2 \) fixes the zero mode fermion profile to be Planck localized.

The original proposals for the embedding of fermions on the WED scenario\(^1\) considered the symmetry \( SU(2)_L \times U(1)_Y \) enlarged to the bulk. The tree level masses for the four dimensional Standard Model fermions are generated under the usual Electroweak Symmetry Breaking mechanism. Five dimensional Yukawa interactions are introduced\(^1\):

\[
L_{\text{Yukawa}} = \lambda_{ij} (\bar{\Psi}_L^i \Psi_R^j H + h.c.) \delta (r_c - \pi r_c), 
\]

(III.3)

where the constants \( \lambda_{ij} \) are dimensionless Yukawa constants for the fermion fields \( i, j \).

\(^1\)Some works addresses the bulk matter space-time back-reaction\(^45\)\(^46\)\(^47\). These references find the deformation on space-time due the presence of matter is very small and consequently the changes on the mass spectra of the kk modes because back-reaction on metric modifications is also small.
From a schematic point of view, if we take the fermions on the diagonalized mass basis, the fermion masses on bulk WED scenario are dependent of both $\lambda_f \equiv \lambda_{ij}$ (where $f$ is the physical fermion) and fermion localization:

$$m_f^2 = \lambda_f v^2 \left( \int d\phi \ f_{\Phi R}(c_{fL}; \phi) f_{\Phi L}(c_{fR}; \phi) \ \delta(\phi r_c - \pi r_c) \right) \quad (\text{III.4})$$

Assuming $c_{fL} = -c_{fR} \equiv c_f$ the Electroweak four dimensional fermion mass is:

$$m_f^2 \sim \lambda_f k \left( c_f - \frac{1}{2} \right) e^{(1-2c_f)\pi kr_c} \quad (\text{III.5})$$

If we take the constant $\lambda_f k \sim 1$ the fermion masses hierarchy is created by fixing the fermion profile localization on the extra dimension, the big hierarchy between fermion masses is exponentially reduced to a hierarchy on the value of the $c_f$ constants. For example, to reproduce the tiny electron mass (assuming all the Yukawa constants $\lambda_f k = 1$) we need to choose $c_e = 0.64$. A complete analysis of the flavour problem on simplest version of WED bulk scenario have been done, for example on [13].

### III.1.1 Electroweak precision tests and the top quark mass

The choice of fermion localization parameters is a more delicate issue on the third family of quarks. The $t_L$ and $b_L$ are zero modes of a bulk $SU(2)_L$ doublet – holding a common profile – there are three geometrical parameters controlling the third family profile localization $c_{3L}$ (for the the left handed doublet) and $c_{tR}$ and $c_{bR}$ (for the right handed singlets). To reproduce both the top and bottom quark masses (with two orders of magnitude distance) we need one of the $L/R$ profiles to be TeV localized.

The SM fermion localizations are chosen in a way to fix the four dimensional couplings to their SM value. By other side the effect of heavy kk modes produces dimension 6 operators, these operators correct SM predictions producing limits on kk-modes mass scale. The pure bosonic operators contribute at tree level to the oblique parameters $S$ and $T$, the fermionic operators contribute non-obliquely correcting processes like $Z \to b_L \bar{b}_L$.

The simplest solution for the third family embedding on the WED bulk (on minimal gauge group setup). Meaning to choose the $t_R$ profile to be TeV localized and the $t_L$ profile to be Planck localized ($c_{tR} \to 1$ and $c_{tL} \to 0$). This scenario leads to very stringent limits on the kk-fermion masses ($O(10)$ TeV), and on the kk-gauge masses ($O(20)$ TeV) [8, 12].

The authors of references [48, 49] realized that enlarge the gauge symmetry on the WED bulk to a $SU(2)_L \times SU(2)_R$ custodial symmetry causes a partial cancellation of the beyond tree level corrections caused by the heavy kk-modes on electroweak corrections from both gauge [48] and fermion sector [49], lowering down the indirect limits on the mass scale of matter kk-modes.

These more realistic scenarios of quark bulk fields embedding makes possible to fix the profiles localization differently from the simplest scenario, specially for the third family of quarks, without hurt Standard Model precision tests.
One should note however that enlarging custodial Symmetry to the WED bulk opens the possibility for the existence new light fermions with exotic charge \([49]\), the phenomenology of such exotic particles was exploited for example by \([50]\). In this note we are ignoring the phenomenology of any BSM particle other than the gravity particles.

