PROCEEDINGS OF SCIENCE

KK Gluons at NLO

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The Kaluza-Klein (KK) excitation of the gluon is one of the important signatures of the Bulk Randall-Sundrum (Bulk RS) model. After briefly reviewing the model and the significance of the KK gluon search in the context of this model, the production of KK-gluons from gluon initial states at next-to-leading order in QCD is discussed.

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1. Introduction

The Randall-Sundrum (RS) model [1] is a 5-dimensional model which provides a new way of addressing the hierarchy problem. In its simplest version this model is five-dimensional model and the fifth dimensionstrong curvature, is compactified on a $\mathbf{S}^1/\mathbb{Z}_2$ orbifold. The radius of compactification R_c is small and of the order of the Planck length. At the orbifold fixed points $\phi = 0,\ \pi$, are located two branes, the Planck brane and the TeV brane, respectively. The Standard Model fields are localised on the TeV brane while the gravitons exist in the full five-dimensional spacetime. The model uses a warped metric:

$$ds^2 = e^{-kR_c\phi}\eta_{\mu\nu}dx^\mu dx^\nu \ + \ R_c^2 d\phi^2\,; \tag{1.1}$$

k is a mass scale related to the curvature; $\exp(-kR_c\phi)$ is the warp factor which rescales masses of fields localised on the TeV-brane. The electroweak hierarchy $\frac{M_P}{M_{EW}} \sim 10^{15}$ can be generated by an exponent of order 30 and thus the model provides a solution to the hierarchy problem provided the radius R_c is stabilised against quantum fluctuations. This is achieved by a stabilising potential generated by introducing a bulk scalar field [2]. Kaluza-Klein (KK) excitations of the graviton constitute the signature for this model and several collider implications of these graviton resonances have been studied in the literature [3].

The AdS/CFT correspondence [4] relates the RS model (which is a gravity theory in AdS) to a dual theory – a strongly-coupled gauge theory in four dimensions [5]. In the dual description, the fields localised on the TeV brane are TeV-scale composites of the strongly interacting theory making the RS model dual to a composite extension of the SM. Such a composite theory is unviable and leads one inexorably to consider modifying the original RS model. The simplest possibility is to modify the model so that only the Higgs field is localised on the TeV brane while the rest of the SM fields are in the bulk [6, 7] so that the Higgs in such a model is composite while the other SM fields inherit a degree of compositeness via their interaction with the Higgs.

In constructing such variants of the RS model, one has to carefully avoid the constraints coming from flavour hierarchy, electroweak precision tests and flavour-changing neutral currents [8]. In particular, in order to avoid an unacceptably large contribution to the electroweak T parameter an enhanced symmetry in the bulk like $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{(B-L)}$ may be required. Moreover, to acquire a large Yukawa coupling through a larger overlap with the Higgs wavefunction, the profiles of the heavier fermions, like the top, need to be peaked closer to the TeV-brane. But this cannot be done naively: the left-handed doublet, $(t,b)_L$, cannot be close to the TeV brane because that induces non-universal couplings of the b_L to the Z strongly constrained by R_b. The solution is to have the doublet as far away from the TeV brane as required by R_b, while the t_R is localised close to the TeV brane to account for the large Yukawa of the top. Even with this choice of profiles the bounds on the masses of the KK gauge bosons, coming from $Z \rightarrow b\bar{b}$ are found to be in the region of 5 TeV. A custodial symmetry can be invoked to relax this constraint and it also allows other choices of profiles for the t_R and $(t,b)_L$. With this custodial symmetry and for appropriate choices of the profiles for the for the t_R and $(t,b)_L$ it is found that gauge boson masses as low as 2-3 TeV can be consistent with the constraint from $Z \rightarrow b\bar{b}$. [8] A review of the literature on this subject can be found in reference [9].

