The specificity of the interactions of electroweak gauge bosons coming from extra dimensions

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Abstract

We discuss the specificity of the interactions of the electroweak gauge boson excitations in models with warped extra dimensions. In particular, we show that the couplings of the gauge boson excitations $W'$, $Z'$, and $\gamma'$ to the SM gauge bosons treated as the zero modes of the 5D gauge fields are either exactly equal to zero or very much suppressed. In the former case, the three-particle and four-particle interaction Lagrangians of the SM gauge bosons and their lowest excitations are found explicitly. Meanwhile, the couplings of $W'$, $Z'$, and $\gamma'$ to the SM fermions are non-zero allowing for their production and decays. These are the characteristic features of the gauge boson excitations in models with warped extra dimensions, which distinguish them from the gauge boson excitations in other models beyond the SM.

1 Introduction

Nowadays the Standard Model (SM) is capable of describing very well a great amount of experimental facts and results. However, there is a number of serious problems such as the hierarchy problem, dark matter, and CP violation, which cannot be consistently explained in its framework. To explain them, a large number of various models beyond the SM and scenarios of new physics have been put forward. Almost all of the SM extensions predict the existence of new particles, in particular, the existence of massive charged and neutral vector particles besides the gauge bosons of the SM. These extra vector bosons appear either because of an extension of the SM gauge group (see e.g., [1–8]), or as excitations of the SM gauge bosons (see e.g., [9–13]). The lowest excitations of $W$, $Z$, and $\gamma$ are usually denoted by $W'$, $Z'$, and $\gamma'$ or para-photon. Depending on a particular model, the physical properties and interactions of these extra vector bosons may be rather different.

If extra vector bosons are found at the LHC, there arises the problem of specifying the theory beyond the SM, to which they may belong. To solve this problem one has to study the characteristic features of the extra vector bosons in different models and the specificity of their production channels and decay modes.

In the present paper we pay closer attention to the interactions of the excitations of the electroweak SM gauge bosons in various models with extra dimensions. Such models have been widely discussed in the literature in the brane-world set-up either in the context of the flat bulk (UED models) [14,17], or in the Randall-Sundrum bulk [18,30]. In fact, in the brane world
models there appear excitations of two neutral gauge bosons $Z'$, and $\gamma'$, while in many BSM models based on extensions of the SM gauge group only additional $Z'$ bosons appear. This property is already a characteristic feature of extra dimensional models. However the feature might not be very pronounced since a typical mass splitting between $Z'$ and $\gamma'$, as it will be discussed below, is expected to be very small of the order of $m_{Z}^{2}/(2m_{\gamma})$. Such small mass splitting might be very difficult to resolve experimentally.

It is a common knowledge that, in models with flat universal extra dimensions, there exists the so-called Kaluza-Klein number conservation, which is a trivial consequence of the properties of the Fourier transform on the circle reflecting the multidimensional energy-momentum conservation law. It means that in such models there is no single production at tree level of the Kaluza-Klein excitations, and a Kaluza-Klein excitation cannot decay at tree level into the SM particles. Nevertheless, such processes can take place due to loop corrections [16], which are usually very much suppressed. This well-known property leads to rather specific collider signatures with cascade decays and stable lightest state similar to SUSY signatures but with different spin correlations in the decay chains [31].

The brane-world models with the Randall-Sundrum bulk are rather different from the UED models with the flat bulk [14–17], because the fields of different tensor type have different wave function profiles. For this reason, this scenario does not necessarily lead to the KK number conservation. In this case the production of single KK states and their decays are possible if kinematically allowed.

The extra vector bosons usually have interactions similar to those of the SM gauge bosons and can mediate the same processes with SM particles. In this case non-trivial interference between the contributions, for example, of $W$ and $W'$, to various processes [32–37] could influence the experimental observation of the latter and the exclusion limits for their masses [38–41]. Interference of $W'$, $Z'$ and $\gamma'$ with the SM bosons and with the rest of corresponding KK-towers coming from extra dimensions leads to certain changes in invariant mass and $P_t$ distributions [42, 43]. However this specific feature of the EW gauge boson KK excitations is also rather delicate for experimental detection.

