Probing the radion-Higgs mixing at photon colliders

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Abstract
In the Randall-Sundrum model, the radion-Higgs mixing is weakly suppressed by the effective electroweak scale. A novel feature of the gravity-scalar mixing would be a sizable three-point vertex of $h^{(n)}_{\mu\nu} h - \phi$. We explore the potential of photon colliders, achieved by the laser backscattering technique, in probing the radion-Higgs mixing via the associated production of the radion with the Higgs boson. The advantage of photon colliders is the capability of adjusting the polarization of the incoming photons such that the signal of the spin-2 graviton exchange can be largely enhanced. The enhancement factor is shown to be about 5, except for small-$\xi$ region. We also study the corresponding backgrounds step-by-step in detail.

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I. INTRODUCTION

The standard model (SM) has been extraordinarily successful up to now in explaining all experimental data on the electroweak interactions of the gauge bosons and fermions. However, the master piece of the SM, the Higgs boson, which is responsible for the electroweak symmetry breaking, still awaits experimental discovery [1]. The direct search has excluded the SM Higgs boson mass below about 114 GeV [2], while the indirect evidences from the electroweak precision data put an upper bound for a SM Higgs of 208 GeV at 95% C.L. [3]. The precision measurements and the direct searches are getting into the situation that they begin to contradict each other. There have been various studies which can ease the situation. One possibility is that the SM Higgs boson mixes with another scalar boson such that the Higgs branching ratio into $b\bar{b}$ becomes smaller and thus escapes the limit of direct search so far. Disentangling the nature of this new scalar state is very significant and challenging.

It has been pointed out that the radion of the Randall and Sundrum (RS) model [4] can play the role of such a scalar boson. The RS model consists of an additional spatial dimension of a $S^1/Z_2$ orbifold introduced with two 3-branes at the fixed points. A geometrical suppression factor, called the warp factor, explains the huge hierarchy between the electroweak and Planck scale with moderate values of model parameters. The presence of a radion, the quantum degree with respect to the fluctuation of the brane separation, naturally emerges from the stabilizing process [5, 6]. As various stabilization mechanisms suggest, the radion is generically much lighter than the Kaluza-Klein (KK) states of any bulk field. In the literature, phenomenological aspects of the radion have been studied such as its decay modes [7, 8], its effects on the electroweak precision observations [9], and its signatures at present and future colliders [10].

The radion-Higgs mixing is originated from the gravity-scalar mixing term, $\xi R(g_{\text{vis}}) \tilde{H}^\dagger \tilde{H}$, where $R(g_{\text{vis}})$ is the Ricci scalar of the induced metric $g_{\text{vis}}^{\mu\nu}$, and $\tilde{H}$ is the Higgs field in the five-dimensional context. It has been shown that the radion-Higgs mixing can induce significant deviations to the properties of the SM Higgs boson [11, 12, 13, 14, 15]. A complementary way to probe the radion-Higgs mixing is the direct search for the new couplings exclusively allowed with a non-zero mixing parameter $\xi$. One good example is the tri-linear vertex among the KK graviton field $h_{\mu\nu}^{(n)}$, the Higgs boson $h$, and the radion $\phi$. In Refs. [16, 17], we have shown that probing the vertex $h_{\mu\nu}^{(n)}h-\phi$ through the $h\phi$ production at $e^+e^-$ colliders
and hadronic colliders can provide very useful information on the radion-Higgs mixing, irrespective of the mass spectrum of the Higgs boson and the radion.

In this work, we turn to photon colliders achieved by the laser backscattering technique\cite{18}. The process that we investigate is

\[ \gamma \gamma \rightarrow h^{(n)}_{\mu \nu} \rightarrow h \phi, \]

where \( h^{(n)}_{\mu \nu} \) denotes the \( n \)-th KK state of the RS graviton. Since the polarization of incoming photons can be adjusted by tuning the polarization of the electron or positron beam and the laser beam, the signal can be largely enhanced because the exchanged graviton is a spin-2 particle. This is the biggest advantage of photon colliders in this regard. The observation of the rare decay of a KK graviton into \( h \phi \) is then the direct and exclusive signal of the radion-Higgs mixing. In addition, the characteristic angular distribution could reveal the exchange of massive spin-2 KK gravitons.