In figure (III.1) we draw a schematic view of the profiles of the zero modes matter fields along the extra dimension \(\phi\) in contrast to the kk-graviton.

![Figure III.1: Scheme of matter localization on the different WED scenarios. Left: Bulk scenario. Right: RS1 scenario. The combination of Exponential and Bessel functions makes the kk-Graviton profile to be very TeV localized. To do not overload the figure we do not show the profile of zero mode Graviton, neither of the third generation of fermions.](image)

IV KK-graviton couplings to matter

We derive the perturbative kk-graviton couplings to matter expanding the metric perturbations (equation (II.2)) on the Lagrangian (equation (III.1)). The four dimensional effective couplings are proportional to the integral of the involved field profiles \(d_i\) and suppressed by \(\Lambda_G \equiv \frac{m_G}{x_i k}\).

\[
\mathcal{L} = -\frac{x_i k}{m_G} h^{\mu\nu(1)} \times d_i \frac{f_i^2}{\sqrt{g}}
\]  

(IV.1)

Where \(T_{\mu\nu} \equiv -\frac{2}{\sqrt{g}} \frac{\delta \mathcal{L}}{\delta g^{\mu\nu}}\) is the four dimensional canonical energy-momentum tensor for the field \(i\) \([51]\) and \(d_i \equiv \int d\phi f_{Gr} f_i^2\) is the integral overlap between the profiles of the fields \(i\) with kk-graviton.

Being the kk-Graviton extremely TeV localized the parameter \(d_i\) indicates how much the profile of the field \(i\) is away from TeV brane, in other words how much the coupling of the field with kk-Graviton is volume suppressed.

On bulk scenario all the profile functions are normalized, the parameter \(d_i\) can only be one or less. On RS1 scenario all the \(d_i\)'s parameters are approximately one, therefore
the strength of the couplings between kk-graviton and SM matter are democratic with each field degree of freedom.

**IV.1 Higgs doublet**

To keep the hierarchy of scales derived on section (I.1) the Higgs doublet (higgs singlet plus the Goldstone bosons) is kept on TeV brane, on practical words $d_H = 1$.

The Higgs doublet is confined to TeV brane so the volume suppression parameter is:

$$d_H = f_h(\phi = \pi)$$  \hspace{1cm} (IV.2)

Where $f_h(\phi = \pi) \sim 1$ is the value of kk-graviton profile on TeV brane (see equation [1.5]). To simplify notation we assume there is no volume suppression on its couplings to kk-graviton $f_h(\phi = \pi) \rightarrow 1$. The couplings of the Higgs doublet components with kk-graviton are defined only by its energy momentum tensor:

$$d_H T^H = \partial_\mu H \partial_\nu H - \frac{1}{2} g_{\mu\nu} \left( \partial_\alpha H \partial^\alpha H - m_H^2 H^2 \right)$$ \hspace{1cm} (IV.3)

The decay rate of kk-Graviton to a pair of physical higgses is \[52\]:

$$\Gamma(h_{\mu\nu} \rightarrow hh) = \frac{m^3}{960\pi} \left( \frac{\tilde{k} x_1}{m_G} \right)^2 \left( 1 - \frac{4m_H^2}{m_G^2} \right)^{5/2}$$ \hspace{1cm} (IV.4)

**IV.2 Massless gauge bosons**

As outlined on figure (III.1), the profile of the massless zero mode of spin-1 bulk bosons is fixed to be flat with respect to the kk-graviton profile, suppressing the four dimensional kk-graviton coupling to massless gauge bosons by the factor:

$$d_g = \frac{2}{k\pi r_c} \frac{(1 - J_0(x_1))}{x_1^2 |J_2(x_1)|}$$ \hspace{1cm} (IV.5)

This coupling suppression holds for both gluons, photons and weak bosons on $R_\xi$ gauge. The decay rate of kk-Graviton to a pair of SM non-massive vectors - photon and gluon - are \[52\]:

$$\Gamma(h_{\mu\nu} \rightarrow gg, \gamma\gamma) = 2d_g^2 \times N_{g,\gamma} \times \frac{m_G}{160\pi} \left( \tilde{k} x_1 \right)^2$$ \hspace{1cm} (IV.6a)

$$N_g = 8, N_\gamma = 1$$ \hspace{1cm} (IV.6b)

On figure (IV.1) we draw the dependence of the $d_g$ parameter with the graviton mass $M_{Gr}$, fixing the dimensionless parameter $\tilde{k} = 0.2$. 

