Pramana – J. Phys. (2017) 89:58
DOI 10.1007/s12043-017-1459-z

Exploring the warped bulk

K SRIDHAR

Department of Theoretical Physics, Tata Institute of Fundamental Research, Homi Bhabha Road, Colaba, Mumbai 400 005, India
E-mail: sridhar@theory.tifr.res.in

Published online 5 October 2017

Abstract. A review of the class of models that go under the name bulk Randall–Sundrum models is presented here. The issue of localization of quantum fields in the five-dimensional bulk and the profiles of the zero modes and the Kaluza–Klein excitations are discussed. The zero modes of these bulk fields are, in general, partially composite. The degree of compositeness of the different fields is discussed and this provides the basis for realizing a Standard Model in the bulk, albeit partially composite. The viability of this model and its extensions when confronted with electroweak precision measurements is also discussed. Two such extensions are: (1) models with a bulk custodial symmetry and (2) models with a deformed metric. The signatures of these models that we expect at collider experiments are discussed and also the search for the Kaluza–Klein excitation of the gluon as the most important of these signatures.

Keywords. Beyond Standard Model Physics; extra dimensions; Randall–Sundrum model; collider searches.

PACS Nos **73.21.Hb; 73.21.La; 73.50.Bk**

1. Introduction

The Randall–Sundrum model (RS model) [1] is an elegant approach to address the gauge-hierarchy problem using the idea of brane-world inspired TeV scale extra dimensions. The RS model, as originally formulated, is a five-dimensional model where the fifth dimension y is compactified on an S^1/Z^2 orbifold of radius, R. At the orbifold fixed points $y = 0$ and $y = \pi R \equiv L$ two branes, the UV and the IR brane respectively, are located.

The novelty of the model is that it uses a warped metric:

$$ds^2 = \mathrm{e}^{-2A(y)}\eta_{\mu\nu}\mathrm{d}x^\mu \mathrm{d}x^\nu - \mathrm{d}y^2 \tag{1}$$

and with this metric one can show that the solutions to the Einstein equation imply

$$A(y) = \pm k|y|, \tag{2}$$

where $k^2 \equiv -(\Lambda/12M^3)$ with M being the Planck scale.

The RS model suggests a way out of the gauge-hierarchy problem by localizing the graviton field in the bulk and very close to the UV brane whereas the Standard Model (SM) fields are all IR-localized. There is only one scale in the problem: the Planck scale but an electroweak scale at around 250 GeV materializes because of the exponential suppression provided by the warp factor. The hierarchy $v/M \sim 10^{-16}$ (where v is the vacuum expectation value of the SM Higgs field) can be obtained by choosing the exponent of the warp factor, kL, to be of the order of 30 or so. While the Higgs vev is warped down to 250 GeV, the UV localization of the gravitons ensures that the Planck scale does not get affected by the warping. One gets, in fact,

$$M^2 = \frac{M^3}{k}[1 - \mathrm{e}^{-2kL}]. \tag{3}$$

The problem, however, is that the suppression that one obtains for the Higgs vev is effective for all the fields localized on the IR brane. Thus, mass scales which suppress dangerous higher-dimensional operators responsible for proton decay or neutrino masses also become small which spells a disaster for the RS model. One way out of this is to realize that to solve the gauge-hierarchy problem one needs to only localize the Higgs on the IR brane but all the other SM fields could be in the bulk [2–4]. In fact, even the Higgs need not be sharply localized on the IR brane but only somewhere close to it. In this way, it is possible to make viable variations of the RS model, now collectively known as bulk RS models. These models yield a bonus: localizing fermions at different positions in the bulk gives different overlaps of their profiles with the Higgs field, which is localized on or close to the IR brane. This gives rise in a natural way to the Yukawa-coupling hierarchy [3].