In the present paper we show that the interaction properties of EW gauge boson excitations in models with warped extra dimensions are essentially different from the decay properties of these excitations in other models. The interaction properties of the excitations of the neutral EW gauge bosons and gluons in the unstabilized Randall-Sundrum model [18] have been already touched upon in papers [22, 23], the greater emphasis having been put on the properties of the gluon excitations. Here we will study the interaction properties of EW gauge boson excitations in more detail taking into consideration the excitations of the charged SM gauge bosons. It will be demonstrated explicitly that a simple common property of the EW gauge bosons excitations $W'$, $Z'$, and $\gamma'$ in brane-world models is that their couplings to the SM gauge fields treated as the zero modes of the 5D gauge fields are either exactly equal to zero or very much suppressed. At the same time the couplings of $W'$, $Z'$, and $\gamma'$ to the SM fermions are non-zero allowing for their production and decays.
2 Electroweak gauge fields in the bulk

Without loss of generality we consider a brane-world RS model stabilized by the bulk scalar field \([44,45]\). Such stabilized brane-world models can solve the hierarchy problem of the gravitational interaction and give rise to the masses of KK excitations in the TeV energy range \([46,47]\), which is smaller than the limit of the order of 20 TeV following from the EW precision data \([20]\) for the unstabilized case. It is worth to mention that in the considered brane-world RS models there are no flavour changing neutral currents strongly suppressed by the present-day experimental data, because the neutral currents have the same diagonal structure as in the SM.

Let us consider the electroweak gauge fields in a five-dimensional space-time with the coordinates
\[
x^M = \{x^\mu, y\}, \quad M = 0, 1, 2, 3, 4.
\]
The compact extra dimension forms the orbifold \(S^1/Z_2\),\[\[\text{which can be represented as the circle of circumference } 2L \text{ with the coordinate } -L \leq y \leq L \text{ and the points } -y \text{ and } y \text{ identified.\[}\]

The background metric is assumed to have the standard form, which is often used in brane world models:
\[
ds^2 = \gamma_{MN}(y) dx^M dx^N = e^{2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu - dy^2.
\]
(1)

This metric gives rise to a usual brane-world model, i.e. it is a solution to the equations of motion for five-dimensional gravity, two branes with tension and the stabilizing bulk scalar field. The explicit form of the solution for \(\sigma(y)\) is unimportant for our considerations and we do not specify it.

We consider the following standard action of the \(SU(2) \times U(1)\) gauge fields in this background:
\[
S = \frac{1}{2L} \int d^4xdy \sqrt{\gamma} \left( -\frac{1}{4} W^{i, MN} W_{i, MN} - \frac{1}{4} B^{MN} B_{MN} \right),
\]
(2)

where \(\gamma = \det \gamma_{MN}\), and the factor \(1/2L\) in front of the integral is introduced for convenience and chosen so that the dimensions of the bulk gauge fields and coupling constants be the same as in the four-dimensional theory, the field strength tensors are given by
\[
W_{i, MN} = \partial_M W^{i}_N - \partial_N W^{i}_M + g \epsilon^{kl} W^{k}_{i, M} W^{l}_{i, N}, \quad B_{MN} = \partial_M B_N - \partial_N B_M,
\]
(3)

and the fields satisfy the orbifold symmetry conditions
\[
W^{i}_{i,}(x,-y) = W^{i}_{i,}(x,y), \quad W^{i}_{4,}(x,-y) = -W^{i}_{4,}(x,y),
\]
\[
B_{i,}(x,-y) = B_{i,}(x,y), \quad B_{4,}(x,-y) = -B_{4,}(x,y).
\]

Next we will study the excitations of the gauge bosons. To this end we pass to the axial gauge, where the components \(W_i^4, B_4\) of the vector fields are equal to zero, and consider only the four-vector components of the five-dimensional gauge fields, whose zero modes must play the role of the SM gauge bosons. From action (2) it is easy to get the following action for the four-vector components of the five-dimensional gauge fields:
\[
S = \frac{1}{2L} \int d^4xdy \left( -\frac{1}{4} \eta^{\mu\nu} \eta^{\alpha\beta} W_{i, \mu\alpha} W^{i, \nu\beta} + e^{2\sigma(y)} \frac{1}{2} \eta^{\mu\nu} \partial_\mu W^{i, \nu\alpha} \partial_\alpha W^{i} + \frac{1}{4} \eta^{\mu\nu} \eta^{\alpha\beta} B_{\mu\alpha} B_{\nu\beta} + e^{2\sigma(y)} \frac{1}{2} \eta^{\mu\nu} \partial_\mu B_{\nu\alpha} \partial_\alpha B_{\nu} \right).
\]
(5)
Now, making the standard redefinition of the gauge fields