This paper is organized as follows. In Section II, we calculate the production cross section of \( \gamma \gamma \rightarrow h^{(n)}_{\mu \nu} \rightarrow h \phi \) folded with the photon luminosity function. Section III deals with the feasibility of detecting the \( h \phi \) final states by considering specific decay channels of the Higgs boson and the radion. We summarize at the end of Section III. Note that we shall use \( G^{(n)} \) or \( h^{(n)}_{\mu \nu} \) to denote the \( n \)-th KK graviton state interchangeably.

II. CALCULATION OF \( \gamma \gamma \rightarrow h^{(n)}_{\mu \nu} \rightarrow h \phi \)

The RS scenario is based on a five-dimensional spacetime of a \( S^1/\mathbb{Z}_2 \) orbifold which has the finite size of \( b_0 \). The warped factor, \( \Omega_0 = e^{-m_0 b_0/2} \), with a moderate value of \( m_0 b_0/2 \simeq 35 \) can solve the gauge hierarchy problem. In terms of the KK graviton field \( h^{(n)}_{\mu \nu} \) and the canonically normalized radion field \( \phi_0 \), the four-dimensional effective Lagrangian is then

\[ L = -\frac{\phi_0}{\Lambda_{\phi}} T^\mu_\mu - \frac{1}{\Lambda_W} T^{\mu \nu}(x) \sum_{n=1}^{\infty} h^{(n)}_{\mu \nu}(x), \]

where \( \Lambda_{\phi} \) is the vacuum expectation value (VEV) of the radion field, \( T^\mu_\mu \) is the trace of the symmetric energy-momentum tensor \( T^{\mu \nu} \), and \( \Lambda_W \) is the effective electroweak scale. Both effective interactions are suppressed by the electroweak scale, not by the Planck scale. The gravity-scalar mixing term of \( S_\xi = \xi \int d^4x \sqrt{g_{\text{vis}}} R(g_{\text{vis}}) \bar{H} H \) is allowed as it respects all the SM symmetries and Poincare invariance. Here \( g_{\text{vis}} \) is the induced metric on the visible
brane, $R(g_{\text{vis}})$ is the Ricci scalar, $H_0 = \Omega_0 \bar{H}$, and the dimensionless parameter $\xi$ of order one denotes the size of the mixing term. This $\xi$-term mixes the $h_0$ and $\phi_0$ fields into the mass eigenstates of $h$ and $\phi$ fields [14]:

$$
\begin{pmatrix}
  h_0 \\
  \phi_0
\end{pmatrix} =
\begin{pmatrix}
  d & c \\
  b & a
\end{pmatrix}
\begin{pmatrix}
  h \\
  \phi
\end{pmatrix}.
$$

(3)

We refer the detailed expressions for $a$, $b$, $c$, and $d$ to Ref. [17].

All phenomenological signatures of the RS model are then determined by five parameters of

$$
\xi, \, \Lambda_\phi, \, \frac{m_0}{M_{\text{Pl}}}, \, m_\phi, \, m_h,
$$

(4)

which in turn determine $\hat{\Lambda}_W = \Lambda_\phi/\sqrt{3}$ and KK graviton masses $m_{G^{(n)}} = x_n m_0 \hat{\Lambda}_W/(M_{\text{Pl}} \sqrt{2})$ with $x_n$ being the $n$-th root of the first order Bessel function of the first kind. The ratio $m_0/M_{\text{Pl}}$ is assumed in $[0.01, \, 0.1]$ to avoid too large bulk curvature [19]. In what follows, we fix the ratio $m_0/M_{\text{Pl}} = 0.1$. The $\Lambda_\phi$ or $\hat{\Lambda}_W$ is constrained by the Tevatron Run I data of Drell-Yan process and by the electroweak precision data: For $m_0/M_{\text{Pl}} = 0.1$, $m_{G^{(1)}} \gtrsim 500$ GeV yields $\Lambda_\phi \gtrsim 3.2$ TeV [20]. Therefore, we consider the case of $\Lambda_\phi = 3.5$ TeV and $m_0/M_{\text{Pl}} = 0.1$, of which the effect of radion on the oblique parameters is small [9]. Then, the first KK graviton mass is about 547.5 GeV. The radion mass is expected to be light as one of the simplest stabilization mechanisms predicts $m_{\phi_0} \sim \hat{\Lambda}_W/40$ [3]. In addition, the Higgs boson mass is set to be 120 GeV throughout the paper.