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IV.3 Weak bosons

The masses for the weak bosons are generated by the Higgs mechanism: The W and Z masses are proportional to the vacuum expectation value of the Higgs doublet (given on equation I.11). The W and Z longitudinal degrees of freedom results from absorbing the Goldstone bosons. At high energies it is possible to realize analytically the direct relation between the Goldstones bosons and the longitudinal polarization of weak bosons [53, 54], what is known as equivalence theorem.

The physical weak boson fields are described rotating the theory to unitary gauge, the W and Z fields couplings to kk-graviton have two components: a component related with the transverse polarizations (volume suppressed by the $d_g$ parameter, as a massless gauge boson) and the component that arises from ElectroWeak Symmetry Breaking, that is not volume suppressed:

$$d_V T_{\mu\nu}^{VV} = -\left( m_W^2 \frac{\delta(g_{\alpha\beta}W_\alpha W_\beta)}{\delta g_{\mu\nu}} + \frac{m_Z^2}{2} \frac{\delta(g_{\alpha\beta}Z_\alpha Z_\beta)}{\delta g_{\mu\nu}} \right) - d_g T_{\mu\nu}^{YM}(W^\pm, Z, \gamma), \quad (IV.7)$$

where $V$ represents the vector fields on theory and $T_{\mu\nu}^{YM}(W^\pm, Z, \gamma)$ is the canonical Yang Mills energy-momentum tensor for $SU(2)_L \times U(1)_Y$ (massless part) of gauge fields.

The decay rates of kk-Graviton to a pair of physical SM massive vectors - W and Z - are [52]:

$$\Gamma(h_{\mu\nu} \rightarrow ZZ) = \frac{m_G^3}{960\pi} \left( \frac{\tilde{k} x_1}{m_G} \right)^2 (A_Z + d_g B_Z + d_g^2 C_Z) \sqrt{1 - 4 \frac{m_i}{m_G^2}}, \quad (IV.8a)$$

$$\Gamma(h_{\mu\nu} \rightarrow WW) = \frac{m_G^3}{480\pi} \left( \frac{\tilde{k} x_1}{m_G} \right)^2 (A_W + d_g B_W + d_g^2 C_W) \sqrt{1 - 4 \frac{m_i}{m_G^2}}, \quad (IV.8b)$$

$^1$The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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Figure IV.1: Dependence of volume suppression of kk-graviton and bulk massless gauge boson the coupling with the mass of the first kk-Graviton excitation $m_{Gr}$, fixing $\tilde{k} = 0.2$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$. 

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The Yang Mills energy momentum tensor of a gauge field $A_\mu$ and field strength $F_{\mu\nu}^A$, in unitary gauge, is given by $T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F_A, \mu\nu + F_{\alpha\mu}^A F^{A,\beta\mu}$.
where for each $i = W, Z$:

$$A_i = \left(1 + 12 \frac{m_i^2}{m_G^2} + 56 \frac{m_i^2}{m_G^2}\right)$$

(IV.9a)

$$B = 80 \left(1 - \frac{m_i^2}{m_G^2}\right) \frac{m_i^2}{m_G^2}$$

(IV.9b)

$$C = 12 \left(1 - 3 \frac{m_i^2}{m_G^2} + 6 \frac{m_i^4}{m_G^4}\right)$$

(IV.9c)

These expressions holds to bulk scenario, where $d_g$ follows equation (IV.5). But also to RS1 scenario, if we take $d_g = 1$.

On figure (IV.2) we show the split of the kk-graviton decay rate to weak bosons on the different polarizations, resultant of the volume suppression (equation IV.5). On RS1 scenario the kk-graviton decays to preferentially to transverse polarized modes while on bulk scenario it decays preferentially to longitudinally polarized modes. This change of behaviour is a direct consequence of the volume suppression $d_g$.