\[ Z_{\mu} = \frac{1}{\sqrt{g^2 + g'^2}} (gW_{\mu}^3 - g'B_{\mu}), \quad A_{\mu} = \frac{1}{\sqrt{g^2 + g'^2}} (gB_{\mu} + g'W_{\mu}^3), \quad W_{\mu}^\pm = \frac{1}{\sqrt{2}} (W_{\mu}^1 \mp iW_{\mu}^2), \]

where \( g \) and \( g' \) are the coupling constants of the gauge groups \( SU(2) \) and \( U(1) \) respectively, we pass to the physical degrees of freedom of the theory.

Let us rewrite action (5) in terms of these redefined fields and decompose it into the quadratic part and the interaction Lagragian. It takes the form

\[
S = \frac{1}{2L} \int d^4x dy \left( -\frac{1}{2} \eta^{\mu\nu} \eta^{\alpha\beta} W_{\mu\alpha} W_{\nu\beta} - \frac{1}{4} \eta^{\mu\nu} \eta^{\alpha\beta} A_{\mu\alpha} A_{\nu\beta} - \frac{1}{4} \eta^{\mu\nu} \eta^{\alpha\beta} Z_{\mu\alpha} Z_{\nu\beta} + e^{2\sigma} \eta^{\mu\nu} \partial_4 W_{\mu}^+ \partial_4 W_{\nu}^- + e^{2\sigma} \frac{1}{2} \eta^{\mu\nu} \partial_4 A_{\mu} \partial_4 A_{\nu} + e^{2\sigma} \frac{1}{2} \eta^{\mu\nu} \partial_4 Z_{\mu} \partial_4 Z_{\nu} + L_{WWV} + L_{WWVV} \right),
\]

where \( W_{\mu\nu}^\pm = \partial_\mu W_{\nu}^\pm - \partial_\nu W_{\mu}^\pm, \quad A_{\mu\nu} = \partial_\mu A_{\nu} - \partial_\nu A_{\mu}, \quad Z_{\mu\nu} = \partial_\mu Z_{\nu} - \partial_\nu Z_{\mu} \), \( L_{WWV} \) and \( L_{WWVV} \) are the gauge boson three-particle and four-particle self interaction 5D Lagrangians respectively. The three-particle self interaction Lagrangian is explicitly given by

\[
L_{WWV} = -i g \left[ (W_{\mu}^+ W_{\mu}^- - W_{\mu}^+ W_{\mu}^-) (A_{\nu} \sin \theta_W + Z_{\nu} \cos \theta_W) \right. + \left. W_{\mu}^+ W_{\mu}^- (A_{\nu} \sin \theta_W + Z_{\nu} \cos \theta_W) \right],
\]

and the four-particle self interaction Lagrangian looks like

\[
L_{WWVV} = -\frac{g^2}{2} \left[ (W_{\mu}^+ W_{\mu}^-)^2 - (W_{\mu}^+ W_{\mu}^-) (W_{\mu}^+ W_{\mu}^-) \right. + \left. 2 (W_{\mu}^+ W_{\mu}^-) (A_{\nu} \sin \theta_W + Z_{\nu} \cos \theta_W) (A_{\nu} \sin \theta_W + Z_{\nu} \cos \theta_W) \right. \right. + \left. \left. 2 W_{\mu}^+ (A_{\nu} \sin \theta_W + Z_{\nu} \cos \theta_W) W_{\mu}^- (A_{\nu} \sin \theta_W + Z_{\nu} \cos \theta_W) \right]\right].
\]

In what follows, we assume that the bulk gauge symmetry \( SU(2) \times U(1) \) is spontaneously broken by the bulk or brane localized Higgs field. Here we will not go into details of this mechanism \[15\]. We will just suppose that, as a result, mass terms for the fields \( W_{\mu} \) and \( Z_{\mu} \) are generated so that the masses of their zero modes are given by the standard expressions

\[
m_W = \frac{g v}{2}, \quad m_Z = \frac{\sqrt{g^2 + g'^2}}{2} v,
\]

\( v \) denoting the standard vacuum value of the Higgs field.

