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Bulk antisymmetric tensor fields coupled to a dilaton in a Randall-Sundrum model

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A string-inspired three-form-dilaton-gravity model is studied in a Randall-Sundrum brane world scenario. As expected, the rank-3 antisymmetric field is exponentially suppressed. For each mass level, the mass spectrum is bigger than the one for the Kalb-Ramond field. The coupling between the dilaton and the massless Kaluza-Klein mode of the three-form is calculated, and the coupling constant of the cubic interactions is obtained numerically. This coupling are of the order of Tev^{-1}; therefore, there exists a possibility to find some signal of it at Tev scale.

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I. INTRODUCTION

The core idea of extra dimension models is to consider the four-dimensional universe as a hyper-surface embedded in multidimensional manifold. The appeal of such models is the determination of scenarios where membranes have the best chances to mimic the standard model's characteristics. In particular, the standard model presents interesting topics to study, such as the hierarchy problem and the cosmological constant problem that can be treated by the above-mentioned models. For example, the Randall-Sundrum model [1,2] provides a possible solution to the hierarchy problem and show how gravity is trapped to a membrane.

There are studies that emphasize the general properties of membranes, for example, one that makes an analysis of the space-time singularities that arise in braneworld models with four-dimensional (4D) Poincare invariance [3]. Others with cosmological implications consider our universe initially as an empty p-brane embedded in a manifold, whose bulk is populated only with a dilaton field that couples itself to the brane, producing a rich and interesting evolutionary cosmology [4]. The presence of other fields in the bulk introduces the problem of field localization in the brane, which is an important tool in order to build up the standard model in the membrane setup. For such, several ingredients had to be added: the gravitational field, spinor fields, scalar fields, and gauge fields. A lot of aspects have been studied related to these topics. These investigations include both the smooth models described by solitonlike membranes and the pure Randall-Sundrum models containing space-time singularities. For these, there is an extensive list of references in literature [5–16].

Despite its achievements, that model does not favor the localization of zero modes of gauge vector fields. Such localization is important to construct the standard model in a membrane [17]. That happens because the gauge vector field theory is conformal in $D = 4$: this makes the warp factor in the pure Randall-Sundrum model disappear from the effective action. The issue of localization of standard model fields on the three-brane has been widely studied [18–22]. Despite this, we must consider a scenario in which it is assumed *a priori* that the matter fields are constrained to live on the three-brane.

Some studies about higher rank tensor fields have been made showing its relation with the AdS/CFT conjecture [23]. Besides this, string theory shows the naturalness of higher rank tensor fields in its spectrum [24,25].

In mathematical terms, the presence of one more extra dimension ($D = 5$) provides the existence of many antisymmetric fields, namely, the two-, three-, four-, and five-forms. However, the only relevant ones for the visible brane are the two- and three-form. This is due to the fact that when the number of dimensions increase, the number of gauge freedom also increases. This can be used to cancel the dynamics of the field in the visible brane. It is important to point out that this happens only because we have an interacting theory. If this was not so, we could cancel more degrees of freedom using the equations of motion and the three-form would not have dynamics [26]. The mass spectrum of the two- and three-form have been studied, for example, in Refs. [27,28].

The coupling between the two-form and the dilaton, with cosmological consequences, has been studied in [29]. There are also authors that investigate that coupling in the domain of standard model physics [30]. Such coupling, inspired in string theory, can provide us with a process that, in principle, could be observed in the LHC. That may happen through a Drell-Yang process, in which a pair of quark-antiquark can give rise to a three-form field, mediated by a dilaton. This kind of process can appear in a scenario where the dilaton is considered as the Higgs field [31–33]. This raises the question if the coupling between the three-form and the dilaton give a similar process, and this is the goal of this article.

The question to be addressed in this piece of work is related to other kinds of gauge fields. We present a full action, including the gravity, the dilaton, and the three-form field action. The dimensional reduction is done, and the effective action at the visible brane is obtained. Then, we analyze carefully the dilaton sector, find the mass spectrum, and find the complete solution for dilaton coefficients. This result will be needed when the coupling with the three-form is analyzed. After, we analyze the three-form sector and, like in the dilaton case, we find the mass spectrum and solve the equation to find the coefficients of the expansion. The mass spectrum of the three-form is compared with that of the two-form, the dilaton, and the gravity field. This gives a hint that the three-form must not be seen at LHC, because of the mass of the massive Kaluza-Klein modes. Finally, we study the coupling with the dilaton and find how we can conclude that a signal of the massless three-form can be found at LHC.

