Higgs pair production due to radion resonance in Randall-Sundrum model: prospects at the Large Hadron Collider

Prasanta Kumar Das 1
Harish-Chandra Research Institute,
Chhatnag Road, Jhusi, Allahabad-211019, India.

Biswarup Mukhopadhyaya 2
Harish-Chandra Research Institute,
Chhatnag Road, Jhusi, Allahabad-211019, India.

Abstract

We consider Higgs pair production at the Large Hadron Collider (LHC) in a Randall-Sundrum scenario containing a radion field. It is shown that the enhanced effective coupling of the radion to gluons, together with contributions from a low-lying radion pole, can provide larger event rates compared to most new physics possibilities considered so far. We present the results for both an intermediate mass Higgs and a heavy Higgs, with a detailed discussion of the background elimination procedure.

1E-mail:pdas@mri.ernet.in
2E-mail: biswarup@mri.ernet.in
1 Introduction

The standard electroweak theory still awaits the discovery of the Higgs boson. After the Large Electron Positron (LEP) collider has set a lower limit of about 114.5 GeV on its mass [1], the responsibility of finding the Higgs now rests mostly on the Large Hadron Collider (LHC). At the same time, puzzles such as the naturalness problem make a strong case for physics beyond the standard model (SM), just around or above the mass scale where the Higgs boson is expected to be found. It is therefore of supreme interest to see if the collider signals of the Higgs contain some imprint of new physics. This necessitates detailed quantitative exploration of a variety of phenomena linked to the production and decays of the Higgs.

In this paper, we have studied pair production of the Higgs boson at the LHC as a possible channel for uncovering new physics effects. In particular, we show that the rate of such production receives a large enhancement in a class of theories with extra compact dimensions, namely, the Randall-Sundrum (RS) type of models containing a radion field.

As has been mentioned above, the large hierarchy between the electroweak scale $M_{EW}$ and the Planck scale $M_{Pl}$ is somewhat puzzling. Whereas theories like supersymmetry and technicolour, each with its own phenomenological implications and constraints, have addressed this question, theories with extra spatial dimensions proposed as a resolution of this problem have recently drawn a lot of attention [2]. Basically, such theories postulate that all hitherto known particles and their standard interactions are confined to our familiar $(3+1)$ dimensional spacetime (on a ‘brane’), while gravity propagates in the ‘bulk’ including additional spacelike compact dimensions. On compactification, the $(3+1)$ dimensional projection of gravity gives rise to both massless and massive graviton modes, where the latter interacts with the standard particles with a strength that can be perceptible in TeV-scale accelerator experiments.

There are two major variants of the above approach. The first one of them [3] proposes a factorizable metric, large compact extra dimensions, and a compatification scheme that gives rise to a continuum of gravitonic modes on the Brane. The integrated effects of this continuum in observable processes have been the subject of a large number of phenomenological studies. However, given the fact that we aim to have a bulk gravity scale in the $TeV$
range as a natural cut-off to the SM (and therefore as a solution to the naturalness problem), such an approach fails to explain a fresh hierarchy that it introduces, namely, one between the higher dimensional gravity scale and the inverse of the large compactification radius.

The other approach is the one proposed by Randall and Sundrum [4]. In this model the fifth dimension corresponds to an $S^1/Z_2$ orbifold and the (4+1) dimensional world is described by the following ‘warped’ metric

$$ds^2 = e^{-2KR_c|\theta|} \eta_{\mu\nu}dx^\mu dx^\nu - R_c^2d\theta^2$$

where $K$ is the bulk curvature constant and $R_c$ is the radius of the extra dimension. The theory postulates two $D_3$ branes, one located at $\theta = 0$ where gravity peaks and the other at $\theta = \pi$ where the SM fields reside. The factor $e^{-2KR_c|\theta|}$ appearing in the metric is known as the warp factor.