The gravity-scalar mixing modifies the couplings among the $h$, $\phi$ and $h_{\mu\nu}^{(n)}$. In particular, a non-zero $\xi$ gives rise to new tri-linear vertices of

$$
\begin{align*}
   h_{\mu\nu}^{(n)} - h - \phi, \quad h_{\mu\nu}^{(n)} - \phi - \phi, \quad h - \phi - \phi, \quad \phi - \phi - \phi.
\end{align*}
$$

(5)

Due to the suppressed coupling of a photon with a Higgs boson or a radion, the $\gamma\gamma$ collider is expected to access the $h_{\mu\nu}^{(n)} - h - \phi$ or $h_{\mu\nu}^{(n)} - \phi - \phi$ vertex directly. In addition, the coupling strength of $h_{\mu\nu}^{(n)} - \phi - \phi$ is much smaller than that of $h_{\mu\nu}^{(n)} - h - \phi$, by a factor of $\gamma \equiv v/\Lambda_\phi \ll 1$. Here $v$ is the VEV of the Higgs boson, which is 246 GeV. Therefore, the channel $h_{\mu\nu}^{(n)} \rightarrow h\phi$ is the most effective in probing the radion-Higgs mixing, with the vertex denoted by

$$
\langle \phi \mid h_{\mu\nu}^{(n)} \mid \phi \rangle \equiv i g_\phi \frac{2k_{1\mu} k_{2\nu}}{\hat{\Lambda}_W},
$$

(6)

\[ -4 - \]
\( \gamma \gamma \rightarrow h \phi \)

where \( \hat{g}_{h\phi} = 6 \gamma \xi [a(\gamma b + d) + bc] + cd \) and \( k_{1,2} \) is the four-momentum of the scalar particles.

Then, the partial decay width of \( h_{\mu\nu}^{(n)} \rightarrow h\phi \) is given as

\[
\Gamma(h_{\mu\nu}^{(n)} \rightarrow h\phi) = \frac{\hat{g}_{h\phi}^2}{240 \pi \Lambda^2} \beta \left[ 1 - \left( \sqrt{\mu_{hG}} + \sqrt{\mu_{\phi G}} \right)^2 \right] \left[ 1 - \left( \sqrt{\mu_{hG}} - \sqrt{\mu_{\phi G}} \right)^2 \right], \tag{7}
\]

where \( \mu_{xy} = \left( \frac{m_x}{m_y} \right)^2 \), \( \beta = \frac{\lambda_1}{2} \left( 1, m_h^2 / \hat{s}, m_\phi^2 / \hat{s} \right) \hat{g}_{h\phi} \), and \( \lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc \).

In Ref. 17, it was shown that the branching ratio \( \text{Br}(h_{\mu\nu}^{(n)} \rightarrow h\phi, \phi\phi) \), which would vanish in the limit \( \xi \rightarrow 0 \), is of the order of \( \mathcal{O}(10^{-3}) \).

For the process of

\( \gamma(q_1, \lambda_1) + \gamma(q_2, \lambda_2) \rightarrow h(k_1) + \phi(k_2), \tag{8} \)

the Feynman diagrams are shown in Fig. 1. Here \( \lambda_{1,2} \) is the polarization of the high energy photons. The helicity amplitudes \( \mathcal{M}_{\lambda_1 \lambda_2} \) including the \( h \) and \( \phi \) mediation are