II. DILATON AND ANTISYMMETRIC TENSOR FIELDS IN RANDALL-SUNDRUM FRAMEWORK

As the number of antisymmetric tensor fields increase with dimensions, these fields has to be considered. In fact these fields have been taken into account in the literature. As said in the introduction, the antisymmetric tensors relevant to the visible brane are that of rank two and three [28]. We focus our attention here in the second tensor field cited. As usual, capital Latin index represent the coordinates in the bulk and Greek index, those on the brane. The metric is given by

$$ds^2 = e^{-2\sigma}\eta_{\mu\nu}dx^\mu dx^\nu + r_c^2 dy^2, \tag{1}$$

where $\sigma = kr_c y$, y is the coordinate of fifth dimension, r_c is the compactification radius for that dimension, k is a constant of the order of the higher dimensional Planck mass M, and $\eta_{\mu\nu}$ is the 4D Minkowski metric.

The coupling between the two-form and the dilaton was considered recently by Mukhopadhyaya [30]. We consider here that the three-form couples in the same way to the dilaton. The motivation for this is that, if our solution is lifted to six dimensions, upon compactification in S^1, we obtain the same coupling for the matter fields. This also justifies the name dilaton for ϕ [14]. The action to be considered is therefore given by

$$S = S_{\rm grav} + S_X + S_{\rm dil}, \tag{2}$$

where

$$S_{\rm grav} = \int d^5x\sqrt{-G}2M^3R,$$
$$S_X = \int d^5x\sqrt{-G}[-e^{\phi/M^{3/2}}2Y_{MNLP}Y^{MNLP}],$$

and

$$S_{\rm dil} = \int d^5x\sqrt{-G}\Big[\frac{1}{2}\partial_M\phi\partial^M\phi - m^2\phi^2\Big]. \tag{3}$$

Defining the three-form field X_{MNO} and fixing the gauge $X_{\mu\nu y} = 0$, we have two possibilities for the strength tensor

$$Y_{\mu\nu\alpha\beta} = \partial_{[\mu}X_{\nu\alpha\beta]}$$

or

$$Y_{y\nu\alpha\beta} = \partial_{[y}X_{\nu\alpha\beta]}.$$

We must here focus only in the dilaton and the antisymmetric field X_{MNO}. Thus, for the metric in (1), we have

$$S_X = \int d^4x\int dy r_c e^{-4\sigma}\{A + B + C\},$$

where

$$A = -2e^{\phi/M^{3/2}}e^{8\sigma}\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\delta}\eta^{\gamma\tau}Y_{\alpha\mu\lambda\gamma}Y_{\beta\nu\delta\tau}, \tag{4}$$

$$B = -2e^{\phi/M^{3/2}}\frac{4}{r_c^2}e^{6\phi}\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma}\partial_y X_{\alpha\mu\lambda}\partial_y X_{\beta\nu\gamma}, \tag{5}$$

and

$$C = r_c e^{-4\sigma}\Big[\frac{e^{2\sigma}}{2}\eta^{\mu\nu}\partial_\mu\phi\partial_\nu\phi + \frac{1}{2r_c^2}\partial_y\phi\partial_y\phi - m^2\phi^2\Big]. \tag{6}$$

As the mass term of the dilaton decouples the visible brane, we must consider the kinetic term only. Performing an integration by parts in the second terms of the dilaton and X fields, respectively, we get

$$S = \int d^4x\int dy\Big[r_c e^{-4\sigma}A + e^{\phi/M^{3/2}}\frac{8}{r_c}\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma}X_{\alpha\mu\lambda}\partial_y(e^{2\phi}\partial_y X_{\beta\nu\gamma}) + \frac{r_c e^{-2\sigma}}{2}\eta^{\mu\nu}\partial_\mu\phi\partial_\nu\phi - \frac{1}{2r_c}\phi\partial_y(e^{-4\sigma}\partial_y\phi)\Big].$$