The interesting feature of such a theory is that projections of the graviton on the brane at $(\theta = \pi)$ gives, apart from the usual zero mode, a discrete massive tower with spacing on the TeV scale. In addition, the coupling of these massive modes to the SM fields is suppressed not by the Planck mass $M_P$ but by a mass $M_P e^{-KR_c\pi}$. In fact, it can be shown that all quantities with the dimension of mass undergo such an exponential suppression on this brane, thereby offering a tangible solution to the naturalness problem with $KR_c \simeq 12$, i.e. with a combination of masses that are not too widely separated.

The length $R_c$ in this scenario is called the brane separation, and can be related to the vacuum expectation value (vev) of some modulus field $T(x)$ which corresponds to the fluctuations of the metric over the background geometry given by $R_c$. Replacing $R_c$ by $T(x)$, we can rewrite the RS metric at the orbifold point $\theta = \pi$ as

$$ds^2 = g^{vis}_{\mu\nu}dx^\mu dx^\nu - T(x)^2d\theta^2$$

where $g^{vis}_{\mu\nu} = e^{-2\pi KT(x)}\eta_{\mu\nu} = \left(\frac{\Phi(x)}{f}\right)^2 \eta_{\mu\nu}$. Here $f^2 = \frac{24M_5^2}{K}$ and $M_5$, the 5-dimensional Planck scale.

One is thus left with a scalar field $\phi(x)$. This field is called the the radion field [5]. Of course, the modulus field has no potential to start with. Thus one needs to generate a stable vacuum for $T(x)$ at $R_c$, which in turn can give $\phi(x)$ a non-zero vev. This is done in terms of the Goldberger-Wise mechanism [6], using a bulk scalar field with non-zero vev’s at the
two branes, whereupon a potential for the modulus field is generated, and one ends up with a radion field whose mass \(m_\phi\) and vev \(\langle \phi \rangle\) are both within the TeV scale. In particular, the radion can be lighter than the other low-lying gravitonic degrees of freedom. Thus it can very well act as the first messenger of a scenario with compact extra dimensions, and reveal itself in collider experiments. Several studies on the observable implications of the radion are available in the literature [7],[8].

Here we wish to focus on the pair-production of Higgs bosons mediated by the radion at hadron colliders. The two features that can be instrumental in enhancing the signal in this channel are (a) the accessibility of the radion resonance for \(m_\phi > 2m_H\), and (b) the relatively enhanced radion coupling with a pair of gluons at LHC energies. Before we come to the details of the predicted signal, however, let us briefly recapitulate the various interactions of the radion in a theory of the above kind.

We list the relevant interactions of the radion with the SM fields in the next section. The general features of Higgs pair-production via radion are discussed in section 3. Sections 4 and 5 contain discussions on the predicted signals for \(m_h < 2m_W\) and \(m_h > 2m_W\) respectively. We conclude in section 6.

2 Effective interaction of radion with the SM fields

Radion interactions with the SM fields on the TeV brane (i.e. the one located at \(\theta = \pi\)) are governed by 4-dimensional general covariance. The radion essentially couples to the trace of the energy-momentum tensor of the SM fields in the following manner:

\[
\mathcal{L}_{int} = \frac{\phi}{\langle \phi \rangle} T^\mu_\mu(SM)
\]

where \(\langle \phi \rangle\) is the radion vev. There are phenomenological limits on the \(m_\phi - \langle \phi \rangle\) parameter space, from which the lower bound on \(\langle \phi \rangle\) can range from about the electroweak symmetry breaking scale to about a TeV, while \(m_\phi\) can in principle be even lighter than \(m_W\).