The equations for the wave functions \( \chi_{V,n} \) and the masses \( m_{V,n} \), \( V = A, W, Z \), of the Kaluza-Klein modes can be derived from action (12) (here and below the subscript \( n \) denotes the number of the corresponding Kaluza-Klein mode). When we take into account the mass terms generated by spontaneous symmetry breaking, they look like

\[
-m_{A,n}^2 \chi_{A,n} - \partial_4 (e^{2\sigma} \partial_4 \chi_{A,n}) = 0,
\]

\[
-m_{W,n}^2 \chi_{W,n} - \partial_4 (e^{2\sigma} \partial_4 \chi_{W,n}) + m_{W,n}^2 \chi_{W,n} = 0,
\]

\[
-m_{Z,n}^2 \chi_{Z,n} - \partial_4 (e^{2\sigma} \partial_4 \chi_{Z,n}) + m_{Z,n}^2 \chi_{Z,n} = 0.
\]
As usually, we assume that the lowest (zero) Kaluza-Klein modes of the 5D gauge fields coincide with the four-dimensional SM gauge fields. It follows from eq. (11) that the solution for the lowest mode of the field $A_\mu$ (the photon) is $m_{A,0} = 0$ and $\chi_{A,0}(y) \equiv \text{const} = 1$ (the latter equality is due to our normalization of the bulk gauge fields), i.e. its wave function does not depend on the coordinate of the extra dimension. This property of the solution guarantees the universality of the electromagnetic charge \[49\]. The solutions of eqs. (12) and (13) for the wave functions of the lowest modes have the same property, if $m_{W,0} = m_W$ and $m_{Z,0} = m_Z$.

The only case, where the zero mode sector of a five-dimensional model exactly coincides with the electroweak gauge boson sector of the SM, is the one, where the wave functions $\chi_{W,0}(y)$ and $\chi_{Z,0}(y)$ do not depend on the coordinate of the extra dimension. However, to this end the vacuum profile of the 5D Higgs field should be equal to $v \exp(-\sigma(y))$, i.e. it should be fine-tuned with background solution [50] for the metric. In this case the self-coupling constants of the massive gauge bosons are defined in terms of the constants $g$ and $g'$ exactly in the same way as in the ordinary SM. Also in this case the wave functions $\chi_{V,n}(y)$ of the gauge boson excitations defined by eqs. (11), (12), (13) are all equal and below will be denoted by $\chi_n(y)$. The expansions of the 5D gauge fields in KK-modes look like

$$V_\mu(x,y) = \sum_{n=0}^{\infty} V^{(n)}(x) \chi_n(y), V = A, W, Z,$$

and it is easy to check that the following relation holds for these wave functions for an arbitrary KK-number $n > 0$ and an arbitrary power $l > 0$ of the zero mode wave function

$$\int_{-L}^{L} \chi_n(y) (\chi_0(y))^l dy = \int_{-L}^{L} \chi_n(y) \chi_0(y) dy = 0.$$  

Below we consider the case, where the masses of the zero modes of the bulk fields $W_\mu$ and $Z_\mu$ are given by (10) and their wave functions are equal to unity due to our normalization of the bulk gauge fields in action (2). In this case the masses of the first excitations $W', Z'$, and $\gamma'$ are given by

$$m_{W'} = m_{W,1} = \sqrt{m_{A,1}^2 + m_W^2} \simeq m_{\gamma'} + \frac{m_W^2}{2m_{\gamma'}}$$

$$m_{Z'} = m_{Z,1} = \sqrt{m_{A,1}^2 + m_Z^2} \simeq m_{\gamma'} + \frac{m_Z^2}{2m_{\gamma'}},$$

where $m_{A,1} = m_{\gamma'}$ denotes the mass of $\gamma'$ and we have taken into account that $m_{\gamma'} \gg m_{W,Z}$. We emphasize that due to eq. (15) their wave functions $\chi_1(y)$ are orthogonal to the wave functions of the gauge boson zero modes, which are constants. To find the three-particle interactions of the first gauge boson excitations with the SM gauge bosons, we substitute the mode decompositions of the 5D gauge fields $W_\mu(x,y)$, $Z_\mu(x,y)$ and $A_\mu(x,y)$ given by (14) into interaction Lagrangian (8), integrate with respect to the extra dimension coordinate $y$ over the orbifold $S^1/Z_2$ and retain only the terms with both $W', Z'$ or $A'$ and the SM gauge bosons. The resulting effective three-particle interaction 4D Lagrangian of $W_\mu'$, $Z_\mu'$, $A_\mu'$ and the SM
gauge bosons is given by