Note that the derivative in the exponential of the dilaton would give a term suppressed by the Planck mass and therefore is irrelevant for the effective action. At this point, we must consider the Kaluza-Klein decomposition of the fields

$$X_{\mu\nu\alpha} = \sum_{n=0}^{\infty}X^n_{\mu\nu\alpha}(x)\frac{\chi^n(y)}{\sqrt{r_c}}, \qquad \phi = \sum_{n=0}^{\infty}\phi^n(x)\frac{\psi^n(y)}{\sqrt{r_c}},$$

where

$$\int e^{4\sigma}\chi^m(y)\chi^n(y)dy = \delta^{mn},$$
$$\int e^{-2\sigma}\psi^m(y)\psi^n(y)dy = \delta^{mn}. \tag{7}$$

In terms of the above projections, the effective action on the visible brane is given by

$$S = \int d^4x \sum_{n=0}^{\infty}\sum_{m=0}^{\infty}\int dy \bigg\{ e^{\phi/M^{3/2}}\bigg[-2e^{4\sigma}\chi^m\chi^n\eta^{\alpha\beta}\eta^{\mu\nu} \times \eta^{\lambda\delta}\eta^{\gamma\tau}Y^n_{\alpha\mu\lambda\gamma}Y^m_{\beta\nu\delta\tau} + \frac{8}{r_c^2}\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma}X^n_{\alpha\mu\lambda}X^m_{\beta\nu\gamma}\chi^m \times \frac{d}{dy}\bigg(e^{2\phi}\frac{d}{dy}\chi^n\bigg)\bigg] + \bigg[\frac{e^{-2\sigma}}{2}\psi^m\psi^n\eta^{\mu\nu}\partial_\mu\phi^m\partial_\nu\phi^n - \frac{1}{2r_c^2}\phi^m\phi^n\psi^m\partial_y(e^{-4\sigma}\partial_y\psi^n)\bigg]\bigg\}.$$

The action above is the action of the rank three antisymmetric tensor field, or three-form, coupled to a dilaton. The first bracket represents the three-form coupled to a dilaton and the second the dilaton itself. As there is no coupling of the terms in the second brackets and the three-form fields, it is a free dilaton action. So we must concentrate in the free dilaton action and later go on to the main objective, which is the antisymmetric field.

III. BULK DILATON FIELD

At this section, we must consider the dilaton sector. Here we must expand the dilaton exponential

$$e^{\phi/M^{3/2}} = 1 + (\phi/M^{3/2}) + (\phi/M^{3/2})^2/2! + \cdots$$

and concentrate only in the first term which will give the usual kinetic and mass term for the dilaton field provided, so we have

$$-\frac{1}{r_c^2}\partial_y(e^{-4\sigma}\partial_y\psi^n) = (m_n^d)^2e^{-2\sigma}\psi^n. \tag{8}$$

The other terms of the expansion will be considered later as interaction terms in the effective action. Using the above orthonormality conditions, we finally get

$$S_{\rm dil} = \int d^4x \sum_{n=0}^{\infty}\frac{1}{2}[\eta^{\mu\nu}\partial_\mu\phi^n\partial_\nu\phi^n + (m_n^d)^2\phi^n\phi^n].$$

Therefore, we obtain a standard dilaton action with masses given by the solutions of Eq. (8). The solution is given in [30] and we will basically repeat it here. The easiest solution is for the massless dilaton, where we have to solve the equation

$$-\frac{1}{r_c^2}\partial_y(e^{-4\sigma}\partial_y\psi^n) = 0$$

with the obvious solution

$$\psi^0 = \frac{C_1}{4kr_c}e^{4\sigma} + C_2.$$

The constants can be obtained from orbifold condition $\psi^0(-\pi) = \psi^0(\pi)$ and from orthonormality conditions, again. We find $C_1 = 0$ and, considering also that $\exp(-2kr_c\pi) \ll 1$, $C_2 = \sqrt{kr_c}$. The final solution is

$$\psi^0 = \sqrt{kr_c}.$$

In order to solve the equations for the massive modes first, we must make the change of variables $z_n = \frac{m_n^d}{k}e^{\sigma}$ and $f^n = \psi^n e^{-2\sigma}$ to obtain the second order equation

$$\bigg[z_n^2\frac{d^2}{dz_n^2} + z_n\frac{d}{dz_n} + (z_n^2 - 4)\bigg]f_n = 0,$$