The trace of the energy-momentum tensor of the SM fields is given by

\[
T^\mu_\mu(SM) = \sum_\psi \left[ \frac{3i}{2} (\bar{\psi} \gamma_\mu \partial_\nu \psi - \partial_\nu \bar{\psi} \gamma_\mu \psi) \eta^{\mu\nu} - 4m_\psi \bar{\psi} \psi \right] - 2m^2_W W^+ W^- \mu - m^2_Z Z_\mu Z^\mu + (2m^2_h h^2 - \partial_\mu h \partial^\mu h) + ...
\]

...
The photon and the gluons couple to the radion via the usual top-loop diagrams; an added source of enhancement of the coupling is the trace anomaly term \[9\]. This term arises from the fact the radion couples to the trace of the energy-momentum tensor, which can be equated with the four-divergence of the dilatation current associated with scale invariance of the theory. Scale invariance is preserved at the tree level in sectors of the theory which are massless and have no dimensionful couplings. However, quantum corrections can break such invariance, thereby giving a value to the four-divergence, which is proportional to the relevant beta-function. This term augments the coupling of the radion to the trace of the energy-momentum tensor. In general, the contribution of this term to the interaction Lagrangian can be expressed as

\[ T^\mu(SM)_{\text{anom}} = \sum_a \frac{\beta_a(g_a)}{2g_a} C^a_{\mu\nu} G^{a\mu\nu} \]  

For gluons, \( \beta_s(g_s)/2g_s = -[\alpha_s/8\pi] b_{QCD} \) where \( b_{QCD} = 11 - 2n_f/3 \), \( n_f \) being the number of quark flavours. On the whole, the effective \( \phi gg \) interaction is given by

\[ \frac{i\alpha_s\delta_{ab}}{2\pi\langle\phi\rangle} [b_{QCD} + I_{QCD}] (p_1 \cdot p_2 \eta_{\mu\nu} - p_{1\mu}p_{2\nu}) \]  

where \( p_1 \) and \( p_2 \) are the 4-momenta of the gluons \( G^a_\mu \) and \( G^b_\nu \). The function \( I_{QCD} \) is dominated by the top quark loop, and is given by

\[ I_{QCD} = \frac{2}{9} x_t \left[ 2 + 3\sqrt{x_t - 1} \lambda(x_t) - 2(x_t - 1)^2 \right] \]

where \( x_t = \frac{4m_t^2}{m_\phi^2} \), and

\[ \lambda(x_t) = \begin{cases} \sin^{-1} \left( \frac{1}{\sqrt{x_t}} \right), & x_t \geq 1 \\ \frac{1}{2} \left[ \pi + i \ln \left( \frac{1 + \sqrt{1-x_t}}{1 - \sqrt{1-x_t}} \right) \right], & x_t < 1 \end{cases} \]

It is important to note that for \( x_t \geq 1 \), \( n_f = 5 \) and hence \( b_{QCD} = 23/3 \), while for \( x_t < 1 \), \( n_f = 6 \) and the corresponding \( b_{QCD} = 21/3 \).

3 Higgs pair production at LHC

The pair production of a neutral Higgs boson from gluon-gluon fusion at the LHC has already been studied in the context of the SM \[10\] as well as several of its extensions such as the
minimal supersymmetric standard model (MSSM) [11]. As has been already mentioned, there is an expectation of some trace of physics beyond the SM being found in signatures of the Higgs, and this can only be possible by studying Higgs production and decays in as many channels as possible. From this standpoint, the pair production of Higgs deserves attention, because the predicted rate of such production is very small in the SM [10], and thus any excess can be interpreted as the signature of new physics. The mediation of the heavier neutral Higgs in SUSY has been shown to be the source of some enhancement in a region of the parameter space. Theories with extra dimensions, too, have been studied in this context, both in the ADD and RS models [12], where it has been reported that the mediation of gravitons can boost the Higgs pair production rates. What we wish to emphasize, however, is the fact that the presence of a radion in the RS context is of particular significance here. This is because (a) whereas the graviton resonance in the RS scenario is usually at too high a mass to be significant, a relatively less massive radion can be within kinematic reach, and also (b) the enhancement of radion coupling to a pair of gluons via the trace anomaly term jacks up the contributions. We have computed the predicted rates for both the cases of \( m_H < 2m_W \) and \( m_H > 2m_W \), and analyzed the viability of the resulting signals with appropriate event selection strategies. We have used the CTEQ4L parton distribution function, setting the renormalisation scale at the radion mass. We have checked that the predicted results are more or less unchanged if this scale is set, say, at the partonic subprocess centre-of-mass energy.