![Figure IV.2: Separation of polarizations on decay product of $h_{\mu\nu}^{(1)} \rightarrow ZZ$. The same behaviour holds to $h_{\mu\nu}^{(1)} \rightarrow WW$ decays. The continuous red line represents the portion of decay to pure longitudinal gauge boson modes (LL), the dot-dashed blue is the portion of decay to pure transverse modes (TT) and the dashed green the portion of decay to longitudinal-transverse mixed final states (LT).](image)

**IV.4 Fermions**

The kk-graviton couplings to each zero mode mode are weighted by by the overlap of the profiles of the fields.

We separate the coupling of kk-Graviton with the physical fermions in tree parts: two parts devoted to the (massless) Dirac energy momentum tensor of each chirality $T_{\mu\nu}^{\psi_L/R}$, and other part proportional to the EWSB mass:

$$d_\psi T_{\mu\nu}^{\psi} = d_{fL} T_{\mu\nu}^{\psi_L} + d_{fR} T_{\mu\nu}^{\psi_R} + d_{fL,R} g_{\mu\nu} \lambda_f v^2 (\bar{\psi}_R \psi_L + h.c.)$$

(IV.10)

Where $\psi$ is the four dimensional fermion field and $d_{LR} = \int f_{GR} f_{fR} f_{fL}$. For each generation the tree localization parameters are important to define kk-graviton couplings.
On section (III.1) we had noted there are freedom to choose the profiles localization of different fermion chiralities, this freedom is diffused on freedom to fix kk-graviton couplings to bulk fermions.

For bulk leptons and light quarks (first two generations) we need to choose both chiralities to be Planck localized, resulting the constants $d_{f_L} = d_{f_R} = d_{f_{LR}} \sim 0$, meaning we can ignore kk-graviton couplings to bulk light fermions.

On realistic realizations for bulk fermions on WED it is necessary to go beyond minimal embedding of SM bulk fermions. Different realizations of bulk fermion embedding modifies the allowed localization of the fermions extra dimensional profiles, specially for the third family, slightly modifying the kk-graviton couplings to the top and botton quarks. It is not our scope to discuss the model building hypothesis inside the microscopic WED scenario driven by Standard Model Precision Tests.

For practical reasons we consider the case where only the SM gauge symmetry $SU(2)_L \times U(1)_Y$ is enlarged to the bulk\cite{55}, we will refer to this case as consider a elementary top quark. In the elementary top case the decay rate of kk-Graviton to a pair of top quarks with mass $m_T$ is:

$$\Gamma(h_{\mu\nu} \rightarrow tt) = \frac{m_T^3}{80\pi} \left(\bar{k}x_1\right)^2 \left(\frac{3}{4} - \frac{7}{2} \frac{m_T^2}{m_G^2} + 2 \frac{m_T^4}{m_G^4}\right) \sqrt{1 - 4 \frac{m_T^2}{m_G^2}} \quad (IV.11)$$

We exemplify the deviations kk-graviton phenomenology can suffer due model building on fermion sector by comparing the total width and decay rates of the scenario \cite{55} with the extreme hypothesis of consider the profiles of both top chiralities to be Planck localized (≡ ignore the kk-graviton couplings to top quark).

V LHC production and decay

On this section we present calculations of cross sections, total width and branching ratios taking into account the model building reviewed on last section. We show results for benchmark values of cross sections with the parameters $\tilde{k} = 0.2$ and $kr_c = 11$.

For the calculation of RS1 cross section we used the Madgraph5 \cite{56} output of the RS1 model \cite{57} from Feynrules \cite{58} database. We also checked the branching ratios of the model with the analytical expressions presented here. The results for RS1 kk-graviton production cross sections also agrees with the WED implementation default on Pythia6 \cite{59}.

LHC production

There are four production modes for the kk-Graviton on LHC.

1 $d_{fL} \sim d_{f_{LR}} \sim 0$ and $d_{fR} \sim 1$. 

\footnote{1 $d_{fL} \sim d_{f_{LR}} \sim 0$ and $d_{fR} \sim 1$.}
Two of them are inclusive: gluon fusion and quark fusion (Drell Yan); The coupling of kk-Graviton to bulk light fermions is negligible, for bulk scenario we consider only gluon fusion as inclusive production. For completeness we draw the diagrams for kk-graviton inclusive production on figure (V.1). There is no specific generation cut to be added to define inclusive production.

![Feynman diagrams for inclusive KK-graviton production.](image)

**Left:** Gluon fusion.  
**Right:** Drell Yan. To bulk scenario the Drell Yan production

The other two LHC production modes, Weak Boson Fusion (WBF) and weak boson associated production (with the weak boson decaying hadronically), involves the presence of two additional light quarks (that we will call jets) on the hard process. We draw the diagrams of WBF and weak boson associated production on figure V.2.

