SEARCHES FOR $W'$ AND $Z'$ IN MODELS WITH LARGE EXTRA DIMENSIONS

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We discuss the characteristic features of processes mediated by intermediate gauge bosons in the framework of theories with large extra dimensions and show that if gauge bosons propagate in the entire multidimensional space, then a destructive interference arises not only between $W$ and $W'$ (or $Z$ and $Z'$) but also between $W'$ and $Z'$ and the respective Kaluza–Klein towers of higher excitations of $W$ and $Z$ bosons. We perform and graphically present calculations for the LHC with the center-of-mass energy 14 TeV.

Keywords: extra dimension, brane, Kaluza–Klein mode

1. Introduction

In recent years, models with “universal extra dimensions” have been widely discussed in the literature [1]–[11]. In this case, the standard model (SM) fields except the Higgs field can propagate in the whole multidimensional space–time. This leads to some interesting phenomenological predictions, which we discuss in this paper.

The characteristic feature of theories with compact extra dimensions is the presence of towers of Kaluza–Klein (KK) excitations of the bulk fields with all the excitations of a bulk field having the same type of coupling to the SM fields. We assume that for a bosonic bulk field $\phi$ (or a set of fields) of arbitrary tensor type in a $(4+d)$-dimensional space–time, the relevant part of the action has the form

$$ S = \int \sqrt{-\gamma} d^{4+d}x \, L(\phi) + \int_{\text{brane}} d^4 x \left( L_{(\text{SM}-\phi)} + g M^{-d/2} J_{\text{SM}}^* \phi \right), $$

where $\gamma_{MN}$ ($M, N = 0, 1, 2, 3, \ldots, 3 + d$, $\text{sgn} \gamma = +, -, \ldots, -$) denotes the background metric in the bulk, $L(\phi)$ is the bulk Lagrangian of the field $\phi$, the Lagrangian of the SM fields, which do not propagate in the bulk, is denoted by $L_{(\text{SM}-\phi)}$, the interaction term $J_{\text{SM}}^* \phi$ is the scalar product of the corresponding current of the SM fields $J_{\text{SM}}$ and the field $\phi$ on the brane, $g$ is a four-dimensional (in general, dimensional) coupling constant, and $M$ is the fundamental energy scale of the $(4+d)$-dimensional theory defined by the gravitational interaction and is assumed to be in the TeV energy range.

It is well known that the bulk field $\phi(x, y)$, $x = \{x^i\}, y = \{x^i\}, i = 4, \ldots, 3 + d$, can be expanded in KK modes with definite masses $\phi^{(n)}(x)$ and their wave functions in the space of extra dimension $\psi^{(n)}(y)$ as

$$ \phi(x, y) = \sum_n \psi^{(n)}(y) \phi^{(n)}(x), \quad n = (n_1, \ldots, n_d). $$

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The current $J_{\text{SM}}$ and the coupling constant $g$ are completely defined by the interaction of the zero mode $\phi^{(0)}(x)$, which is a field of the SM or the graviton field, with the fields of the four-dimensional SM according to

$$gJ_{\text{SM}} = \frac{\delta I_{\text{int}}^{\text{SM}}}{\delta \phi^{(0)}}$$

It is easy to show (see a detailed derivation in [12]) that if we consider this theory for the energy or momentum transfer much smaller than the masses of the KK excitations $\phi^{(n)}$, $n \neq 0$, we can pass to the effective “low-energy” theory, which can be obtained by the standard procedure. Namely, we must drop the momentum dependence in the propagators of the heavy modes and integrate them out in the functional effective “low-energy” theory, which can be obtained by the standard procedure. We thus obtain a contact interaction of the SM fields

$$S = \int d^4x \frac{1}{2} \partial_\mu \phi^{(0)} \ast \partial^\mu \phi^{(0)} - \frac{1}{2} M_0^2 \phi^{(0)} \ast \phi^{(0)} + L_{\text{int}}^{(0)} +$$

