Associated production of a Kaluza-Klein excitation of a gluon with a $t\bar{t}$ pair at the LHC

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ABSTRACT

In Randall-Sundrum models, the Kaluza-Klein (KK) excitations of the gluon, $g_{KK}$ have enhanced couplings to the right-handed quarks. In the absence of a $ggg_{KK}$ coupling in these models, the single production of a $g_{KK}$ from an initial $gg$ state is not possible. The search for other production mechanisms at the LHC, therefore, becomes important. We suggest that the associated production of a $g_{KK}$ with a $t\bar{t}$ pair is such a mechanism. Our study shows that through this process the LHC can probe KK gluon masses in the range of 2.8 – 2.9 TeV.

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The Randall-Sundrum (RS) model [1] is a five-dimensional model with the fifth dimension \( \phi \) compactified on a \( S^1/Z^2 \) orbifold. The compactification radius \( R_c \) is somewhat larger than \( M_P^{-1} \), the Planck length. The fifth dimension is a slice of anti-de Sitter spacetime and is strongly curved. At the fixed points \( \phi = 0, \pi \) of the orbifold, two D-3 branes are located and are known as the Planck brane and the TeV brane, respectively. The Standard Model fields are localised on the TeV brane while gravitons exist in the full five-dimensional spacetime. The five-dimensional spacetime metric is of the form

\[
ds^2 = e^{-\mathcal{K} R_c \phi} \eta_{\mu\nu} dx^\mu dx^\nu + R_c^2 d\phi^2.
\]

Here \( \mathcal{K} \) is a mass scale related to the curvature. The factor \( \exp(-\mathcal{K} R_c \phi) \) is called the warp factor and serves to rescale masses of fields localised on the TeV-brane. For example, \( M_P = 10^{19} \) GeV for the Planck brane at \( \phi = 0 \) gets rescaled to \( M_P \exp(-\mathcal{K} R_c \pi) \) for the TeV brane at \( \phi = \pi \). The warp factor generates \( \frac{M_P}{M_{EW}} \sim 10^{15} \) by an exponent of order 30 and solves the hierarchy problem. For this mechanism to work, one will have to ensure that the radius \( R_c \) is stabilised against quantum fluctuations and this can be done by introducing a bulk scalar field which generates a potential that allows for the stabilisation [2]. The model predicts a discrete spectrum of Kaluza-Klein (KK) excitations of the graviton and these couple to the Standard Model fields with a coupling that is enhanced by the warp factor to be of the order of electroweak strength. Several collider implications of these gravitons resonances have been studied in the literature [3].

The AdS/CFT correspondence [4] allows us to get an understanding of the RS model in terms of a dual theory – a strongly coupled gauge theory in four dimensions [5]. This four-dimensional theory is conformal all the way from the Planck scale down to the TeV scale and it is only the presence of the TeV brane that breaks the conformal symmetry. The KK excitations as well as the fields localised on the TeV brane are TeV-scale composites of the strongly interacting theory. Since in the RS model, all SM fields are localised on the TeV brane, the AdS/CFT correspondence tells us that the RS model is dual to a theory of TeV-scale compositeness of the entire SM. Such a composite theory is clearly unviable: but is there a way out? There seems to be – and 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].

Flavour hierarchy, consistency with electroweak precision tests and avoidance of flavour-changing neutral currents can be used as guiding principles in constructing such models [7]. In particular, in order to avoid an unacceptably large contribution to the electroweak $T$ parameter an enhanced symmetry in the bulk like $SU(2)_L \times SU(2)_R \times U(1)_{(B-L)}$ may be used. The heavier fermions need to be closer to the TeV brane so as to get a large Yukawa coupling i.e. overlap with the Higgs. In other words, the profiles of the heavier fermions need to be peaked closer to the TeV-brane. Conversely, the fermions close to the Planck brane will have small Yukawa couplings. However, while the large Yukawa of the top demands proximity to the TeV brane, the left-handed electroweak doublet, $(t, b)_L$, cannot be close to the TeV brane because that induces non-universal couplings of the $b_L$ to the $Z$ constrained by $Z \to b\bar{b}$. So the doublet needs to be as far away from the TeV brane as allowed by $R_b$ whereas the $t_R$ needs to be localised close to the TeV brane to account for the large Yukawa of the top. We stress that this is one model realisation; a different profile results, for example, in models that invoke other symmetry groups in the bulk [8]. It has been found that in order to avoid huge effects of flavour-changing neutral currents (FCNCs) and to be consistent with precision tests of the electroweak sector, the masses of the KK modes of the gauge bosons have to be strongly constrained. The resulting bounds on the masses of the KK gauge bosons are found to be in the region of 2-3 TeV [7] though this bound can be relaxed by enforcing additional symmetries. A review of the literature on this subject can be found in Ref. [9].