\[ L_{W,W}^{eff} = - ig \left[ (W_{\mu}^r W_{-\mu}^r - W_{\mu}^l W_{-\mu}^l) (A^\nu \sin \theta_W + Z^\nu \cos \theta_W) + \right. \]
\[ + W_{-\mu}^r W_{\mu}^r (A^{\nu \mu} \sin \theta_W + Z^{\nu \mu} \cos \theta_W) \]
\[ - ig \left[ (W_{\mu}^l W_{-\mu}^l - W_{-\mu}^l W_{\mu}^l) (A^\nu \sin \theta_W + Z^\nu \cos \theta_W) \right. \]
\[ + W_{-\mu}^l W_{\mu}^l (A^{\nu \mu} \sin \theta_W + Z^{\nu \mu} \cos \theta_W) \]
\[ - ig \left[ (W_{\mu}^l W_{-\mu}^r - W_{-\mu}^r W_{\mu}^l) (A^\nu \sin \theta_W + Z^\nu \cos \theta_W) \right. \]
\[ + W_{-\mu}^r W_{\mu}^l (A^{\nu \mu} \sin \theta_W + Z^{\nu \mu} \cos \theta_W) \right]. \] (18)

Similarly we can find the effective four-particle interaction 4D Lagrangian, which turns out to be very bulky (see Appendix).

Both these Lagrangians have the property that, due to orthogonality condition (15), a lowest excitation of the SM gauge bosons cannot interact at tree level with two or three SM gauge bosons. In particular, it means that the decays at tree level of \( W^r, Z^r, \) and \( \gamma^r \) into two or three SM gauge bosons are forbidden. However, these bosons can decay into SM fermions, because the wave functions of the zero modes of the 5D fermions are not constant \([50]\), and the corresponding coupling is defined by the overlap integral of two fermion zero mode wave functions and the wave function \( \chi_1(y) \) of the first gauge boson excitations. These overlap integrals are, of course, model dependent and generally not equal to zero. For example, in papers \([21, 22]\) they have been calculated for the case of the unstabilized RS-model and found to be non-zero. This property also means that the lowest excitations of the SM gauge bosons can decay into two or three SM gauge bosons via triangle or box loop diagrams with SM fermions running in the loops, although the decays are very much suppressed.

However, in the general case, where the vacuum solution for the 5D Higgs field is not fine-tuned, the solutions for the wave functions of the zero modes of the bulk fields \( W_\mu \) and \( Z_\mu \), which correspond to the SM massive gauge bosons, are not necessarily constant, and these decays can also take place due to deviations of the zero mode gauge boson wave functions from unity. In this case eq. (11) remains the same, whereas eqs. (12), (13) take the form:

\[ - (m_{W,n}^2 - m_W^2) \chi_{W,n} - \partial_4 (e^{2\sigma} \partial_4 \chi_{W,n}) + \Delta M_W^2 (y) \chi_{W,n} = 0, \] (19)

\[ - (m_{Z,n}^2 - m_Z^2) \chi_{Z,n} - \partial_4 (e^{2\sigma} \partial_4 \chi_{Z,n}) + \Delta M_Z^2 (y) \chi_{Z,n} = 0, \] (20)

where the extra terms \( \Delta M_V^2 (y) \), \( V = W, Z \), depend on the vacuum profile of the bulk Higgs field and result in deviations of the wave functions \( \chi_{W,0}(y) \) and \( \chi_{Z,0}(y) \) from constant. This has the following well-known consequences. Indeed, the self-coupling of massive gauge bosons comes from the term \( W_{-\mu\nu} W_{\mu\nu} \) and the corresponding coupling constants are defined only by the structure constants of the SM gauge group. In the five-dimensional theory under consideration the self-coupling terms also come from the same term of action (5), but now the coupling constants also include the overlap integrals of the wave functions \( \chi_{W,0}(y) \) and \( \chi_{Z,0}(y) \) over the space of extra dimension. Moreover, it is well known that, in the general case, a modification of the shapes of the zero mode gauge boson wave functions has an influence on the electroweak
observables, which was discussed in detail in [51][52]. It is shown in these papers that, for example, in the case of the unstabilized Randall-Sundrum model [18], a noticeable deviation of the zero mode wave functions from constant may lead to restrictions on the value of the five-dimensional energy scale, which put the theory out of the reach of the present day experiments. For this reason the corrections to the masses and wave functions of the $W$- and $Z$-bosons arising from the terms $\Delta M^2_{V}(y)$, $V = W, Z$, must be very small in order not to influence noticeably the electroweak observables. In this case we can use the standard perturbation theory to find approximate solutions to eqs. (19), (20). To first order of perturbation theory the masses of $W$- and $Z$-bosons and the masses of their first excitations $W'$ and $Z'$ look like