$$+ gM^{-d/2} \psi^{(0)}(y_b)J_{\text{SM}} \ast \phi^{(0)} + L_{(\text{SM} - \phi)} +$$

$$+ \frac{1}{2} g^2 M^{-d} \left( \sum_{n \neq 0} \frac{\langle \psi^{(n)}(y_b) \rangle^2}{M_n^2} \right) J_{\text{SM}} \ast \Delta \ast J_{\text{SM}},$$

where $M_n$ is the mass of the $n$th mode and $\Delta$ is the tensor structure (the numerator) of the propagator with the momentum equal to zero, which is the same for all modes, and $\{y_b\}$ denotes the coordinates of the brane in the space of extra dimensions. We thus obtain a contact interaction of the SM fields

$$\lambda J_{\text{SM}} \ast \Delta \ast J_{\text{SM}}, \quad \lambda = \frac{1}{2} g^2 M^{-d} \left( \sum_{n \neq 0} \frac{\langle \psi^{(n)}(y_b) \rangle^2}{M_n^2} \right),$$

and the sum of all the other terms in (4) is the Lagrangian of the SM $L_{\text{SM}}$. We see that the Lagrangian structure is fixed by the corresponding structure of the SM currents $J_{\text{SM}}$ and the spin-density matrix $\Delta$, which is defined by the type of the field $\phi$ as shown in formula (4).

## 2. Effective Lagrangian for the gauge interaction

In this section, we discuss the case of the contact interactions due to the $SU(2) \times U(1)$ gauge fields in the bulk. These fields are described in the bulk by vector potentials $W_M$ and $B_M$, which yield four-dimensional vector and scalar fields. The latter are in the trivial and the adjoint representations of $SU(2)$ and cannot break $SU(2) \times U(1)$ to $U(1)_{\text{em}}$, as is necessary in the SM. We therefore assume that the gauge symmetry is broken standardly by the Higgs field on the brane. It is useful to introduce the charged vector fields

$$W_{\mu}^\pm = \frac{W_{\mu}^1 \mp W_{\mu}^2}{\sqrt{2}}$$

and the standard mixing of the neutral vector fields

$$Z_{\mu} = W_{\mu}^3 \cos \theta_W - B_{\mu} \sin \theta_W,$$

$$A_{\mu} = W_{\mu}^3 \sin \theta_W + B_{\mu} \cos \theta_W.$$

After the spontaneous symmetry breaking, the neutral component of the brane Higgs field acquires a vacuum value $v/\sqrt{2}$, and there arises a quadratic interaction of the vector fields of the form

$$\frac{g^2 v^2}{4} M^{-d} \sum_{m,n} \psi_m(y_b) \psi_n(y_b) \eta^\mu \eta^\nu W_{\mu}^{(m)} W_{\nu}^{(n)} +$$

$$+ \frac{(g^2 + g'^2) v^2}{8} M^{-d} \sum_{m,n} \psi_m(y_b) \psi_n(y_b) \eta^\mu \eta^\nu Z_{\mu}^{(m)} Z_{\nu}^{(n)},$$

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where \( \psi_n(y_b) \) denotes the wave functions of the KK modes of the fields \( W^\pm_\mu \) and \( Z_\mu \) on the brane. Because of this interaction, the KK modes are no longer the mass eigenstates, which are now superpositions of the modes [13]. But if the mass scale \( g v \) generated by the Higgs field is much smaller than the mass of the first KK excitation (exactly the scenario that we study here), then this mixing of KK modes can be neglected [13]. The coupling of the KK modes to the fields of the SM is defined by that of the zero mode and has the form

\[
L_{\text{int}} = \frac{g}{\sqrt{2}} M^{-d/2} \sum_{n>0} \psi_n(y_b) (J^{+\mu}_\mu W^{(n)-}_\mu + J^{-\mu}_\mu W^{(n)+}_\mu) + \frac{g}{\cos \theta_W} M^{-d/2} \sum_{n>0} \psi_n(y_b) J^{\mu}_\mu A^{(n)}_{\mu},
\]