To discuss the localization of fields in the bulk, we take the example of a scalar field. We start with the action for a 5-d scalar field with a bulk mass ak and a boundary mass term bk and vary the action to get the equations of motion. To separate the five-dimensional part, we use separation of variables in the Kaluza–Klein (KK) decomposition to write

$$\phi(x^\mu, y) = \sum_{n=0}^{\infty} \phi^{(n)}(x^\mu) f_\phi^{(n)}(y). \quad (4)$$

With this in hand, we can solve for the profiles $f_\phi^{(n)}(y)$. For the zero mode we get a profile $f_\phi^{(0)}(y) = C e^{(b-1)ky}$ only if $b = 2 - \sqrt{4+a}$. For $b < 1$ the zero mode is localized towards the UV brane and for $b > 1$, towards the IR brane. The scalar field zero mode can be localized anywhere in the bulk whereas the KK modes, which are given by combinations of Bessel functions, are localized towards the IR brane and have masses of the order of the infrared scale $ke^{-\pi kR}$.

For bulk fermions, similarly, we get for the zero mode $f_L^{(0)}(y) \sim e^{(\frac{1}{2}-c)ky}$, where c is the bulk mass parameter. When $c > \frac{1}{2}$ the zero mode is towards UV brane, $c < \frac{1}{2}$ towards IR brane, i.e. like in the scalar case, the zero mode of the bulk fermion can be localized anywhere in the bulk. For bulk gauge fields, the zero mode is flat in the bulk, while the graviton zero mode is UV localized. All the massive KK modes given by Bessel functions are IR localized. The profiles for the different fields are shown in figure 1.

The AdS/CFT correspondence, when worked out for a slice of AdS space–time, also provides both a motivation for and an understanding of bulk models. (For a review, see ref [5].) In the full AdS space–time, the AdS/CFT correspondence provides an exact equivalence between a Type IIB superstring theory on $AdS_5 \times S_5$ and a superconformal $\mathcal{N} = 4$ $U(N)$ gauge theory. This correspondence implies the existence of a precise map between the correlation functions in the two theories. To go to the weak coupling limit of the string theory so that in the limit we approach a gravity theory, we need to work in the strong coupling regime of the gauge theory. A set of fields $\phi(x^\mu, y)$ in the bulk acquire a value on the boundary $\phi(x^\mu)$, a 4-d field which is a Schwinger source field for operators $\mathcal{O}$ of the conformal field theory. It is then possible to study n-point functions of the strongly coupled conformal field theory in terms of the five-dimensional effective action.

We need to, however, discuss the AdS/CFT correspondence not in the full AdS space–time but only in a segment of it – the segment enclosed between the two branes in the RS model. The UV brane in five dimensions shows up as a UV cut-off in the dual theory and the source field ϕ_0 becomes dynamical. The IR brane leads to a spontaneous breakdown of conformal symmetry which gives rise to CFT bound states. The mixing