which admits a Bessel function of order 2 as a solution. The general solution is

$$\psi^n = \frac{e^{2\sigma}}{N_n}[J_2(z_n) + \alpha_n Y_2(z_n)],$$

where N_n and α_n are constants to be determined. First of all, we must use the continuity condition for the derivative of ψ^n at $y = 0$. Remembering that J_2 and Y_2 are Bessel and Neumann functions of order 2, we obtain

$$\alpha_n = -\frac{J_1(\frac{m_n^d}{k})}{Y_1(\frac{m_n^d}{k})}.$$

As the masses m_n^d are expected to be of order of Tev on the brane, we have $m_n^d \ll k$, and expanding the right-hand side of the above equation, we obtain

$$\alpha_n = -\frac{\pi}{4}\bigg(\frac{m_n^d}{k}\bigg)^2 \ll 1.$$

Using the boundary condition at $y = \pi$, we obtain

$$J_1(x_n) = 0, \tag{9}$$

where we have defined $x_n = z_n(\pi) = \frac{m_n^d}{k}e^{kr_c\pi}$.

If we find the solutions of the equation above, we can find x_n and therefore the mass spectrum. Now the orthogonality condition can be used to determine N_n, and we obtain for the solution

$$\psi^n(z_n) = \sqrt{kr_c}\frac{e^{2\sigma}}{e^{kr_c\pi}}\frac{J_2(z_n)}{J_2(x_n)}.$$

Now we have the complete solution for the dilaton given by

$$\phi = \sqrt{kr_c}\phi^0 + \sum_{n=1}^{\infty}\sqrt{k}\frac{e^{2\sigma}}{e^{kr_c\pi}}\frac{J_2(z_n)}{J_2(x_n)}\phi^n$$

with masses given by Eq. (9). Posteriorly, we will compare the mass of the dilaton with the graviton one. The complete solution for the dilaton will be necessary for studying its coupling to the antisymmetric field.

IV. BULK THREE-FORM FIELD

Now we must consider the three-form field. The action was constructed before and we have

$$S_X = \int d^4x \sum_{n=0}^{\infty}\sum_{m=0}^{\infty}\int dy e^{\phi/M^{3/2}}\Big[-2e^{4\sigma}\chi^m\chi^n\eta^{\alpha\beta}\eta^{\mu\nu} \times \eta^{\lambda\delta}\eta^{\gamma\tau}Y^n_{\alpha\mu\lambda\gamma}Y^m_{\beta\nu\delta\tau} + \frac{8}{r_c^2}\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma}X^n_{\alpha\mu\lambda}X^m_{\beta\nu\gamma} \times \chi^m\frac{d}{dy}\Big(e^{2\phi}\frac{d}{dy}\chi^n\Big)\Big]. \quad (10)$$

Again we must expand of the dilaton exponential and consider only the first term. The mass spectrum is obtained from

$$-\frac{1}{r_c^2}\frac{d}{dy}\Big(e^{2\phi}\frac{d}{dy}\chi^n\Big) = (m_X^n)^2\chi^n e^{4\sigma},$$

and again the other terms of the expansion will be considered later as interaction terms in the effective action. The simplest and most important solution of the above equation is obtained for the massless case, where we have the equation

$$-\frac{1}{r_c^2}\frac{d}{dy}\Big(e^{2\phi}\frac{d}{dy}\chi^n\Big) = 0$$

with obvious solution

$$\chi^0 = -\frac{C_1}{2kr_c}e^{-2\sigma} + C_2.$$

Using continuity of the derivative at $y = \pi$ gives us $C_1 = 0$ and from orthogonality, we finally get

$$\chi^0 = \sqrt{2kr_c}e^{-2kr_c\pi}.$$

For the massive modes, similarly to the dilaton case, we perform the redefinitions $z'_n = \frac{m_n}{k}e^{\sigma}$ and $f'_n = e^{\sigma}\chi^n$ to obtain the equation

$$\Big[z_n'^2\frac{d^2}{dz_n'^2} + z'_n\frac{d}{dz'_n} + (z_n'^2 - 1)\Big]f'_n = 0$$

with solution given by a Bessel function of order 1. Therefore,

$$\chi^n = e^{-\sigma}f'_n = \frac{e^{-\sigma}}{N'_n}[J_1(z'_n) + \alpha'_n Y_1(z'_n)],$$