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The t_R localised close to the TeV brane has an enhanced coupling to the first KK excitation of the gluon in the bulk and as a result we expect that $g_{KK} \to t\bar{t}$ (where g_{KK} represents the first KK excitation of the gluon) will be a significant decay mode from a discovery perspective. The decay also lends itself to identification via spin determination, since the enhanced coupling to t_R over t_L means the top quarks from KK gluon decays will be polarised [10]. There will be additional challenges for identification of the $t\bar{t}$ pairs because they will be highly boosted in the lab frame, but this channel remains a promising search channel.

From a hadron collider perspective, we are interested in the production process $pp \to g_{KK}$. The subprocess $q\bar{q} \to g_{KK}$ has already been investigated in some detail for the LHC [10, 11], as well as for the Tevatron [12], but because in many models the light quarks have a relatively suppressed coupling to the g_{KK}, it is worth also considering the process $gg \to g_{KK}$, even though this process is one-loop at leading order. This has been studied in [13] and the details of this calculation is reported here [1].

2. The NLO calculation

The theoretical strategy that we adopt in computing the loop diagrams needed to study the process $gg \to g_{KK}$ closely follows what is detailed in Refs. [16, 17]. We work with the reduced amplitude $F_{\mu\nu\rho}(p,q)$ related to the the matrix element $\mathscr{M}(p,q)$ via:

$$\mathscr{M}(p,q) = \varepsilon^{\rho *}_{g_{KK}}(r)\varepsilon^{\mu}(p)\varepsilon^{\nu}(q)F_{\mu\nu\rho}(p,q)\,. \tag{2.1}$$

where $\varepsilon^{\mu}(p)$ and $\varepsilon^{\nu}(q)$ are the polarisations of the incoming gluons and $\varepsilon^{\rho}_{g_{KK}}(r)$ is that of the outgoing KK-gluon. We may simplify the calculation by deriving a general form that must be taken by the reduced amplitude using QCD current conservation which implies the following:

$$p^{\mu}F_{\mu\nu\rho} = 0\,, \tag{2.2}$$

$$q^{\nu}F_{\mu\nu\rho} = 0\,. \tag{2.3}$$

Expanding about $p = 0$ and $q = 0$, and after some calculation using the above relations it can be shown that

$$\begin{aligned} F_{\mu\nu\rho} = {} & A\left(\eta_{\mu\nu}\, p\cdot q - q_{\mu}p_{\nu}\right)p_{\rho} + B\varepsilon_{\mu\nu\gamma\delta}p^{\gamma}q^{\delta}p_{\rho} + \\ & + C\left(\varepsilon_{\mu\nu\rho\gamma}p^{\gamma}\, p\cdot q - \varepsilon_{\mu\rho\gamma\delta}p^{\gamma}q^{\delta}p_{\nu}\right) + D\left(\varepsilon_{\mu\nu\rho\gamma}q^{\gamma}\, p\cdot q - \varepsilon_{\nu\rho\gamma\delta}p^{\gamma}q^{\delta}q_{\mu}\right), \end{aligned} \tag{2.4}$$

where A, B, C and D are constants. The problem of calculating the amplitude reduces to the problem of calculating A, B, C and D.

[1] Other, tree-level, processes, involving gg fusion to g_{KK} in association with additional top quark production, have also been considered previously [14, 15].