$$
m_{W,0} = \sqrt{m_W^2 + (\Delta M^2_{W})_{00}} \quad (21)$$

$$
m_{Z,0} = \sqrt{m_Z^2 + (\Delta M^2_{Z})_{00}} \quad (22)$$

$$
m_{W'} = m_{W,1} = \sqrt{m_W^2 + m_W^2 + (\Delta M^2_{W})_{00}} \simeq m_{\gamma'} + \frac{m_W^2 + (\Delta M^2_{W})_{00}}{2m_{\gamma'}} \quad (23)$$

$$
m_{Z'} = m_{Z,1} = \sqrt{m_Z^2 + m_Z^2 + (\Delta M^2_{Z})_{00}} \simeq m_{\gamma'} + \frac{m_Z^2 + (\Delta M^2_{Z})_{00}}{2m_{\gamma'}}, \quad (24)$$

the wave functions of the lowest modes $\chi_{V,0}(y)$ and $\chi_{V,1}(y)$, $V = W, Z$, are given by

$$
\chi_{V,0}(y) = \chi_0(y) - \frac{(\Delta M^2_{V})_{00}}{m_{\gamma'}^2} \chi_1(y) - \sum_{n=2}^{\infty} \frac{(\Delta M^2_{V})_{n0}}{m_{A,n}^2} \chi_n(y), \quad (25)
$$

$$
\chi_{V,1}(y) = \chi_1(y) + \frac{(\Delta M^2_{V})_{01}}{m_{\gamma'}^2} \chi_1(y) + \sum_{n=2}^{\infty} \frac{(\Delta M^2_{V})_{n1}}{m_{\gamma'}^2 - m_{A,n}^2} \chi_n(y), \quad (26)
$$

where the matrix elements $(\Delta M^2_{V})_{mn} = (\Delta M^2_{V})_{nm}$ are defined as

$$
(\Delta M^2_{V})_{mn} = \frac{1}{2L} \int_{-L}^{L} \chi_m(y) \Delta M^2_{V}(y) \chi_n(y) dy.
$$

Since $m_{W,0}$ and $m_{Z,0}$ should lie within the experimental uncertainties $\Delta m_W$, $\Delta m_Z$ from the standard masses $m_W$ and $m_Z$, the matrix elements of the perturbations must satisfy the conditions

$$
(\Delta M^2_{W})_{00} \simeq 2m_W(\Delta m_W) \sim 2\text{GeV}^2, \quad (\Delta M^2_{Z})_{00} \simeq 2m_Z(\Delta m_Z) \sim 0.4\text{GeV}^2. \quad (27)
$$

Each perturbation term $\Delta M^2_{V}(y)$, $V = W, Z$, being proportional to a dimensional parameter, the other matrix elements are also of the same order, which means that the corrections to the wave functions in formulas (25), (26) are really very small.

The wave functions $\chi_{V,0}(y)$ and $\chi_{V,1}(y)$, $V = W, Z$, are normalized to unity up to terms of second order in the perturbations, which can be neglected. Due to the orthogonality of the system of the unperturbed wave functions $\{\chi_n(y)\}$ the calculation of the overlap integrals of one wave function $\chi_{V,1}(y)$ (25) with two or three wave functions $\chi_{V,0}(y)$ (26) is very easy and gives the results of the order $(\Delta M^2_{V})_{01} / m_{\gamma'}^2$, which is extremely small because of (27). This means that, in this case, also the decays at tree level of $W'$ and $Z'$ to two or three SM gauge bosons are very much suppressed.
3 Conclusion