\[
\mathcal{M}_{++} = -\frac{\hat{s}}{2\nu \Lambda} \left( \hat{c}_{\gamma h} \hat{g}_h \mathcal{D}_h + \hat{c}_{\gamma \phi} \hat{g}_\phi \mathcal{D}_\phi \right), \tag{9}
\]
\[
\mathcal{M}_{+-} = -\frac{\hat{s}}{\Lambda^2} \lambda(1, m_h^2 / \hat{s}, m_\phi^2 / \hat{s}) \hat{g}_{h\phi} \mathcal{D}_G \sin^2 \theta^*, \tag{10}
\]

where \( \hat{s} = (q_1 + q_2)^2 \), \( \theta^* \) is the scattering angle of the Higgs boson in the \( \gamma\gamma \) c.m. frame, \( \hat{c}_{\gamma h} \) and \( \hat{c}_{\gamma \phi} \) are

\[
\hat{c}_{\gamma h} = -\frac{\alpha}{2\pi} \left( (d + \gamma b) \sum_i e_i^2 N_i^h F_i(4m_i^2 / m_h^2) + \frac{11}{3} \gamma b \right), \tag{11}
\]
\[
\hat{c}_{\gamma \phi} = -\frac{\alpha}{2\pi} \left( (c + \gamma a) \sum_i e_i^2 N_i^\phi F_i(4m_i^2 / m_\phi^2) + \frac{11}{3} \gamma a \right).
\]

We refer the expressions for \( F_{1/2}, F_1 \), and \( \hat{g}_{h,\phi} \) to Ref. 17. The propagator factors of the

FIG. 1: Feynman diagrams for the process \( \gamma\gamma \rightarrow h\phi \).
KK-graviton, the Higgs boson, and the radion are given by

\[ D_G = \sum_{n=1}^{\infty} \frac{\hat{s}}{\hat{s} - m_{G(n)}^2 + i m_{G(n)} \Gamma_{G(n)}} , \quad D_{h,\phi} = \frac{\hat{s}}{\hat{s} - m_{h,\phi}^2 + i m_{h,\phi} \Gamma_{h,\phi}} . \]  

(12)

Note that \( M_{++} = M_{--} \) and \( M_{+-} = M_{-+} \) are guaranteed by CP invariance. In principle, the photon polarization can separate the contribution of the scalar mediation from that of KK gravitons even though leading contribution comes from the KK graviton mediation.

Brief comments on \( \gamma\gamma \) colliders are in order here \[18\]. From the head-on collisions between the laser and energetic electron (or positron) beams, high energy photons are produced. If we denote the fraction of the photon beam energy to the initial electron beam energy by \( x = \frac{E_\gamma}{E_e} \), its maximum value is \( x_{\text{max}} = z/(1 + z) \) where \( z = 4E_e \omega_0/m_e^2 \). Here \( E_\gamma, E_e, \omega_0 \) are the photon, electron and laser beam energies, respectively. Usually, \( z \) is optimized to be \( 2(1 + \sqrt{2}) \) to avoid the \( e^+ e^- \) pair production through the interactions of the laser beam and the backward scattered photon beam. In the numerical analysis, we consider the following range for \( x_{1,2} \):

\[ \sqrt{0.4} \leq x_{1,2} \leq x_{\text{max}} \bigg|_{z = 2(1 + \sqrt{2})} . \]  

(13)

With the given polarizations of the laser and parent electron (positron) beams, their Compton back-scattering leads to the differential cross section

\[ \frac{d\sigma}{d\cos\theta^*} = \frac{1}{32\pi s_{ee}} \int \int dx_1 dx_2 \frac{f(x_1) f(x_2)}{x_1 x_2} \lambda^{1/2} \left( \frac{1}{s}, \frac{m_h^2}{\hat{s}} \right) \lambda \left( \frac{m_{h,\phi}^2}{\hat{s}}, \frac{m_{\phi}^2}{\hat{s}} \right) \times \left[ \left( 1 + \xi_2(x_1) \xi_2(x_2) \right) |M_{J_z=0}|^2 + \left( 1 - \xi_2(x_1) \xi_2(x_2) \right) |M_{J_z=2}|^2 \right] , \]  