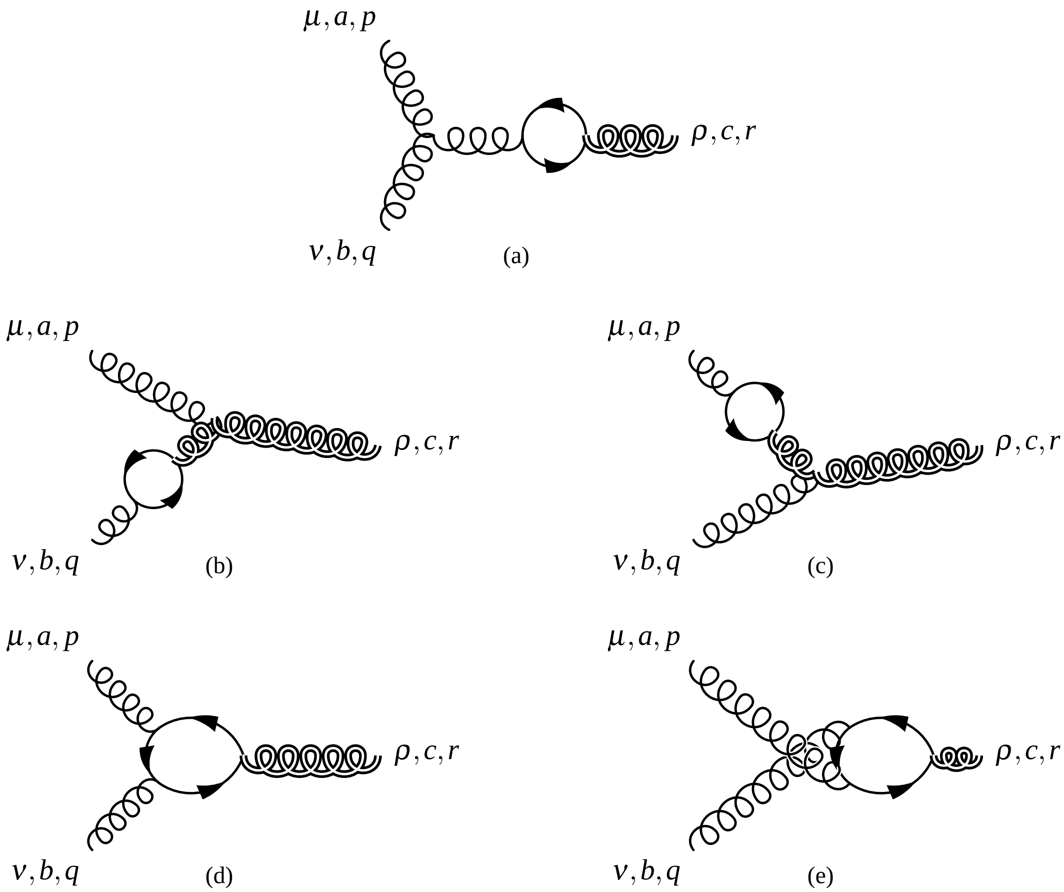

Figure 1: Feynman diagrams for the process that involve a quark in the loop.

Figure 1 contains diagrams for the process that have a quark in the loop. We may write the contributions to the amplitude from the individual diagrams as

$$F^{(q:a)}_{\mu\nu\rho} = -\frac{1}{2} i g^2 g^{(1q)} f^{adb} \delta^{de} \mathrm{Tr}(t^c t^e) \left[\eta_{\mu\alpha}(p+r)_\nu + \eta_{\alpha\nu}(-r-q)_\mu + \eta_{\mu\nu}(q-p)_\alpha\right] \times$$
$$\times\, \eta^{\alpha\beta} \frac{1}{r^2} \int \frac{d^4 l}{(2\pi)^4} \frac{\mathrm{Tr}\left[\gamma_\beta(\not{l} - m_q)\gamma_\rho(1\pm\gamma^5)(\not{l} + \not{r} - m_q)\right]}{[l^2 - m_q^2][(l+r)^2 - m_q^2]}\,, \tag{2.5}$$

$$F^{(q:b)}_{\mu\nu\rho} = -\frac{1}{2} i g^2 g^{(1q)} \mathrm{Tr}(t^b t^d) \delta^{de} f^{aec} \frac{\eta^{\alpha\beta}}{q^2 - M_{KK}^2} \left[\eta_{\mu\beta}(p-q)_\rho + \eta_{\beta\rho}(q+r)_\mu + \eta_{\rho\mu}(-r-p)_\beta\right] \times$$
$$\times \int \frac{d^4 l}{(2\pi)^4} \frac{\mathrm{Tr}\left[\gamma_\nu(\not{l} - m_q)\gamma_\alpha(1\pm\gamma^5)(\not{l} + \not{q} - m_q)\right]}{[l^2 - m_q^2][(l+q)^2 - m_q^2]}\,, \tag{2.6}$$