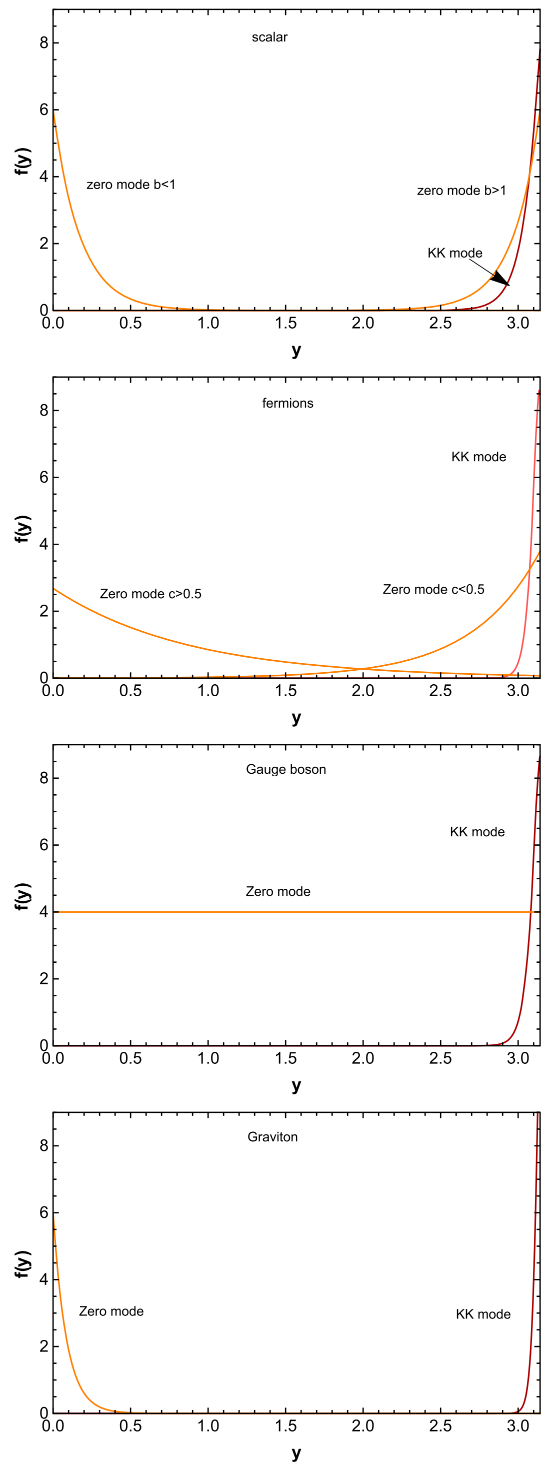

Figure 1. Profiles for the zero modes and KK modes of scalar, fermion, gauge boson and graviton fields in the bulk.

$\phi_0\mathcal{O}$ between the source field (elementary) and the CFT bound states (composites) gives rise to the mass eigenstates which are from a 5-d perspective KK excitations of the bulk fields.

Because of the mixing alluded to above, the zero modes and the KK modes of the bulk fields are partially composite objects. To study this partial compositeness, instead of doing a KK expansion, it is more useful to expand the bulk field in terms of the source field and CFT composites. Such an expansion is referred to as a holographic expansion. In general, there will be both kinetic and mass mixing but after diagonalization and matching with the KK decomposition, one can identify the degree of compositeness of any bulk field. The graviton zero mode which is UV localized is pure elementary and has a tiny mixing with the composite states but the KK modes of the graviton are purely composite. The gauge field zero mode has a flat profile. So mixing is expected to be large. But the coupling is marginal and the zero modes are elementary. Again the KK modes are purely composite.

Light fermions like electrons ($c > 1/2$) are UV localized and so they are elementary. For $(t,\ b)_L$, c is 0.3–0.4. For these states, zero modes contain significant fraction of the composite CFT states though it is largely elementary. On the other hand, the right-handed top ($c < -1/2$) has a large composite component and the Higgs which is IR localized has also a purely composite zero mode. The KK modes of the fermions are all composite, localized as they are close to the IR brane.

2. Precision electroweak constraints

It is then possible to construct a bulk extension of the SM by having the gauge and fermion fields in the bulk, the Higgs localized on or near the IR brane and with a suitable mechanism to make the model successfully confront constraints from electroweak precision measurements [5]. Such a model has interesting features. Other than providing a framework for addressing the question of fermion mass hierarchy, it also naturally results in small mixing angles in the Cabibbo–Kobayashi–Maskawa (CKM) matrix, gauge-coupling universality and suppression of flavour-changing neutral currents to experimentally acceptable values [6–10].