The collider implications of this scenario has been studied recently [10]. While some of these studies have focussed on graviton production [11], the interesting signals for this scenario is the production of KK gauge bosons and, for the LHC in particular, the production of KK gluons. The KK gluon couples strongly to the $t_R$, with a strength which is enhanced by a factor $\xi$ compared to the QCD coupling where $\xi \equiv \sqrt{\log(M_{pl}/\text{TeV})} \sim 5$. Consequently, it decays predominantly to tops if produced. To the left-handed third-generation quarks, the KK gluon couples with the same strength as the QCD coupling whereas to the light quarks its couplings are
suppressed by a factor $1/\xi$. The problem in producing the KK gluon at a collider, however, is that its coupling to two gluons vanishes because of the orthogonality of the profiles of these particles and, therefore, the gluon production mechanism at a hadron collider cannot produce the KK gluon at leading order. The KK gluon can, therefore, be produced by annihilation of light quarks and this production mechanism has been studied in the context of the LHC [12]. The same mechanism has also been studied in the context of Tevatron to derive a model-independent bound of 770 GeV from the Tevatron top cross-section [13].

In this paper, we study the production of KK gluons in association with a $t\bar{t}$ pair. A similar process has been recently discussed in Ref. [14]. In this process the $t\bar{t}$ pair can be produced from both $gg$ and $q\bar{q}$ initial states through the usual QCD processes and the KK gluon, $g_{KK}$ can then be radiated from one of the heavy-quark legs. The fact that the $gg$ initial state contributes to the associated production process makes it appealing. Also the process directly probes the coupling of the $g_{KK}$ to the tops which is an important feature of the new dynamics.

The Feynman diagrams for the $q\bar{q} \rightarrow g_{KK}t\bar{t}$ and the $gg \rightarrow g_{KK}t\bar{t}$ subprocess are shown in Fig. 1. We have computed the matrix elements for these subprocesses using FORM. The $g_{KK}$ is produced on-shell and we ignore virtual effects. The produced $g_{KK}$ decays into a $t\bar{t}$ pair yielding two pairs of $t\bar{t}$ in the final state. The background to this signal of two non-resonant $t\bar{t}$ pairs coming from QCD processes has been computed using ALPGEN [15]. The squared-matrix elements for the signal are available in the form of a Fortran code but the expressions are too lengthy to reproduce here. A KK gluon with a mass just a little above the $t\bar{t}$ threshold has a very large branching into top pairs: the branching ratio is about 92.5% [12]. Since we are interested in KK gluon masses well above the $t\bar{t}$ threshold we will assume that the produced $g_{KK}$ decays with this branching ratio into a $t\bar{t}$ pair.

For the signal kinematics, we have used that originally proposed by Gottschalk and Sivers [16] for three-jet production suitably modified to take into account the fact that all three final-state particles in our case are massive. In this description of the kinematics, the $z$-axis of the co-ordinate system is chosen to be the direction of one of the final-state particles rather than the initial beam axis. We choose this
Figure 1: The Feynman diagrams for the processes: (a) $q\bar{q} \rightarrow g_{KK}t_R\bar{t}_R$ and, (b) $gg \rightarrow g_{KK}t_R\bar{t}_R$.

particle (labelled $p_5$) as the $g_{KK}$. The momentum assignments that we start with are:

\[
\begin{align*}
 p_1 & : \frac{\sqrt{s}}{2} (1, \sin \theta \cos \phi, \sin \theta \sin \phi, \cos \theta) \\
 p_2 & : \frac{\sqrt{s}}{2} (1, -\sin \theta \cos \phi, -\sin \theta \sin \phi, \cos \theta) \\
 p_3 & : \frac{\sqrt{s}}{2} x_3 (1, \beta_3 \cos \theta_{35}, \beta_3 \sin \theta_{35}, 0) \\
 p_4 & : \frac{\sqrt{s}}{2} x_4 (1, \beta_4 \cos \theta_{45}, \beta_4 \sin \theta_{45}, 0) \\
 p_5 & : \frac{\sqrt{s}}{2} x_5 (1, \beta_5, 0, 0) .
\end{align*}
\]

In the above equation, $p_1$ and $p_2$ are the 4-momenta of the initial partons, $p_3$ and $p_4$ are the 4-momenta of the $t$ and the $\bar{t}$ and $p_5$ is the 4-momentum of the KK gluon.
The $\beta_i$'s are given by:

$$\beta_i = \sqrt{1 - \frac{4m_i^2}{x_i^2 s}}$$

(3)

Energy conservation implies

$$x_3 + x_4 + x_5 = 2,$$

(4)

and, using 3-momentum conservation, one can get

$$\cos\theta_{35} = \frac{x_3^2\beta_3^2 - x_4^2\beta_4^2 - x_5^2\beta_5^2}{2x_3x_5\beta_3\beta_5}$$

$$\cos\theta_{45} = \frac{x_3^2\beta_3^2 - x_4^2\beta_4^2 - x_5^2\beta_5^2}{2x_4x_5\beta_4\beta_5}$$

(5)

Using the above, all relevant momenta and angles may be constructed. For example, the transverse momentum of the $g_{KK}$ is given by

$$p_T(g_{KK}) = \frac{\sqrt{s}}{2}x_5\beta_5\sqrt{\cos^2\theta + \sin^2\theta\sin^2\phi}$$

(6)

The kinematics for the decay of the $g_{KK}$ into a $t\bar{t}$ pair is the standard two-particle decay kinematics.