$$F^{(q:d)}_{\mu\nu\rho} = -\frac{1}{2} g^2 g^{(1q)} \mathrm{Tr}(t^a t^c t^b) \times$$
$$\times \int \frac{d^4 l}{(2\pi)^4} \frac{\mathrm{Tr}\left[\gamma_\mu(\not{l} - \not{p} - m_q)\gamma_\rho(1\pm\gamma^5)(\not{l} + \not{q} - m_q)\gamma_\nu(\not{l} - m_q)\right]}{[(l-p)^2 - m_q^2][(l+q)^2 - m_q^2][l^2 - m_q^2]}\,, \tag{2.7}$$

$$\tag{2.8}$$

denoting the contribution from quark loop diagram a by $F^{(q:a)}_{\mu\nu\rho}$, etc. and where $F^{(q:c)}_{\mu\nu\rho}$ is obtained from $F^{(q:b)}_{\mu\nu\rho}$ by $p \leftrightarrow q$ as is $F^{(q:e)}_{\mu\nu\rho}$ from $F^{(q:d)}_{\mu\nu\rho}$.

There are no diagrams with a gluon in the loop, since there is no vertex containing gluons and a single KK gluon. However, there are diagrams with KK gluons (and KK ghosts) in the loop which have all been calculated in Ref. [13] but which we do not present here in the interests of brevity. However, we go through the quark loops discussion in some detail.

We observe that there are four counterterm diagrams and that the Lorentz structure of the counterterms is derived from consideration of the underlying Lagrangian term. In particular, each term in the three-point counterterm vertex has one momentum factor carrying a Lorentz index (the other two being carried by a metric factor), and the two-point counterterm vertex has the sum of a term where the metric carries both external Lorentz indices and a term where there are two momentum terms each carrying an external Lorentz index. We note that this latter term always contains a momentum factor that contracts with an external polarisation vector to give zero. It is therefore the case that none of the counterterm diagrams contains a term where there is more than one momentum factor carrying an external Lorentz index. (There is also no term with a Levi-Civita tensor carrying an external Lorentz index.)

We can simplify our calculation significantly by using the general form derived in equation (2.4) to justify disregarding many diagrams and we illustrate how this is done to handle the quark-loop diagrams.

Firstly, we note that the only diagrams capable of producing a Levi-Civita tensor are those containing a trace of a γ^5, i.e. the diagrams with quark loops. Of the diagrams with quark loops, we note that the loop integrals for diagrams (a), (b) and (c) only contain the loop momentum and one other momentum, and have as a maximum two factors of the momentum on the numerator (both of which contract with a trace of gamma matrices). This means that, even taking reparametrisation of the integrand into account, the only possible terms in the numerator contain either

- Two identical momenta contracted with a Levi-Civita tensor, which gives zero since the Levi-Civita tensor is antisymmetric, or
- One loop momentum and one other momentum contracted with a Levi-Civita tensor, which gives zero since such a term is odd in the loop momentum and the loop momentum integral is over all of space-time, or
- No loop momenta, but such a term does not yield a Levi-Civita tensor, since the trace involving a γ^5 term contains only two other gamma matrices, and this is zero.

So the only contributions to amplitude coefficients B, C and D come from $F^{(q:d)}_{\mu\nu\rho}$ and $F^{(q:e)}_{\mu\nu\rho}$.

Secondly, we note that in evaluating the contribution to amplitude coefficient A, we may sum the coefficients of either the term $\eta_{\mu\nu} p \cdot q p_\rho$ or the term $-q_\mu p_\nu p_\rho$. We choose the latter term.