Using Breit-Wigner approximation for the resonant production of a radion, one finds the following expression for the cross-section for \( pp \rightarrow \phi \rightarrow hh \) (where the dominant subprocess is \( gg \rightarrow \phi \rightarrow hh \)):

\[
\sigma_{hh}(pp \rightarrow \phi \rightarrow hh) = \int dx_1 \int dx_2 \, g_1(x_1) \, g_2(x_2) \hat{\sigma}_{hh}(gg \rightarrow \phi \rightarrow hh)
\]  

(9)

where,

\[
\hat{\sigma}_{hh}(gg \rightarrow \phi \rightarrow hh) = \frac{\Gamma[\phi \rightarrow gg]}{128 \, m_\phi^3 (\phi)^2} \frac{\hat{s} [\hat{s}^2 + 4 m_H^2 \hat{s} + 4 m_h^4]}{[(\hat{s} - m_H^2)^2 + m_H^4 \Gamma_H^2]}
\]  

(10)

where \( \hat{s} \) is centre-of-mass energy for the partonic subprocess. \( \hat{s} = x_1 x_2 s \), \( s(= 14 \, TeV) \) being the proton-proton centre-of-mass energy, and \( x_1, x_2 \), the momentum fractions of the first and
second partons (gluons) respectively. $m_\phi$ and $\Gamma_\phi$ denote the mass and total decay width of the radion [13].

Away from the radion resonance, the term proportional to $\Gamma_\phi^2$ is of little consequence, and the total rate falls as $1/\langle \phi \rangle^4$ if the radion vev is increased. This is because the decay width of the radion in any channel is proportional to $1/\langle \phi \rangle^2$. Near resonance, on the other hand, the contribution is dictated by the term with $\Gamma_\phi$ in the radion propagator. Given the above dependence of the decay width on $\langle \phi \rangle$, the net contribution near resonance becomes practically independent of the radion vev.

4 Event selection strategy and results: $m_h < 2m_W$

The signal of Higgs pair production depends on the final states produced by Higgs decays. We shall broadly consider two mass ranges here, one corresponding to $m_h < 2m_W$, and the other, to $m_h > 2m_W$.

In the first case, the dominant decay channel for each Higgs is $H \rightarrow b\overline{b}$, so that the signal consists of four b-jets. Normally, such events are beset with a huge background arising from the following sources:

- The QCD production of $(\overline{b}b)(\overline{b}b)$,
- The QCD production of $(\overline{b}b)(jj)$ with the two non-$b$ jets misidentified as $b$,
- The production of two($\overline{b}b$) pairs from different partonic collisions, arising either from the same pair of protons or from two different pairs.
- The electroweak production of $Z(\overline{b}b)$ and $W(\overline{b}b)$.

A very relevant discussion of these backgrounds and the cuts that can be used to eliminate them can be found in [14]. These cuts have been adapted to our context in this study. Basically, the contributions from mistagged b-jets and different partonic collisions from the same proton pair are relatively small, the ‘pile-up’ effects arising from different proton pairs can be also handled with the help of effective b-tagging. For this, the following b-tagging
efficiencies have been assumed in different transverse momentum ($p_T$) ranges:\footnote{15}:

\[
\begin{align*}
\epsilon_b & = 0.6, \quad \text{for} \quad p_T > 100 \text{ GeV} \\
& = 0.1 + p_T/(200 \text{GeV}), \quad \text{for} \quad 40 \text{ GeV} \leq p_T \leq 100 \text{ GeV} \\
& = 1.5p_T/(100 \text{GeV}) - 0.3, \quad \text{for} \quad 25 \text{ GeV} \leq p_T \leq 40 \text{ GeV}
\end{align*}
\] (11)

where we also assume that $b$-jets are tagged only in the pseudorapidity region $\eta_b \leq 2$.