![Examples of Feynman diagrams for KK-graviton production in association with two jets.](image)

**Left:** Weak boson fusion topology.  
**Right:** Associated production with a weak boson.

The common cuts we use on di-jet system are:

\[ p_{t,j} > 20 \text{ GeV}, \quad \Delta R_{jj} > 0.1, \quad (V.1) \]

where \( p_{t,j} \) is the transverse momenta of each one of the jet (final state light quark) and \( \Delta R_{jj} \) is the angular distance between the hard jets, The relax of \( \Delta R_{jj} \) cut is designed to take into account the possibility of tag the two hard jets as coming from the weak boson on the associated production.

To calculate processes with Madgraph5 simulation we can only specify initial and final state\[.\] When we generate quark initiated process \( pp \rightarrow jj h_{\mu \nu}^{(1)} \) (di-jet associated production) the full process will cover both VBF like and the weak boson associated topologies. We separate the two topologies adding cuts on the final state di-jet system.

\[\text{It is also possible to chose the order of the interaction being generate, and forbid some particular particle to participate the internal process, but this does not help to separate the considered topologies.}\]
The WBF cross section is defined asking for the quark initiated process $pp \to jjh_{\mu\nu}^{(1)}$ and imposing a cut $M_{jj} > 100$ GeV. The weak boson associated production is defined with the same process under the cut $M_{jj} < 100$ GeV.

The production cross sections of kk-bulk graviton for all production modes are calculated we adapting the Feynrules model [57] to match the bulk scenario: we remove all the bulk kk-graviton couplings to all fermions[1] and rescale the its couplings to massless gauge bosons (both photon, gluon and transverse part of weak bosons) with the volume suppression $d_g$ (equation [IV.5]).

We found agreement on the production cross section results between the model mentioned above and the bulk WED scenario implemented by the authors [60] on CalcHep [61]. Following the reference [55] this implementation that assumes the right-handed top TeV localized. On all calculations we use CTEQ6L [62] with dynamical PDF scale, on gluon fusion production mode we use the kk-graviton mass for regularization scale.

On figure (V.3) we show the production cross section of kk-graviton on both RS1 and bulk scenario on the exclusive production, weak boson associated and WBF production.

![Figure V.3: Production cross sections with $k/M_{Pl} = 0.2$, in pb. The red curves correspond to the inclusive production, the green ones to Vector boson fusion and blue to vector boson associated production. The dot-dashed, continuous and dashed curves are respectively to LHC14, LHC13 and LHC8. Right: Bulk scenario. Left: RS1 scenario. The cross section values scale with $k^2$ for all production mechanisms.](image)

For the same compactification scale the kk-graviton production cross section on bulk scenario with respect to RS1 ones for all modes is smaller.

On gluon fusion production mode one suffer with the $d_g$ volume suppression. We also see a decrease of cross section of WBF and associated production processes on bulk scenario with respect to RS1 scenario. On the di-jet associated production modes the counting of degrees of freedom assures that when the particle that couples to the weak boson pair is not hierarchical on polarizations most of the weak gauge modes involved

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1 This model ignores kk-Graviton coupling also to top quarks, this one is not important to define production cross section.
are transverse (see for example [63]). The kk-graviton coupling to bulk weab bosons also experience the $d_g$ volume suppression.

**Decays**

On figure (V.4) we show the kk-graviton branching fractions to SM particles, the left part of this figure shows the kk-graviton branching ratios on bulk scenario, assuming the elementary top scenario. The right part shows kk-graviton branching ratios on RS1 scenario.

On this figure we clearly note the couplings of RS1 kk-graviton with SM particles are democratic with respect to the number of degree of freedom of each field, while the couplings of bulk kk-graviton with SM particles are predominantly to the fields localized on TeV brane (massive bosons and higgs field).

The highest branching fraction of a RS1 kk-graviton is to di-jet final states ($gg + qq$). The di-photon channel would also be a good search channel on this scenario, by its cleanness and non negligible branching ratio.