We have already noted that no counterterm diagram contains more than one loop momentum carrying an external Lorentz index, so no counterterm diagram provides a contribution we need to evaluate.

In the quark loop sector, it initially appears that we can obtain a contribution we need to evaluate from each diagram, it being possible to obtain terms with three external momenta carrying external Lorentz indices in each case. However, we note that considering the loop momentum integral, in the case of diagram (a) such a term would have to contain a factor of r_ρ, which contracts with the external polarisation vector to give zero, and similarly such a term in diagram (b) would have to contain a factor of q_ν and such a term in diagram (c) would have to contain a factor of p_μ, both of which contract with external polarisation vectors to give zero. So the only contributions to amplitude coefficient A from the quark loop sector that we need to evaluate come from $F^{(q:d)}_{\mu\nu\rho}$ and $F^{(q:e)}_{\mu\nu\rho}$.

Having identified the diagrams, one can proceed with summing the diagrams and using standard trace manipulations, Feynman parametrisation of the integrals, integral redefinitions etc. one obtains the result that each quark loop contributes a total of $-\frac{2g^2 g^{(1q)} f^{abc}}{(4\pi)^2} I(m_q, M_{KK})$ to the coefficient A in equation (2.4), a total of $\mp\frac{g^2 g^{(1q)} d^{abc}}{(4\pi)^2} K(m_q, M_{KK})$ to the coefficient C and a total of $\pm\frac{g^2 g^{(1q)} d^{abc}}{(4\pi)^2} K(m_q, M_{KK})$ to the coefficient D, where the sign of the contribution varies as the quark is right- or left-handed. where

$$I(m_q, M_{KK}) = \int_0^1 dx \int_0^{1-x} dy \frac{xy(1-x-y)}{m_q^2 - xyM_{KK}^2} \tag{2.9}$$

and

$$K(m_q, M_{KK}) = \int_0^1 dx \int_0^{1-x} dy \frac{xy}{m_q^2 - xyM_{KK}^2} \,. \tag{2.10}$$

When the contributions from the KK gluon and KK ghost loop diagrams are added to that of the quark-loop diagrams, we get

$$\begin{aligned} F_{\mu\nu\rho} = & A\left(\eta_{\mu\nu}\, p\cdot q - q_\mu p_\nu\right) p_\rho + \\ & + C\left(\varepsilon_{\mu\nu\rho\gamma} p^\gamma p\cdot q - \varepsilon_{\mu\nu\rho\gamma} q^\gamma p\cdot q - \varepsilon_{\mu\rho\gamma\delta} p^\gamma q^\delta p_\nu + \varepsilon_{\nu\rho\gamma\delta} p^\gamma q^\delta q_\mu\right), \end{aligned} \tag{2.11}$$

where

$$A = \frac{17}{8}\frac{g^2 g^{(111)}}{(4\pi)^2}(f^{abc} - id^{abc}) I(M_{KK}, M_{KK}) - \sum_{q_L, q_R} \frac{2g^2 g^{(1q)} f^{abc}}{(4\pi)^2} I(m_q, M_{KK}) \,, \tag{2.12}$$

$$C = \sum_{q_R} \frac{g^2 g^{(1q)} d^{abc}}{(4\pi)^2} K(m_q, M_{KK}) - \sum_{q_L} \frac{g^2 g^{(1q)} d^{abc}}{(4\pi)^2} K(m_q, M_{KK}) \,, \tag{2.13}$$

At this stage, we note that the amplitude as calculated so far contains an anomaly in the current associated with the outgoing KK gluon (that is, the on-shell Ward identity $r^\rho F_{\mu\nu\rho} = 0$ is not satisfied). This is because we have taken the gauge $A_4 = 0$, which we do not have the freedom to do in a five-dimensional non-Abelian theory with chiral delocalised quarks if we desire anomaly cancellation [18]. We must therefore now apply a gauge transformation that leaves the four-dimensional theory anomaly-free. We note that from the perspective of our current calculation, this is a technical requirement that does not affect the final result of the on-shell calculation. However, it does have the potential to affect the result for $F_{\mu\nu\rho}$. In Ref. [13] we have calculated the anomaly contribution and find that the change of gauge required for anomaly cancellation does not alter any of the diagrams already considered, and produces only one more diagram (containing a scalar line). This additional diagram does not alter the square of the on-shell matrix element.