(14)

where \( s_{ee} = \hat{s}/(x_1 x_2) \) is the square of the c.m. energy of the parent \( e^+ e^- \) collision, and

\[ |M_{J_z=0}|^2 = \frac{1}{2} \left[ |M_{++}|^2 + |M_{--}|^2 \right] , \]  

\[ |M_{J_z=2}|^2 = \frac{1}{2} \left[ |M_{+-}|^2 + |M_{-+}|^2 \right] . \]  

(15)

Here \( f(x) \) is the photon luminosity function and \( \xi_2(x) \) is the averaged circular polarization of the back-scattered photon beam, both of which depend on the polarizations of the electron \( P_e \) and laser beam \( P_l \). The explicit expressions for \( f(x) \) and \( \xi_2(x) \) are

\[ f(x, P_e, P_l; z) = \frac{1}{\hat{s}_C} C(x) , \]  

(16)

where

\[ \hat{s}_C = \left( 1 - \frac{4}{z} - \frac{8}{z^2} \right) \ln(z + 1) + \frac{1}{2} \left( 1 + \frac{8}{z} - \frac{1}{2(z + 1)^2} \right) \]  

(17)
FIG. 2: Total cross section for $\gamma\gamma \to h\phi$ in fb as a function of $\xi$ for $m_\phi = 30, 70, 150$ GeV at $e^+e^-$ colliders running in the $\gamma\gamma$ mode using laser backscattering. All of the polarizations are set to be zero, $m_h = 120$ GeV, $\Lambda_\phi = 3.5$ TeV, and $\sqrt{s_{ee}} = 500$ GeV.

\[ +P_e P_l \left( \frac{1 + 2}{z} \ln(z + 1) - \frac{5}{2} + \frac{1}{z + 1} - \frac{1}{2(z + 1)^2} \right). \]

and

\[ C(x) \equiv \frac{1}{1-x} + (1-x) - 4r(1-r) - P_e P_l r z(2r-1)(2-x), \]

where $r \equiv x/[z(1-x)]$. The average helicity $\xi_2(x, P_e, P_l; z)$ is given by

\[ \xi_2(x, P_e, P_l; z) = \frac{1}{C(x)} \left\{ P_e \left[ \frac{x}{1-x} + x(2r-1)^2 \right] - P_l \left( 2r - 1 \right) \left( 1 - x + \frac{1}{1-x} \right) \right\}. \]

In Fig. 2 we plot the total cross section as a function of $\xi$ for $m_\phi = 30, 70, 150$ GeV. We set $m_h = 120$ GeV, $\Lambda_\phi = 3.5$ TeV and $\sqrt{s_{ee}} = 500$ GeV. Note that the requirement of positive definiteness of the mass and kinetic terms limits the range of the mixing parameter $\xi$. Here we assume that all beams are unpolarized. For $m_\phi = 30$ GeV, the maximum total cross section can reach about 1 fb, which will produce about 1000 events with 1 ab$^{-1}$ luminosity. For heavier radion mass, the more restricted $\xi$-range reduces the maximum of the total cross section, e.g., to several 10$^{-2}$ fb for $m_\phi = 150$ GeV.

Since the $\gamma\gamma \to h\phi$ process is practically mediated by the massive spin-2 KK graviton, the angular distribution shows its characteristic behavior proportional to $\sin^4 \theta^*$, as shown in Eq. (10). In Fig. 3 we plot the normalized differential cross section $(1/\sigma)d\sigma/d\cos \theta^*$ versus $\cos \theta^*$ for $m_\phi = 30, 70, 150$ GeV. Here we have set $\Lambda_\phi = 3.5$ TeV, and $\xi = 0.26$ which is the allowed maximum value for $m_\phi = 150$ GeV. For $m_\phi = 30$ GeV and $m_\phi = 70$
The average helicity function $\xi_2(x)$ becomes zero for $P_e = P_l = 0$. 

$^{1}$ The average helicity function $\xi_2(x)$ becomes zero for $P_e = P_l = 0$. 

FIG. 3: Normalized differential cross section $\frac{1}{\sigma_{tot}} \frac{d\sigma}{d\cos \theta^*}$ for $m_\phi = 30, 70, 150$ GeV. The $m_\phi$ dependence is mild. All of the polarizations are set to be zero, $m_h = 120$ GeV, $\Lambda_\phi = 3.5$ TeV and $\xi = 0.26$.