Electroweak precision tests provide very strong constraints on bulk models. If, for example, only gauge bosons propagate in the bulk but the fermions are localized on the IR brane, then the couplings of the gauge boson Kaluza–Klein (KK) modes to the IR-localized fermions yield unacceptably large contributions to T and S and this yields a lower bound of 25 TeV on the mass of the first KK mode of the gauge boson. Of course,

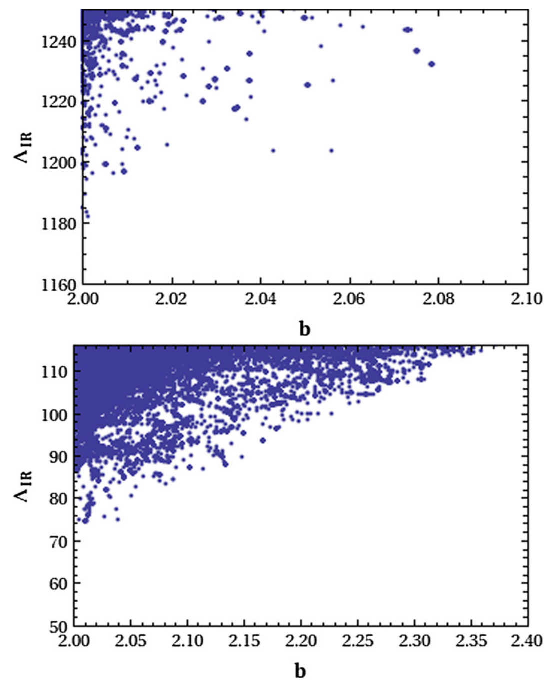

Figure 2. Fits to electroweak precision measurements for the custodial symmetry model and the deformed metric model.

one way to relax this bound is to localize the fermions in the bulk and especially the light fermions close to the UV brane and this significantly reduces the constraint coming from the S-parameter. But the T parameter constraints require further attention. One way to handle this [11,12] is by enlarging the gauge symmetry in the bulk to $SU(3)_\mathrm{c} \times SU(2)_\mathrm{L} \times SU(2)_\mathrm{R} \times U(1)_\mathrm{y}$ which is an enlarged custodial symmetry which is broken on the IR brane to recover the SM gauge group. It turns out that the corrections to the T parameter coming from the dangerous KK gauge boson sector can be tamed using this custodial symmetry and by a judicious choice of the fermion representations under the extended gauge group it is also possible to rein in the non-oblique $Z \to b\bar{b}$ corrections and eventually the bound on the lightest KK gauge boson mode comes down to about 3 TeV [13,14] (see figure 2).

An alternate proposal to address the issue of the T parameter [15,16] uses a deformed metric near the IR brane along with moving the Higgs scalar into the bulk. For this set up, the function $A(y)$ in eq. (2) is then modified to

$$A(y) = ky - \frac{1}{\nu^2}\log\left(1 - \frac{y}{y_s}\right). \quad (5)$$

The UV brane, similar to the RS set-up, is located at $y = 0$. The IR brane is however located at $y = y_1$ with the position of the singularity ($y = y_s$) located behind the IR brane at $y_s = y_1 + \Delta$, where $\Delta \sim 1/k$. y_1 is determined by demanding the solution to the hierarchy problem which requires $A(y_1) \sim 35$. The limit $\nu \to 0$ reverts to the original RS geometry in eq. (2). The deformation of the metric actually causes the Higgs field to be moved further away from the IR brane but the gauge boson KK modes are moved by the deformation towards the IR brane. This differential action causes the overlap of the Higgs and KK gauge boson modes to be reduced and that relaxes the bounds coming from the T parameter and the mass of the first KK gauge boson mode in this model can be as small as 1.5 TeV [14] (see figure 2).

3. Kaluza–Klein gluons and collider searches

In typical bulk RS models, the lightest KK excitations are those of the gauge bosons and searches for these are likely to be the most promising probes of such a model. Of these, the KK gluons, because of their larger couplings, are the most interesting though there are interesting signals from KK excitations of electroweak gauge bosons and fermions.