and again we have to determine the constants N'_n and α'_n using contour conditions. First of all, we must use continuity conditions at $y = 0$ to obtain

$$\alpha'_n = -\frac{J_2(\frac{m_X^n}{k})}{Y_2(\frac{m_X^n}{k})}.$$

As the masses are in Tev scale, we have $m_X^n \ll k$ and expanding the above expression, we get

$$\alpha'_n \sim \frac{\pi}{2^5}\Big(\frac{m_X^n}{k}\Big)^4 \ll 1.$$

Using the above result, the fact that $e^{kr_c\pi} \gg 1$, and the contour condition at $y = \pi$, we obtain

$$J_2(x'_n) = 0 \quad (11)$$

with the definition $x'_n = z'_n(\pi) = \frac{m_X^n}{k}e^{kr_c\pi}$. Therefore, we can obtain the mass spectrum and, in the visible brane, the masses are solutions of the Eq. (11). The effective action is given by

$$S_X = \int d^4x \sum_{n=0}^{\infty} e^{\phi/M^{3/2}}\{-2\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\delta}\eta^{\gamma\tau}Y^n_{\alpha\mu\lambda\gamma}Y^m_{\beta\nu\delta\tau} - 8(m_X^n)^2\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma}X^n_{\alpha\mu\lambda}X^m_{\beta\nu\gamma}\}. \quad (12)$$

We can obtain N_n with the normalization condition to obtain

$$N'_n = \frac{e^{kr_c\pi}}{\sqrt{kr_c}}J_1(x'_n),$$

and we arrive at the final solution for the massive modes given by

$$\chi^n(z_n) = \sqrt{kr_c}\frac{e^{\sigma}}{e^{kr_c\pi}}\frac{J_1(z'_n)}{J_1(x'_n)}.$$

It is interesting to note that, despite the way we have defined it, the mass spectrum of the three-form field is no altered by the dilaton, and the only change will be in the interaction terms, that will be analyzed posteriorly. We list in Table I the value of the masses for the graviton, the dilaton, the two-form, and three-form, per [28,30], using the scale $kr_c = 12$ and $k = 10^{19}$ Gev.

It is obvious from the Table I that, the higher the rank of the tensor is, the higher the mass spectrum is, and the possibility of a finding these fields in the LHC becomes more and more elusive. The possibility for some signal of the X field in LHC will come from its interaction with the dilaton field, and this is analyzed in the next section.

TABLE I. Masses of KK modes where $kr_c = 12$ and $k = 10^{19}$ Gev.

n	1	2	3	4
$m^n_{\rm grav}$ (TeV)	1.66	3.04	4.40	5.77
$m^n_{\rm dil}$ (TeV)	1.66	3.04	4.40	5.77
m^n_{KR} (TeV)	2.87	5.26	7.62	9.99
m^n_X (TeV)	4.44	7.28	10.05	12.79

V. THREE-FORM COUPLED TO DILATON

As we explore here only the coupling between the three-form and the dilaton, we must consider only the relevant part of the action, which is given by

$$S_X = \int d^4x \sum_{n=0}^{\infty} \sum_{m=0}^{\infty} \int dy e^{\phi/M^{3/2}}[-2e^{4\sigma}\chi^m\chi^n\eta^{\alpha\beta}\eta^{\mu\nu} \times \eta^{\lambda\delta}\eta^{\gamma\tau}Y^n_{\alpha\mu\lambda\gamma}Y^m_{\beta\nu\delta\tau} - 8(m^n_X)^2\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma} \times X^n_{\alpha\mu\lambda}X^m_{\beta\nu\gamma}\chi^m\chi^n e^{4\sigma}]. \tag{13}$$

When we expand the exponential of the dilaton, the first term will give us the usual kinetic and mass term studied previously. The next terms of the expansion are

$$\exp(\phi/M^{3/2}) - 1 = \left(M^{-3/2}\sum_{n=0}^{\infty}\phi^n(x)\frac{\psi^n(y)}{\sqrt{r_c}}\right) + \frac{1}{2!}\left(M^{-3/2}\sum_{n=0}^{\infty}\phi^n(x)\frac{\psi^n(y)}{\sqrt{r_c}}\right)^2 + \cdots \tag{14}$$