where \( J^\pm_\mu \) and \( J^{(0)}_\mu \) are the weak charged and neutral currents of the SM particles and \( J^{\mu}_\mu \) is the electromagnetic current of the SM particles. Integrating over the heavy modes, we again obtain the effective Lagrangian of form (5). Recalling that all the masses are proportional to \( M \) and the wave functions are proportional to \( M^{d/2} \), we then obtain the effective Lagrangian for the interaction of the SM fields due to the excitations of the \( SU(2) \times U(1) \) gauge bosons

\[
L_{\text{eff}} = G_F M^0_W \left( C_W J^{+\mu}_\mu J^{(0)}_\mu + C_W J^{-\mu}_\mu J^{(0)}_\mu + C_Z J^{(0)}_\mu J^{(0)}_\mu + C_A J^{\mu}_\mu A^{(0)}_{\mu} \right),
\]

where \( G_F \) denotes the Fermi constant. The constants \( C_W, C_Z, \) and \( C_A \) again depend on the model and can be estimated only in a specific model. In particular, in the simplest model with two branes and one flat extra dimension, the constants can be estimated as

\[
C_W = \frac{\pi^2}{6\sqrt{2}}, \quad C_Z = \frac{\sqrt{2} \pi^2}{6 \cos^2 \theta_W}, \quad C_A = \frac{2 \sqrt{2} \pi^2 \sin^2 \theta_W}{3}.
\]

We now estimate the constants in the effective Lagrangian for the gauge interaction in the case of the Randall–Sundrum bulk [14]. First, because the bulk is five-dimensional, we can pass to the axial gauge, where the components corresponding to the extra dimension are equal to zero [15]. Hence, there are no corresponding scalar fields in the effective four-dimensional theory. The wave functions \( w_n(y) \) of the fields \( A^{(n)}_\mu(x) \) with definite masses are solutions of a Sturm–Liouville eigenvalue problem with Neumann boundary conditions. The wave function of the massless zero mode, unlike the wave function for the tensor zero mode, is therefore constant in the extra dimension, which guarantees the universality of its coupling constant [16]. The wave functions of the excitations on the brane behave as \( w_n(y)\big|_{y=y_b} \sim \sqrt{k} \), i.e., similarly to the wave functions of the tensor modes. The masses of the modes also appear to be in the TeV energy range [15]. We are interested in the cases where the masses of the modes and the mass gaps between the modes are quite large, for example, of the order of a few TeV.

Below, we consider some processes with the KK electroweak gauge bosons at energies accessible at the LHC. We note that the coupling constants and the masses of the modes depend significantly on the particular model under consideration. We also extract the first KK mode from effective Lagrangian (11) and assume that the accessible energy is above the production threshold of the first KK mode; these modes are correspondingly called \( W' \) and \( Z' \). We take all the other modes into account by using the contact effective interactions.

The symbolic and numerical computations, including simulations of the SM background for the LHC, were performed using the CompHEP software package [17]. The corresponding Feynman rules were implemented in the new version of CompHEP.
3. Processes with KK gauge bosons

It was shown in [18] independently of the model that there exists a nontrivial destructive interference between the processes mediated by $W$ and $W'$. If we assume that the gauge bosons propagate in the bulk, then the $W$ boson is just the zero KK mode, the $W'$ boson is the first excitation, and there exists an infinite tower of KK modes above it. The same of course holds for $Z$ and $\gamma$. In this case, we expect that the higher KK modes can also interfere with the zero and the first modes.