With such tagging efficiencies, it is found that the irreducible backgrounds become manageable with a further set of event selection criteria, particularly if the Higgs mass is not in close vicinity of the Z-mass (a fact that is confirmed by LEP data). The exact cuts used in our analysis are listed below.

- **Reconstruction of the Higgs boson mass**: Since we do not differentiate between $b$-and $\bar{b}$-jets, there will be four $b$ in the final state and hence one has three possible ways to pair them up. We choose those pairs which correspond to the smallest invariant mass difference between them. The reconstructed Higgs boson mass is then defined by $M_h = [M_{b_1,b_2} + M_{b_3,b_4}]/2$. Keeping in mind the finite resolution of the reconstruction procedure, this invariant mass has been smeared by a Gaussian distribution of width $\sigma \approx \sqrt{M_h}$. The window for the reconstructed Higgs boson mass in which our search is being carried out is

\[
0.9m_{h,in} - 1.5\sigma \leq M_h \leq 0.9m_{h,in} + 1.5\sigma
\] (12)

- The mass difference between the invariant masses of the two pairs are required to be below the width of Gaussian smear:

\[
\delta M_h = |M_{b_1,b_2} - M_{b_3,b_4}| \leq 2\sigma
\] (13)

- Large azimuthal angles between the two jets of each pair in the transverse plane are demanded:

\[
\delta \phi_{b_1b_2}, \delta \phi_{b_3b_4} > 1
\] (14)
Figure 1: The Signal cross-section(pb) for the process $pp(gg) \rightarrow \phi \rightarrow hh \rightarrow b\overline{b}b\overline{b}$ against the Higgs mass for $m_{\phi} = 250, 375$ and $500$ GeV and $\langle \phi \rangle = 500, 1000$ GeV.

- Events where the difference between the angles $\delta \phi_{b_1b_2}, \delta \phi_{b_3b_4}$ is small are retained:

\[ |\delta \phi_{b_1b_2} - \delta \phi_{b_3b_4}| < 1 \quad (15) \]

- The b-jet are subjected to the following $p_T$-cuts:

\[ p_{T,min} > M_h/4; \quad p_{T,max} > M_h/4 + 2\sigma \quad (16) \]

- In order to further utilise the hardness of the b-jets in eliminating backgrounds, a minimum value is imposed on the $4b$ invariant mass $M_{4b}$:

\[ M_{4b} > 1.9M_h - 3\sigma \quad (17) \]

After applying all the above cuts, the cross-section for the process $pp(gg) \rightarrow \phi \rightarrow hh \rightarrow b\overline{b}b\overline{b}$ is obtained for different radion vevs $\langle \phi \rangle$, radion masses $m_{\phi}$ and Higgs masses $m_h$. In Figure
we have plotted this cross-section against $m_h$ (GeV) ($\approx m_{h,in}$) for $\langle \phi \rangle = 500, 1000$ GeV and for $m_\phi = 250, 375$ and 500 GeV.

The first conclusion to draw from the figures is that there is a substantial enhancement of the total rates over what is predicted in the standard model as well as in the case of MSSM. In addition, it also exceeds the predicted rates in the ADD-and RS-type models when such models assume graviton mediation to be the only new effect. This is particularly evident from the fact that the numbers presented here are after all the cuts have been applied and the b-tagging efficiency has been folded in, which effectively causes well below one per cent of the signal events to survive.

As has been mentioned before, such an enhancement has two main sources, namely (a) the availability of the radion resonance, and (b) the enhanced coupling at the gluon-gluon-radion vertex. It can also be noticed by comparing figures 1(a) and 1(b) that the rates for Higgs masses corresponding to the onset of resonance are almost independent of the radion vev. Since we are showing here only that range of Higgs masses where the decay into the $b\bar{b}$ is dominant, such peaking effect is prominent only for $m_\phi = 250$ GeV; for the higher values of $m_\phi$, $m_h$ becomes large enough for the branching fraction in the $b\bar{b}$ channel to drop drastically, before the peak can be reached.