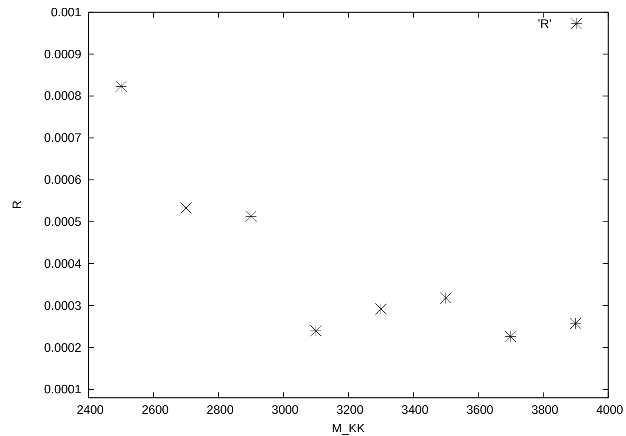

Figure 2: The ratio of the gluon-initiated NLO and the $q\bar{q}$ -initiated LO cross-sections for the production of a KK gluon at the LHC with 14TeV centre of mass energy.

3. The production cross-section

The squared matrix element works out to be

$$|\mathscr{M}|^2 = \frac{M_{KK}^6 g^4}{(4\pi)^4} \Bigg| \frac{4046}{16384} g^{(111)^2} \left[I(M_{KK}, M_{KK})\right]^2 - \\ - \frac{51}{256} g^{(111)} \sum_{q_L, q_R} g^{(1q)} [I(M_{KK}, M_{KK}) I(m_q, M_{KK})] + \\ + \frac{3}{64} \left[\sum_{q_L, q_R} g^{(1q)} I(m_q, M_{KK}) \right]^2 \Bigg|, \tag{3.1}$$

where the expressions resulting from the real and imaginary parts of the matrix element have been combined to obtain this last expression.

Writing $|\mathscr{M}|^2 = (M_{KK}^6 g^4/(4\pi)^4)|\tilde{\mathscr{M}}|^2$, we can write an expression for the cross-section as

$$\sigma = \frac{M_{KK}^2 \alpha_s^2}{8\pi} \int dy x_1 g_a(x_1, M_{KK}^2) x_2 g_b(x_2, M_{KK}^2) |\tilde{\mathscr{M}}|^2, \tag{3.2}$$

where $x_{1,2} = \sqrt{\tau} e^{\pm y}$, with $\sqrt{\tau} = M_{KK}/\sqrt{s}$, y being the rapidity of the KK gluon and $\sqrt{s}$ being the total centre of mass energy of the pp system.

We have used this expression to calculate the cross-section for the KK gluon from the gg-initial state and compared it with the leading order $q\bar{q}$ result (using the LO cross-section presented in Ref. [12]) at the Large Hadron Collider (LHC), assuming a centre-of-mass energy of 14 TeV. The ratio is plotted for some typical values of the KK gluon mass in Fig. 2. The cross-section from the gg NLO subprocesses turns out to be less than a thousandth of the LO cross section. This is, in turn, due to appearance of the large KK gluon mass squared in the denominators of the integral I.

In principle, to complete the full calculation of the KK gluon cross-section at NLO one needs to calculate the $q\bar{q}$-initiated diagrams at NLO. But given that the gg-initiated contribution is tiny, it is expected that the $q\bar{q}$-initiated contribution will be even smaller due to the suppressed couplings of valence quarks and the calculation is, therefore, not of much interest.