We now turn to specific examples. As noted in the preceding section, the coupling constants and the masses of the modes depend significantly on the particular model. For simplicity, we assume that all the KK modes have the same coupling constant as those of the corresponding SM $W$ and $Z$ bosons and photon. The masses of the $W'$ and $Z'$ bosons and of the first KK excitation of the photon are respectively $M_{W'}$, $M_{Z'}$, and $M_{\gamma'}$. The remaining towers of the modes were simulated in CompHEP using auxiliary particles with the masses $M_{W',\text{sum}}$, $M_{Z',\text{sum}}$, and $M_{\gamma',\text{sum}}$ without the four-momenta in the propagators. Indeed, we can schematically write the squared amplitude as

$$\left|\frac{1}{p^2 - M^2} + \frac{1}{p^2 - M'^2} - \sum_{n=2}^{\infty} \frac{1}{M^2_n}\right|^2 = \left|\frac{1}{p^2 - M^2} + \frac{1}{p^2 - M'^2} - \frac{1}{M^2_{\text{sum}}}\right|^2,$$

where $M_n$ corresponds to the masses of the KK modes. This formula shows the origin of the parameters used below. The term $1/M^2_{\text{sum}}$ corresponds to effective contact interaction (11).

We now consider specific processes involving KK gauge bosons. For illustrative purposes, all calculations were performed for the LHC with the center-of-mass energy 14 TeV.

We first consider a process with a $W'$ boson plus the remaining tower of modes, namely, the single top production. We assume that the mass of the first mode is $M_{W'} = 2\,\text{TeV}$ and the effective mass is $M_{W',\text{sum}} = 2.8\,\text{TeV}$. The width of the $W'$ resonance was calculated as $\Gamma_{W'} = 65.7\,\text{GeV}$. The distributions for the process $u\bar{d} \to t\bar{b}$ give the main contribution to the process $pp \to t\bar{b}$ at the LHC presented in Figs. 1 and 2.

We calculated for only the SM $W$ boson, for the SM $W$ boson plus the $W'$ boson, and for the SM $W$ boson plus the $W'$ boson plus the remaining tower of KK modes. It is clear from Fig. 1 that the presence of the $W'$ boson leads to a destructive interference at energies smaller than the mass of the $W'$ resonance.
The $P_T$ distribution for the single top production at the LHC: SM is the dash-dotted line, SM + $W'$ is the solid line, and SM + $W' + KK$ is the dotted line.

Fig. 2. The $P_T$ distribution for the single top production at the LHC: SM is the dash-dotted line, SM + $W'$ is the solid line, and SM + $W' + KK$ is the dotted line.

The invariant-mass distribution for the Drell–Yan process at the LHC: $\gamma + Z$ is the dash-dotted line, $\gamma + Z + Z' + \gamma'$ is the solid line, and $\gamma + Z + Z' + \gamma' + KK$ is the dotted line.

Fig. 3. Invariant-mass distribution for the Drell–Yan process at the LHC: $\gamma + Z$ is the dash-dotted line, $\gamma + Z + Z' + \gamma'$ is the solid line, and $\gamma + Z + Z' + \gamma' + KK$ is the dotted line.

At energies larger than the mass of the $W'$ resonance, we see that there is an increase in the distribution tails due to the existence of $W'$ and the corresponding KK modes in comparison with the case of only the SM $W$ (see Fig. 1). We note that KK modes above $W'$ can lead to a quite considerable modification of the distributions (see Fig. 2).

The second process that we consider is with the $Z'$ boson and the $\gamma'$ boson plus the remaining towers of the modes, namely, the Drell–Yan process with $u$ quarks, which is also dominant at the LHC. We assume that the masses of the first modes are $M_{Z'} = 2.3$ TeV and $M_{\gamma'} = 2$ TeV and the effective masses are $M_{Z',\text{num}} = 3.2$ TeV and $M_{\gamma',\text{num}} = 2.8$ GeV. The widths of the $Z'$ and $\gamma'$ resonances were found to be $\Gamma_{Z'} = 0.026$ TeV and $\Gamma_{\gamma'} = 0.021$ TeV. The corresponding distributions are presented in Figs. 3 and 4. It can be seen that the properties of these distributions are analogous to those of the single top production.