**Significance contours:**

With an integrated luminosity of $100 \, fb^{-1}$, the above rates indicate a rather impressive prospect of detecting pair-produced Higgs bosons if radion mediation is operative. To gauge the actual situation, however, one must also remember that the backgrounds are never totally eliminated, and thus one must examine how the signal fares compared to the surviving backgrounds. This is depicted in figures 2(a) and (b) with contour plots in $S/\sqrt{B}$, where, $S$ and $B$ are the number of events corresponding to the signal and backgrounds respectively with the above luminosity. In calculating the backgrounds, we have taken both statistical and systematic effects into account, assuming the systematic uncertainty to be 2% of the total background and adding it in quadrature to the computed background itself. The contour plots corroborate our expectation that the signal really stands out over a large region of the parameter space. Thus the so-called intermediate mass range acquires a high degree of visibility if a low-lying radion is present in the theory.
Let us now consider the case of a somewhat heavier Higgs, which can decay into a pair of $W$’s or $Z$’s. The signal in such a case consists in $4W$ or $4Z$ final states, where the SM backgrounds are negligibly small. In Figures 3 and 4, we have plotted the cross sections for the processes $pp(gg) \to \phi \to hh \to W^+W^-W^+W^-$ and $pp(gg) \to \phi \to hh \to ZZZZ$ (for cases where the $W$ and the $Z$ decay into electrons and muons with total branching ratios of about 0.2 and 0.06 respectively) against the Higgs mass $m_h$ for different $m_\phi$ and $\langle \phi \rangle$. In addition, an average detection efficiency of 75% per lepton has been assumed. It can be seen from the figures that the $4Z$ final states are very unlikely to be seen at the LHC. The $4W$ final state, however, is quite substantial, especially in view of the absence of backgrounds. One may thus expect to see events ranging in number from a few tens to several thousands, depending on the radion mass, so long as the Higgs mass is within approximately 200 GeV. For a radion mass on the order of 375 GeV, a slight peaking behaviour can be seen around
Figure 3: The Signal cross-section (fb) for the process $pp(gg) \rightarrow \phi \rightarrow hh \rightarrow W^+W^-W^+W^-$ where each of the produced $W$ decay leptonically ($W \rightarrow l\nu_l, l = e, \mu$) against the Higgs mass for $m_\phi = 375$ and 500 GeV and $\langle \phi \rangle = 500, 1000$ GeV.

$m_h \simeq 180$ GeV. For a heavier radion, however, the decay width becomes so large that the peak is washed out. In this case, the $\Gamma_\phi m_\phi^2$-term in the Breit-Wigner propagator gives a substantial contribution even where the subprocess centre-of-mass energy is considerably away from $m_\phi^2$. Such a contribution is responsible for the lack of dependence of the radion vev, following arguments given earlier.

6 Summary and conclusion:

We have looked at the pair production of the Higgs boson enhanced by radion mediation at the LHC. For $m_h < 2m_W$, 4$b$ final states are investigated. It is found that in spite of substantial backgrounds to start with, a careful event selection strategy can lead to a
Figure 4: The Signal cross-section (fb) for the process pp(gg) → φ → hh → ZZZZ where each of the produced Z decay leptonically (Z → l⁺l⁻, l = e, µ) against the Higgs mass for m_φ = 375 and 500 GeV and ⟨φ⟩ = 500, 1000 GeV.

discovery potential at the 5σ or even 10σ level. For m_h > 2m_W, on the other hand, one has to depend on the 4W final states which are almost background-free. There, too, one should be able to see anything between about 100 and 800 events so long as the radion vev is within a TeV and the Higgs mass lies within about 200 GeV. Thus the presence of a radion can boost the very conspicuous phenomenon of Higgs pair production, over a large region of the parameter space, including the entire range of Higgs mass favoured by precision electroweak data.