In the custodial models, the couplings of g_{KK} to the SM states [17] are parametrized in terms of the parameter $\xi \equiv \sqrt{kL} \sim 5$ and relative to the QCD coupling g_s are given as

$$g^{q\bar{q}g_{\mathrm{KK}}} \approx \frac{1}{\xi} g_s, \qquad g^{Q\bar{Q}^3 g_{\mathrm{KK}}} \approx 1 g_s,$$
$$g^{t_R \bar{t}_R g_{\mathrm{KK}}} \approx \xi g_s, \qquad g^{g g g_{\mathrm{KK}}} = 0. \tag{6}$$

These denote the coupling of the first Kaluza–Klein mode of the gluon to light quarks, to the third-generation left-handed doublet, to the right-handed top quark and to the gluon, respectively. Note that g_{KK} couples predominantly to the right-handed top quark and, consequently, the $g_{\mathrm{KK}} \to t\bar{t}$ branching ratio is more than 90%. For the deformed metric, the couplings are similar except for overall factors.

Also, one sees that the coupling of g_{KK} to the zero-mode Standard Model gluons vanishes because of the orthonormality of the Kaluza–Klein modes. This means that, at leading order, g_{KK} production takes place via annihilation of light quarks, to which the coupling of g_{KK} is suppressed and, consequently, the cross-section is small especially since electroweak precision constraints require the mass of g_{KK} to be not less than 2–3 TeV. The produced g_{KK} decays into a $t\bar{t}$ pair and this tiny cross-section has to compete with a huge QCD $t\bar{t}$ production background. The fact that this is a resonant cross-section helps somewhat but then g_{KK} has a very large width, owing to its strong coupling to the tops, and so the resonant bump is not sharp but rather smeared out. The fact that g_{KK} couples chirally to the tops is, however, an advantage and a forward–backward asymmetry to pick out the signal may be used. However, in a pp machine like LHC, this is not easy. Finally, the fact that t and $\bar{t}$ come from the decay of a heavy object not less than 2–3 TeV in mass are highly boosted. These boosted top jets can effectively be used as signal discriminant [17,18].

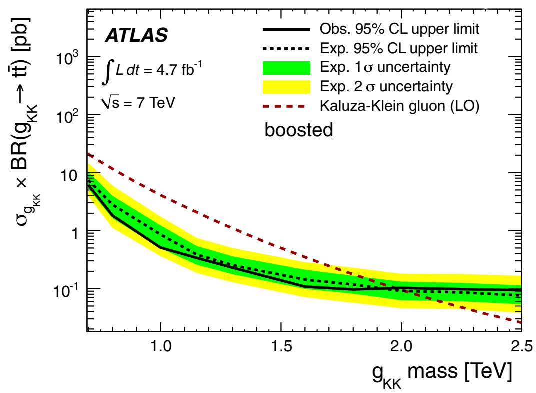

Figure 3. Bound on KK gluon mass from the ATLAS experiment.

Nonetheless, given the smallness of the cross-section, it becomes important to look for other production mechanisms for g_{KK}, especially ones which have gluon initial states. The associated production of g_{KK} with a $t\bar{t}$ pair leading to distinctive four top final states (with two boosted tops) is a process that has been studied with this in mind [19]. The production of g_{KK} through top loops has also been considered [20] though the loop contributions are greatly suppressed.

The ATLAS and CMS experiments at the LHC have searched for a KK gluon and have put constraints on its mass from the non-observation of any signal in the studies undertaken so far. The present bounds on the mass are about 2 TeV in mass as can be seen from a typical search plot presented by the ATLAS experiment in figure 3.

In the following we discuss a new mechanism to search for KK gluons which is the production of a g_{KK} with an associated jet (a light quark or gluon jet). As this appears at an order of α_s more than the leading-order g_{KK} production, one would think that the cross-section is smaller. But this naive expectation is not true because the associated jet production process has both $q\bar{q}$ and qg initial states and a larger number of subprocess contributions. Consequently, the cross-section for this process is comparable to the leading-order g_{KK} production process.

g_{KK} will decay predominantly to tops and these, in turn, will decay to a b and a W which will finally result in a multijet final state with the associated jet being one of these jets. It is useful to note some features of the kinematics here. We are interested in producing g_{KK}, at