The terms beyond the first are highly suppressed by powers of the plank mass; therefore, we must consider only the first term, or

$$\exp(\phi/M^{3/2}) - 1 \sim M^{-3/2}\sum_{n=0}^{\infty}\phi^n(x)\frac{\psi^n(y)}{\sqrt{r_c}} = M^{-3/2}\left(\sqrt{kr_c}\phi^0 + \sum_{n=1}^{\infty}\sqrt{k}\frac{e^{2\sigma}}{e^{kr_c\pi}}\frac{J_2(z_n)}{J_2(x_n)}\phi^n\right). \tag{15}$$

The above term will give a coupling of the form ϕX^2. As we can see from the complete solution of the dilaton and the three-form, the interaction of the massless fields will be suppressed by a factor of $\frac{1}{M_p}$ and therefore is not considered. The interaction between a massless dilaton, a massless two-form, and a massive three-form is like $\phi^0X^0X^n$ and gives us a null result because of the orthogonality relations.

Therefore, the interaction that posses massless dilatons are irrelevant, and we must consider only the massive ones. Taking in account these considerations and rearranging terms in the action, we get for the cubic interactions

$$S_X \supset S_{\text{cubic}} = \sum_{n,m=0}^{\infty}\sum_{p=1}^{\infty}\frac{M^{-3/2}}{\sqrt{r_c}}\int dy e^{4\sigma}\chi^m(y)\chi^n(y)\psi^p(y) \times \int d^4x\phi^p(x)[-2\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\delta}\eta^{\gamma\tau}Y^n_{\alpha\mu\lambda\gamma} \times Y^m_{\beta\nu\delta\tau} - 8(m^n_X)^2\eta^{\alpha\beta}\eta^{\mu\nu}\eta^{\lambda\gamma}X^n_{\alpha\mu\lambda}X^m_{\beta\nu\gamma}]. \tag{16}$$

From the effective action, we can see that the coupling constant for these interactions are given by

$$\frac{M^{-3/2}}{\sqrt{r_c}}\int_{-\pi}^{+\pi} dy e^{4\sigma}\chi^m\chi^n\psi^i.$$

As said before, due the large values of the mass for the Kaluza-Klein (KK) modes, their signal must hardly be seen at LHC. Therefore, we concentrate in the more interesting case, that is the massless mode. Using now the solution for the coefficients above we obtain, for the massless modes

$$\frac{2kr_c}{M_P}e^{-5kr_c\pi}\int_{-\pi}^{+\pi} dy e^{6\sigma}\frac{J_2(z_n)}{J_2(x_n)}.$$

The above integral was solved numerically using the software MATHEMATICA, and the results for the first four KK modes are given in Table II.

TABLE II. Coupling constants in units of Gev^{-1}.

n	Coupling constant
1	0.00143676
2	0.000788648
3	0.000413211
4	0.000249329

The values of the coupling constant in the above table are given in Gev^{-1}. This shows a rather interesting possibility: Despite the fact that the massless mode is extremely suppressed, its coupling with the dilaton rises the possibility of a signal at LHC. As pointed in [30], the Lorenz structure of the interactions and masses involved will change the angular distribution of X fields produced. Furthermore, in the a Drell-Yan process, the production of X zero modes through, for example, the interaction of quark-antiquark pairs, may have different (but significant) rates for sufficient integrated luminosity if we compare with the case of the Kalb-Ramond and gravity fields.

VI. CONCLUSIONS

Extending earlier results about antisymmetric fields [28], we have shown here that the nontrivial coupling between the dilaton and the antisymmetric field do not affect its high suppression in our visible brane due to the warp factor of the Randall-Sundrum scenario.

In the dilaton background, we have studied new interactions between the Kaluza-Klein modes of the dilaton field with the antisymmetric field. These results are important from the phenomenological viewpoint. The interactions that possess massless dilatons are irrelevant, and we must consider only the massive ones. As we have shown, higher order interactions of the massless fields are suppressed by a factor of $\frac{1}{M_p}$ and therefore are not considered. Because of the large mass of the massive modes, we considered only the massless case. In fact, numerical computation of the important interaction is of order of Tev^{-1} and can give us signals of this field in the LHC's searches through Drell-Yan processes mediated by massive dilaton.

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