There is good reason to believe that the NLO corrections do not destroy this interference picture. First,
it is obvious that the corrections to the external lines do not violate the structure of amplitude (12). Of course, the most dangerous terms seem to be those with the self-energy of the gauge bosons. But these self-energy terms are defined such that they vanish on the mass shell and contribute only to the renormalization of the particle widths and masses.

Therefore, our analysis has shown that the KK modes must be taken into account because they can contribute to the amplitudes of the corresponding processes. Of course, single particles $W''$ or $Z''$ can in principle also provide analogous effects, but simultaneous effects with KK gravitons and $W'$ and $Z'$ KK modes can be interpreted in favor of the existence of extra dimensions.

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REFERENCES

1. I. Antoniadis, Phys. Lett. B, 246, 377–384 (1990).
2. T. Gherghetta and A. Pomarol, Nucl. Phys. B, 586, 141–162 (2000); arXiv:hep-ph/0003129v2 (2000).
3. T. Appelquist, H.-C. Cheng, and B. A. Dobrescu, Phys. Rev. D, 64, 035002 (2001); arXiv:hep-ph/0012100v2 (2000).
4. T. G. Rizzo, Phys. Rev. D, 64, 095010 (2001); arXiv:hep-ph/0106336v1 (2001).
5. C. Macesanu, C. D. McMullen, and S. Nandi, Phys. Lett. B, 546, 253–260 (2002); arXiv:hep-ph/0207269v2 (2002).
6. K. Agashe, A. Belyaev, T. Krupovnickas, G. Perez, and J. Virzi, Phys. Rev. D, 77, 015003 (2008); arXiv:hep-ph/0612015v1 (2006).
7. K. Agashe, G. Perez, and A. Soni, Phys. Rev. D, 75, 015002 (2007); arXiv:hep-ph/0606293v2 (2006).
8. K. Agashe, H. Davoudiasl, G. Perez, and A. Soni, Phys. Rev. D, 76, 036006 (2007); arXiv:hep-ph/0701186v3 (2007).
9. A. L. Fitzpatrick, J. Kaplan, L. Randall, and L. T. Wang, JHEP, 0709, 013 (2007); arXiv:hep-ph/0701150v2 (2007).
10. B. Lillie, L. Randall, and L. T. Wang, JHEP, 0709, 074 (2007); arXiv:hep-ph/0701166v1 (2007).
11. G. Burdman, L. Da Rold, O. Eboli, and R. Matheus, Phys. Rev. D, 79, 075026 (2009); arXiv:0812.0368v1 [hep-ph] (2008).
12. E. E. Boos, V. E. Bunichev, M. N. Smolyakov, and I. P. Volobuev, Phys. Rev. D, 79, 104013 (2009); arXiv:0710.3100v5 [hep-ph] (2007).
13. A. Mück, A. Pilaftsis, and R. Rückl, “An introduction to 5-dimensional extensions of the standard model,” in: Heavy Quark Physics (Lect. Notes Phys., Vol. 647, D. Blaschke, M. A. Ivanov, and T. Mannel, eds.), Springer, Berlin (2004), p. 189–211; arXiv:hep-ph/0209371v2 (2002).
14. L. Randall and R. Sundrum, Phys. Rev. Lett., 83, 3370–3373 (1999); arXiv:hep-ph/9905221v1 (1999).
15. H. Davoudiasl, J. L. Hewett, and T. G. Rizzo, Phys. Lett. B, 473, 43–49 (2000); arXiv:hep-ph/9911262v2 (1999).
16. V. A. Rubakov, Phys. Usp., 44, 871–893 (2001).
17. E. Boos, V. Bunichev, M. Dubinin, L. Dudko, V. Edneral, V. Ilyin, A. Kryukov, V. Savrin, A. Semenov, and A. Sherstnev, Nucl. Instrum. Meth. Phys. Res. A, 534, 250–259 (2004); E. Boos et al., PoS A, CAT08, 008 (2008).
18. E. Boos, V. Bunichev, L. Dudko, and M.Perfilov, Phys. Lett. B, 655, 245–250 (2007); arXiv:hep-ph/0610080v3 (2006).