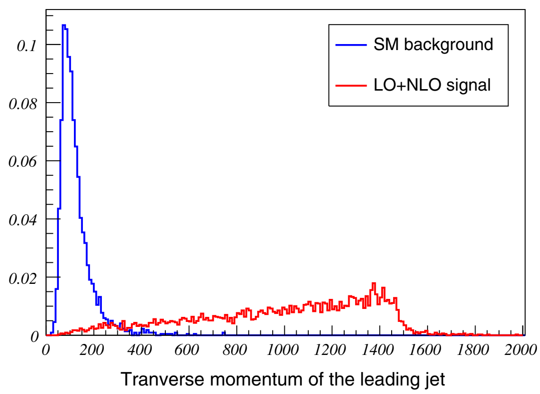

Figure 4. p_T distribution for the leading jet for the signal (red) and the background (blue) for $m_{g_{KK}} = 3$ TeV.

least 2.5 TeV or so in mass, in association with a jet. Even at the highest energies at the LHC now available, the subprocess centre-of-mass energy will be sufficient only to produce g_{KK} with small momentum with the p_T-balancing associated jet also, therefore, possessing a p_T that is not very large. When g_{KK} decays into the $t\bar{t}$ pair, they are produced almost back-to-back. But the $t\bar{t}$ pair so produced will be highly boosted and give rise to very collimated decay products.

The parton-level amplitudes for both the signal and the background were generated using MADGRAPH [21] at 13 TeV centre of mass energy. This output was then matched to PYTHIA 8 [22] using MLM matching scheme. Jets are reconstructed from these partons by employing FASTJET [23,24] using the anti-k_T [25] clustering algorithm.

Top plot in figure 4 displays the minimum luminosity required for a 5σ signal sensitivity for different masses for both normal RS and deformed RS models. Owing to constraints from precision electroweak data we do not consider masses below 2.5 TeV for normal RS model. Due to their larger production cross-section, the lower masses (indicated by blue points in the figure) have better sensitivity in terms of early discovery prospects. We find that even masses as heavy as 4 TeV can be accessed with an integrated luminosity of $\sim$800fb^{-1}. This is particularly useful, as the constraints from indirect measurements become tighter, pushing the masses to higher values.

Lower masses can be admitted in an RS model with a deformity near the IR brane. For the deformed model we choose $\nu = 0.4$ while the curvature radius is chosen to be $L_1 = 0.2/k$ [26]. This scenario however suffers from reduced production cross-section owing to the smaller coupling of g_{KK} to the lighter quarks.

The reach for the lightest point in deformed RS model (indicated by majenta points in figure 5) is slightly peculiar. A 1.5 TeV KK gluon requires a fairly high

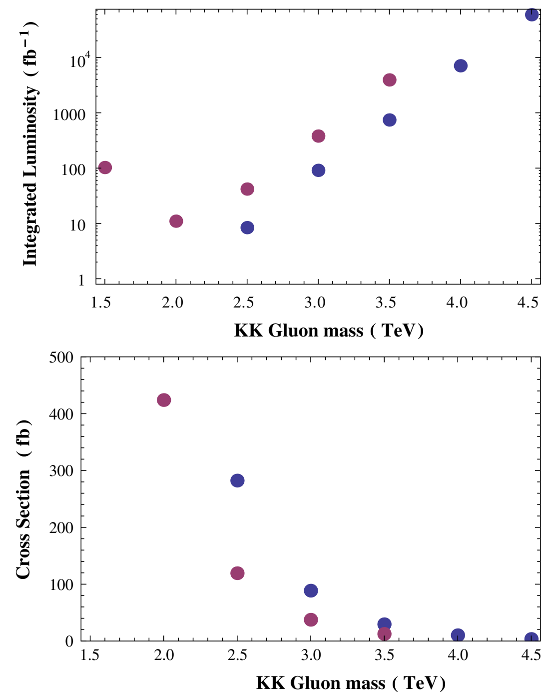

Figure 5. Minimum luminosity required for a 5σ sensitivity for normal RS (blue) and deformed RS (majenta). The bottom plot shows the production cross-section for different masses.

luminosity for its discovery as compared to masses as heavy as 2.5 TeV. This can be attributed to a very hard cut on the transverse momentum of the leading jet, $p_T^{j_0} > 900$. Since p_T of the leading jet is $\sim m_{g_{KK}}/2$, this cut is more effective for the heavier masses as compared to the lighter masses. While this depletes majority of the signal points for 1.5 TeV KK gluon, this is helpful in depleting the $t\bar{t}$ background to a great extent. This is evident in the promising reaches for heavier KK states in the near future run of the LHC.