In the present paper we have studied the interactions of the electroweak gauge boson excitations in models with warped extra dimensions. It has been found that they are rather different from the interaction properties of these excitations in other models. In particular, we have shown that the couplings of the lowest gauge boson excitations $W'$, $Z'$, and $\gamma'$ to the SM gauge bosons treated as the zero modes of the 5D gauge fields are either exactly equal to zero or very much suppressed. In the former case, we have explicitly found the three-particle and four-particle interaction Lagrangians of the gauge boson excitations $W'$, $Z'$, and $\gamma'$ and the SM gauge bosons. At the same time the couplings of $W'$, $Z'$, and $\gamma'$ to the SM fermions are non-zero allowing for their production and decays. These properties of the gauge boson excitations in models with warped extra dimensions distinguish them from the gauge boson excitations in other models beyond the SM. In particular, in the models with an extension of the SM gauge group [4] these excitations can couple to both the SM gauge bosons and fermions with approximately the same strength. Even in the UED models with flat extra dimension [14–17] the couplings of these excitations to both the SM gauge bosons and fermions are different: they are very much suppressed, and the excitations are expected to be long-lived particles. Thus, the interactions of the electroweak gauge boson excitations are rather different in different extensions of the SM and, if extra vector bosons are found at the LHC, their interaction properties may point out a theory beyond the SM, to which they may belong.

Finally, we would like to note that the discussed property of vanishing or strongly suppressed couplings of the first excited KK modes of the electroweak gauge bosons to the SM gauge bosons can be important not only for searching and interpreting the signals at the LHC, but also for analyzing dark matter scenarios with vector mediators arising in models with extra dimensions.

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Appendix

Substituting the mode decompositions of the 5D gauge fields $W_\mu(x, y)$, $Z_\mu(x, y)$ and $A_\mu(x, y)$ (14) into interaction 5D Lagrangian (9), integrating with respect to the extra dimension coordinate $y$ over the orbifold $S^1/Z_2$ and retaining only the terms with both $W'_\mu$, $Z'_\mu$, or $A'_\mu$ and the SM gauge bosons, we get the following effective three-particle interaction 4D Lagrangian of $W'_\mu$, $Z'_\mu$, $A'_\mu$ and the SM gauge bosons:
\[ L_{WWVV} = - \frac{g^2}{T} \left[ \left( W^\mu_\mu W^{-\mu} \right) \left( W^\nu_\nu W^{-\nu} \right) + \left( W^\mu_\mu W^{\nu-} \right) \left( W^\nu_\nu W^{\nu-} \right) \right] \\
+ 2 \left( W^\mu_\mu W^{\nu-} \right) \left( W^\nu_\nu W^{-\nu} \right) - 2 \left( W^\mu_\mu W^{\nu+} \right) \left( W^\nu_\nu W^{-\nu} \right) \\
- \left( W^\nu_\nu W^{\nu+} \right) \left( W^\nu_\nu W^{-\nu} \right) - \left( W^\mu_\mu W^{\mu+} \right) \left( W^\nu_\nu W^{-\nu} \right) \\
+ 2 \left( W^\mu_\mu W^{\nu-} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
+ 4 \left( W^\mu_\mu W^{\nu-} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
+ 4 \left( W^\mu_\mu W^{\nu+} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
+ 2 \left( W^\mu_\mu W^{\nu-} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 2W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 2W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 2W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 2W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 2W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
+ g_{eff} \left[ 2 \left( W^\mu_\mu W^{\nu-} \right) \left( W^\nu_\nu W^{-\nu} \right) + 2 \left( W^\mu_\mu W^{\nu-} \right) \left( W^\nu_\nu W^{-\nu} \right) \right] \\
- 2 \left( W^\mu_\mu W^{\nu+} \right) \left( W^\nu_\nu W^{-\nu} \right) - 2 \left( W^\mu_\mu W^{\nu+} \right) \left( W^\nu_\nu W^{-\nu} \right) \\
+ 6 \left( W^\mu_\mu W^{\nu-} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
+ 3 \left( W^\mu_\mu W^{\nu-} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
+ 3 \left( W^\mu_\mu W^{\nu-} \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 3W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 3W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 3W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \\
- 3W^\mu_\mu \left( A^\nu - Z^\nu \cos \theta_W \right) \left( A^\nu - Z^\nu \cos \theta_W \right) \right],

where

\[ g_{eff} = \frac{g^2}{2L} \int_{-L}^{L} (\chi_1(y))^3 dy. \]
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