Acknowledgment:

We thank Uma Mahanta and Sreerup Raychaudhuri for sharing their radion decay code with us. PKD would like to thank Partha Konar for computational help. The work of BM has been partially supported by the Board of Research in Nuclear Science, Government of India.
References

[1] T. Junk, The LEP Higgs Working Group, at LEP Fest October 10th 2000.

[2] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys.Lett. B429, 263 (1998).

Some precursors of the model are Akama L, Lect. Notes, 176 267 (1982); V. Rubakov
and M. Shaposhnikov, Phys.Lett. B125, 136 (1984); A. Barnaveli and O. Kancheli, Sov.
J. Nucl. Phys. 51, 573 (1990); I. Antoniadis, Phys.Lett. B246, 377 (1990); I. Antoniadis,
C. Muñoz and M. Quiros, Nucl.Phys. B397, 515, (1993); I. Antoniadis, K. Benakli and
M. Quiros, Phys.Lett., B331, 313 (1994); V. Rubakov, Phys.Usp. 44, 871-893 (2001).

[3] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys.Lett B429, 263 (1998); I. Anto-
niadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys.Lett. B463, 257 (1998).

[4] L. Randall and R. Sundrum, Phys. Rev. Lett. 83, 3370 (1999); L. Randall and R.
Sundrum, Phys. Rev. Lett. 83, 4690 (1999); H. Davoudiasl, J. L. Hewett, and T. G.
Rizzo, Phys. Rev. Lett. 84, 2080 (2000).

[5] G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B595, 250 (2001); W. D.
Goldberger and M. B.Wise, Phys. Lett. B475, 275-279 (2000); W. D. Goldberger and
I. Z. Rothstein, Phys. Lett. B491, 339 (2000).

[6] W. D. Goldberger and M. B.Wise, Phys. Rev. Lett. 83, 4922 (1999); W. D. Goldberger
and M. B.Wise, Phys. Rev. D60, 107505 (1999).

[7] K. Cheung, Phys. Rev. D63, 056007, 2001, hep-ph/0009232; C. Csaki, M. Graesser, L.
Randall and J. Terning, Phys. Rev. D62, 045015 (2000); U. Mahanta and A. Datta,
Phys. Lett. B483, 196 (2000); U. Mahanta and S. Rakshit, Phys. Lett. B480, 176
(2000); S. Bae, P. Ko, H. Lee and J. Lee, Phys. Lett. B487, 299 (2000).

[8] P. Das and U. Mahanta, Phys. Lett. B520, 307 (2001), Phys. Lett. B528, 253 (2002),
hep-ph/0110309, hep-ph/0201260, hep-ph/0202193; M. Chaichian., A. Datta, K.Huitu
and Z. Yu, Phys. Lett. B524, 161 (2002).

[9] J. C. Collins, A. Duncan and S. D. Joglekar, Phys. Rev. D16, 438-449, (1977).
[10] T. Plehn, M. Spira and P. M. Zerwas, *Nucl. Phys.* **B479**, 46 (1999); D. A. Dicus, C. Kao and S. Willenbrock, *Phys. Lett.* **B203**, 457 (1988).

[11] A. Belyaev, M. Drees, O. J. Eboli, J. K. Mizukoshi and S. F. Novaes *Phys. Rev.* **D60**, 075008 (1999); A. A. Barrientos Bendezu and B. A. Kniehl *Phys. Rev.* **D64**, 035006 (2001); A. Belyaev, M. Drees, O. J. Eboli, J. K. Mizukoshi and S. F. Novaes, hep-ph/9910400, C. S. Kim, K. Y. Lee and J. Song, *Phys. Rev.* **D64**, 015009 (2001).

[12] C. S. Kim, K. Y. Lee and J. Song, *Phys. Rev.* **D64**, 015009 (2001).

[13] P. K. Das, U. Mahanta and S. R. Raychaudhuri, *in preparation*.

[14] A. Belyaev, M. Drees and J. K. Mizukoshi, *Eur.Phys.J.* C**17**, 337-351 (2000).

[15] CMS collaboration, S. Abdullin *et al.*, *J. Phys.* G**28**, 469 (2002).