References

[1] L Randall and R Sundrum, *Phys. Rev. Lett.* **83**, 3370 (1999)
[2] A Pomarol, *Phys. Lett. B* **486**, 153 (2000)
[3] T Gherghetta and A Pomarol, *Nucl. Phys. B* **586**, 141 (2000)
[4] Y Grossman and M Neubert, *Phys. Lett. B* **474**, 361 (2000)
[5] T Gherghetta, TASI Lectures on a Holographic View of Beyond the Standard Model Physics. Physics of the Large and the Small, *Proceedings of the Theoretical Advanced Study Institute in Elementary Particle Physics* edited by C Csaki and S Dodelson (2010)
[6] G Burdman, *Phys. Lett. B* **590**, 94 (2004)
[7] S J Huber, *Nucl. Phys. B* **666**, 269 (2003)

[8] S Casagrande, F Goertz, U Haisch, M Neubert and T Pfoh, *J. High Energy Phys.* **0810**, 094 (2008)
[9] M Bauer, S Casagrande, U Haisch and M Neubert, *J. High Energy Phys.* **1009**, 017 (2010)
[10] K Agashe, G Perez and A Soni, *Phys. Rev. D* **71**, 016002 (2005)
[11] K Agashe, A Delgado, M J May and R Sundrum, *J. High Energy Phys.* **0308**, 050 (2003)
[12] K Agashe, R Contino, L Da Rold and A Pomarol, *Phys. Lett. B* **641**, 62 (2006)
[13] H Davoudias, S Gopalakrishna, E Ponton and J Santiago, *New J. Phys.* **12**, 075011 (2010)
[14] A M Iyer, K Sridhar, S K Vempati, *Phys. Rev. D* **93**(**7**), 075008 (2016)
[15] J A Cabrer, G von Gersdorff and M Quiros, *Phys. Lett. B* **697**, 208 (2011)
[16] J A Cabrer, G von Gersdorff and M Quiros, *J. High Energy Phys.* **05**, 083 (2011)
[17] K Agashe, A Belyaev, T Krupovnickas, G Perez and J Virzi, *Phys. Rev. D* **77**, 015003 (2008)
[18] B Lillie, L Randall and L-T Wang, *J. High Energy Phys.* **0709**, 074 (2007)
[19] M Guchait, F Mahmoudi and K Sridhar, *Phys. Lett. B* **666**, 347 (2008)
[20] B C Allanach, F Mahmoudi, J P Skittrall and K Sridhar, *J. High Energy Phys.* **1003**, 014 (2010)
[21] J Alwall, M Herquet, F Maltoni, O Mattelaer and T Stelzer, *J. High Energy Phys.* **06**, 128 (2011)
[22] T Sjostrand, S Mrenna and P Z Skands, *Comput. Phys. Commun.* **178**, 852 (2008)
[23] M Cacciari, FastJet: A Code for fast k_t clustering, and more, in: *Proceedings of the 14th International Workshop on Deep Inelastic Scattering* (Tsukuba, Japan, April 20–24, 2006) pp. 487–490
[24] M Cacciari, G P Salam and G Soyez, *Eur. Phys. J. C* **72**, 1896 (2012)
[25] M Cacciari, G P Salam and G Soyez, *J. High Energy Phys.* **04**, 063 (2008)
[26] A Carmona, E Ponton and J Santiago, *J. High Energy Phys.* **10**, 137 (2011)