For the bulk kk-graviton case the highest branching fractions are to massive states, heavy di-bosons and top quarks (on the elementary top scenario). On the case of a trans-TeV resonance, each one of the di-bosons products are boosted and consequently its sub-products collimated. The development of substructure techniques make from the hadronic di-boson channels golden channels on bulk kk-graviton search.

Figure (V.4): KK-graviton branching fractions. **Left:** Bulk scenario, assuming elementary top quark. **Right** RS1 scenario. The dashed line represents the individual RS1 kk-graviton decay rates to a pair of light fermions $f\bar{f}$, where $f$ represents both light quarks ($u,d,s,c,b$) or leptons ($e,\mu,\tau$). – Branching ratios are independent of $k$ parameter.

Figure (V.5) shows a zoom of kk-graviton branching ratios comparing two distinct scenarios for third family WED embedding: The thick curves considers a fully elementary top quark [60] while thin dashed ones we ignores kk-graviton coupling with top quark. It is not a surprise to note that kk-graviton branching ratios to bosons pais ($W/Z/H$) increases when there is not kk-graviton coupling to top quark.

On figure (V.6) we see the total width on RS1 scenario the kk-graviton is around two orders of magnitude bigger than on bulk scenario, given the same geometric parameters. The experimental resolution on the width of the resonance depends on the investigated
Figure V.5: Zoom of the dominant branching ratios for bulk KK-graviton comparing two hypothesis of fermion WED embedding. The thick curves considers a fully elementary top quark ([60]) and thin dashed ones completely ignores kk-graviton coupling with top quark.

channels and also of the mass hypothesis probed. The experimental searches are usually designed using signal models with negligible width (with respect to experimental resolution).

Figure V.6: Total width for kk-graviton with $k/M_{Pl} = 0.2$. The green curve represents RS1 scenario and the red curves represents bulk scenario. We show two different hypothesis of treating kk-graviton couplings to bulk third family of quarks: the continuous curve considers a fully elementary top quark ([60]) while dashed curve ignores kk-graviton coupling with top quark. Total width values scale with $\tilde{k}^2$ parameter.
VI Remarks on phenomenology

The weak boson polarization dominance of kk-graviton coupling (seen on figure IV.2) can be experimentally probed when weak boson are present on both final or/and kk-graviton production modes.

- **On di-boson decays of a heavy (\(\sim \text{TeV}\)) kk-graviton:** it is experimentally possible to distinguish the polarization of a boostet hadronic weak boson by looking at disposition of jet sub-cores [64, 65].

- **On WBF kk-graviton production:** It was already noted that the behaviour of WBF jets reminiscent of production of a resonance is dependent of the spin of such particle [66]. We note however that the topology of WBF jets is also sensible on the underlying model that defines the momentum dependence of couplings.

Concerning the last item we show on figure IV.2 the parton level pseudo-rapidity between WBF jets comparing processes of kk-graviton WBF production on WED bulk and RS1 scenarios with the WBF SM higgs production. The jet system coming from the kk-graviton production tend to be more central than on SM case.

We also note a stronger mass dependence of WBF topology of processes involving bulk scenario kk-gravitons. This is due its couplings to be mainly with longitudinal gauge bosons, making the dominant part of coupling to be more momentum dependent than RS1 case. On this last one the kk-graviton coupling to weak bosons is mainly to transverse modes – the coupling is mainly not dependent of momenta, resulting on no mass dependence on the WBF topology.

More detailed studies about WBF topology in kk-graviton production are beyond the scope of this note. Dedicated studies exploiting the WBF resonance production on HL-LHC are now ion progress [67, 68].
Predictive scenarios that motivates the search for high mass resonances with spin even (0 or 2) on LHC are based on the existence of extra dimensions. We had reviewed the physics of a heavy spin-2 particle interpreting it as kk-Graviton under Warped Extra Dimension scenario. Two hypothesis was exploited: RS1 (where only gravity is allowed to propagate on the extra dimensional bulk) and a bulk scenario where SM matter is also allowed to propagate on the extra dimensional direction, using the same gauge construction of SM. We also commented on some points that allow/urge for further model building on the exploited scenarios, like the bulk fermions embedding.

We had re-calculated its cross sections, total width and branching ratios to put side by side the results for RS1 and bulk scenarios and allow a fair comparison between the hypothesis. On the course of calculation we had cross checked with analytical calculation some of the leading order kk-graviton Monte Carlo implementations available for the community, implemented on Madgraph5, Calchep and Pythia.

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