References

[1] L. Randall and R. Sundrum, *Phys. Rev. Lett.* **83** (1999) 3370, [hep-ph/9905221].

[2] W. D. Goldberger and M. B. Wise, *Phys. Rev. Lett.* **83** (1999) 4922–4925, [hep-ph/9907447]; W. D. Goldberger and M. B. Wise, *Phys. Lett.* **B475** (2000) 275–279, [hep-ph/9911457].

[3] H. Davoudiasl, J. L. Hewett, and T. G. Rizzo, *Phys. Rev.* **D63** (2001) 075004, [hep-ph/0006041]; K. Sridhar, *JHEP* **05** (2001) 066, [hep-ph/0103055]; B. C. Allanach, K. Odagiri, M. A. Parker, and B. R. Webber, *J. High Energy Phys.* **09** (2000) 019, [hep-ph/0006114]; B. C. Allanach, K. Odagiri, M. J. Palmer, M. A. Parker, A. Sabetfakhiri, and B. R. Webber, *J. High Energy Phys.* **12** (2002) 039, [hep-ph/0211205].

[4] J. M. Maldacena, *Adv. Theor. Math. Phys.* **2** (1998) 231–252, [hep-th/9711200].

[5] N. Arkani-Hamed, M. Porrati, and L. Randall, *JHEP* **08** (2001) 017, [hep-th/0012148]; R. Rattazzi and A. Zaffaroni, *JHEP* **04** (2001) 021, [hep-th/0012248].

[6] A. Pomarol, *Phys. Lett.* **B486** (2000) 153, [hep-ph/9911294].

[7] T. Gherghetta and A. Pomarol, *Nucl. Phys.* **B586** (2000) 141–162, [hep-ph/0003129].

[8] K. Agashe, A. Delgado, M. J. May, and R. Sundrum, *J. High Energy Phys.* **08** (2003) 050, [hep-ph/0308036]; K. Agashe, A. Delgado, and R. Sundrum, *Nucl. Phys.* **B643** (2002) 172–186, [hep-ph/0206099]; R. Contino and A. Pomarol, *JHEP* **11** (2004) 058, [hep-th/0406257]; K. Agashe, R. Contino, and A. Pomarol, *Nucl. Phys.* **B719** (2005) 165–187, [hep-ph/0412089].

[9] R. Contino, T. Kramer, M. Son, and R. Sundrum, *JHEP* **05** (2007) 074, [hep-ph/0612180].

[10] B. Lillie, L. Randall, and L.-T. Wang, *J. High Energy Phys.* **09** (2007) 074, [hep-ph/0701166].

[11] K. Agashe, A. Belyaev, T. Krupovnickas, G. Perez, and J. Virzi, *Phys. Rev.* **D77** (2008) 015003, [hep-ph/0612015].

[12] M. Guchait, F. Mahmoudi, and K. Sridhar, *J. High Energy Phys.* **05** (2007) 103, [hep-ph/0703060].

[13] B. C. Allanach, F. Mahmoudi, J. P. Skittrall and K. Sridhar, *J. High Energy Phys.* **1003** (2010) 014 [arXiv:0910.1350 [hep-ph]].

[14] M. Guchait, F. Mahmoudi, and K. Sridhar, *Phys. Lett.* **B666** (2008) 347–351, [arXiv:0710.2234].

[15] A. Djouadi, G. Moreau, and R. K. Singh, *Nucl. Phys.* **B797** (2008) 1–26, [arXiv:0706.4191].

[16] J. F. Nieves and P. B. Pal, *Phys. Rev.* **D72** (2005) 093006, [hep-ph/0509321].

[17] B. C. Allanach, J. P. Skittrall, and K. Sridhar, *J. High Energy Phys.* **11** (2007) 089, [arXiv:0705.1953].

[18] P. Anastasopoulos, M. Bianchi, E. Dudas, and E. Kiritsis, *J. High Energy Phys.* **11** (2006) 057, [hep-th/0605225].