Defining the variables $\tau$ and $y_{boost}$ through the equations:

$$\tau = \frac{\hat{s}}{s} = x_1x_2$$

$$x_1 = \sqrt{\tau} e^{y_{boost}}$$

$$x_2 = \sqrt{\tau} e^{-y_{boost}}$$

(7)

the differential cross-section for $g_{KK}t\bar{t}$ production assumes the form:

$$\frac{d\sigma}{d\sqrt{s}dy_{boost}dx_3dx_4d \Omega} = \int dx_1dx_2 \frac{\alpha_s^2 \Lambda_t^2 \tau}{8\pi \sqrt{s}} \sum_{ij} \frac{1}{1 + \delta_{ij}} \left[ f_i^{(a)}(x_1, Q^2) f_j^{(b)}(x_2, Q^2) |M_{ij}|^2 \right]$$

(8)

where, $f^{(a)}$, $f^{(b)}$ are the parton densities evaluated at the scale $Q^2$, $|M_{ij}|^2$ is the squared-matrix element and $\Lambda_t$ is the coupling of the KK gluon to the $t_R$ and is given by $5\sqrt{4\pi\alpha_s}$.

Since the $g_{KK}$ masses we are interested in are large, we expect the $t\bar{t}$ pair coming from the decay of the $g_{KK}$ to have large momenta. The other $t\bar{t}$ pair is expected to
have more moderate values of momenta. This simple fact may allow one to enhance the quality of the $g_{KK}$ signal over the QCD background. As a first guess, we choose to put a lower cut of 300 GeV on the $p_T$ of the $t$ and the $\bar{t}$ coming from the decay of the $g_{KK}$ and a cut of 50 GeV on each of the other pair. We use these cuts to calculate the cross-section for the associated production of the KK gluon with a $t\bar{t}$ pair.

![Figure 2](image.png)

Figure 2: The cross-section for the production of a KK gluon in association with a $t\bar{t}$ pair at the LHC energy as a function of the KK gluon mass and with $p_T$ cuts as described in the text.

In Fig. 2, we have plotted this cross-section as a function of the mass of the KK gluon, $M$, for $pp$ collisions at the LHC energy of $\sqrt{s} = 14$ TeV. We have used the CTEQ4M densities [17] and the parton distributions are taken from PDFLIB [18]. For the QCD scale, we use $Q = \sqrt{s}/2$. For this choice of parameters and cuts, we have used ALPGEN to compute the background and we find a background cross-section of 0.33 fb. Assuming an integrated luminosity of 100 fb$^{-1}$, we find from
Fig. 2 that a significance ($\equiv S/\sqrt{B}$) of 5 is obtained for $M = 2790$ GeV. The fact that the kinematic reach that this channel provides in searching for the KK gluon at the LHC is of the same order of magnitude as allowed by precision electroweak measurements is encouraging. Note that we have made no attempt to optimise the significance of the signal and a more judicious choice of cuts could conceivably help in increasing the reach.

![Figure 3: The reach, $M^*$, in KK gluon mass at the LHC as a function of the $p_T$ cut.](image)

Since the choice of 300 GeV for the value of the cut on the $p_T$ of the top quarks coming from the decay of the KK gluon was only an educated guess, we also studied the effect of changing this cut on the significance of the signal. In Fig. 3, we have displayed the results of varying the cut on $p_T$ assuming an integrated luminosity of 100 fb$^{-1}$. For different values of the cut we have plotted the value, $M^*$, of the mass of the $g_{KK}$ for which a significance of 5 is obtained. We find that changing the cut from 200 GeV to 600 GeV increases the reach by about 300 GeV. But in choosing a larger $p_T$ cut one loses out on the number of events so that there are hardly a couple
of background events at a $p_T$ cut of 600 GeV. Therefore, one may have to optimise
the cut by choosing it to be around 300 or 400 GeV which leaves us with a sizeable
number of events.

The preferential coupling of $g_{KK}$ to $t_R$ as opposed to $t_L$ can be exploited to in-
crease the significance of this signal. The chiral coupling of the $g_{KK}$ suggests that the
polarization of the top quarks, studied by looking at its decay products, can prove
to be a very useful discriminator between the signal and the background. However,
in the present paper which is based on a parton-level Monte Carlo study, we have
limited ourselves to studying the kinematic reach of the LHC in the associated pro-
duction process because the 4-top final state that we have focussed on here is not
going to be an easy final state to analyze at the LHC experiments given the com-
binatorial backgrounds from this state that would have to be dealt with to extract
a realistic signal. We have deferred a more detailed study of this signal after im-
plementing it in a hadron-level Monte Carlo. This will enable us to present various
kinematic distributions and top-polarization studies. Nevertheless, the results pre-
sent in this paper are interesting enough to urge experimentalists at the LHC to
consider this process seriously.

Acknowledgements:
One of us (K.S) would like to acknowledge fruitful discussions with Kaustubh Agashe,
Abdelhak Djouadi, Riccardo Rattazzi and Bryan Webber.

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