GeV, the distributions are practically the same, proportional to $\sin^4 \theta^*$. For $m_\phi = 150$ GeV, the enhanced contributions from the Higgs boson and radion alter the behavior slightly. Especially, at the end points ($\cos \theta^* = \pm 1$) the cross section comes solely from the Higgs and radion exchanges. One crucial reason is that the heavy radion mass reduces the KK-graviton contribution which has an additional factor of $\lambda^2(1, m_h^2/s, m_\phi^2/s)$ with respect to the $h/\phi$ contribution, as can be seen from Eqs. (9) and (10).

As the general expression of Eq. (14) suggests, the appropriate adjustment of the electron (or positron) and laser beam polarizations can enhance the production rate. As can be seen from Eq. (14), non-zero and negative $\xi_2(x_1)\xi_2(x_2)$ can enhance the graviton contribution. Another merit is that the polarized beams can enhance the energy of back-scattered photon beam and thus increase the signal cross section. Figure 4 presents $\hat{f}(x)/x$ where $\hat{f}(x) = f(x)\xi_2(x)$ for the polarized beam while $\hat{f}(x) = f(x)$ for the unpolarized beam$^1$. Compared to the unpolarized beam case, $(P_e, P_l) = (1, -1)$ combination generates a more energetic photon beam. Moreover, the spin-2 nature of KK graviton prefers opposite polarizations for $e^+$ and $e^-$ beams. Therefore, the optimal polarizations would be $(P_{e-}, P_{e+}, P_{l1}, P_{l2}) = (1, -1, -1, 1)$.

With the optimal polarization combination, we plot the ratio of the polarized total cross section to unpolarized one, as a function of $\xi$ for $m_\phi = 30, 70, 150$ GeV in Fig. 5.
FIG. 4: $\hat{f}(x)/x$ versus $x$ for various polarizations of the electron and laser beams, denoted by $(P_e, P_l)$. The $\hat{f}(x) = f(x)$ for unpolarized beams ($P_e = P_l = 0$), and $\hat{f}(x) = f(x)\xi_2(x)$ for polarized beams with $P_e = 1$. The beam polarizations of $(P_e, P_l) = (1, -1)$ generate the highest energy photon beam from the parent electron beam.

FIG. 5: The ratio of the polarized total cross section to unpolarized one, as a function of $\xi$ for $m_\phi = 30$, 70, 150 GeV. The polarization is set to be $(P_{e-}, P_{e+}, P_{l1}, P_{l2}) = (1, -1, -1, 1)$, $m_h = 120$ GeV, $\sqrt{s_{ee}} = 500$ GeV, and $\Lambda_\phi = 3.5$ TeV.

dependence is negligible in the region of $|\xi| \gtrsim 0.3$. The enhancement is significant and better than the naive estimation of a factor of two from Eq. (14). The enhancement factor can
reach upto 5.1 for $m_\phi = 150$ GeV, and about 4.8 for $m_\phi = 30, 70$ GeV. Even for the $m_\phi = 70$ GeV case, the total cross section can be about 1 fb. Under conservative assumption for the electron and positron beam polarization, we set $(P_{e^-}, P_{e^+}, P_{l1}, P_{l2}) = (0.8, -0.6, -1, 1)$: the maximum of the enhancement factor is about 3.3 for $m_\phi = 30, 70$ GeV, and about 3.6 for $m_\phi = 150$ GeV.

### III. DECAYS AND DETECTION OF THE RADION-HIGGS PAIR

In this section, we consider the feasibility of detecting $h\phi$ pair production. For a Higgs boson of mass around 120 GeV, the major decay mode is into $b\bar{b}$. The partial decay rate into $WW$ will begin to grow at $m_h \gtrsim 140$ GeV. Therefore, we shall focus on the $b\bar{b}$ mode for the Higgs boson decay. A light radion, on the other hand, has the major decay mode of $gg$ because of the QCD trace anomaly, followed by $b\bar{b}$ (a distant second). When the radion mass gets above the $WW$ threshold, the $WW$ mode becomes dominant.

The major background comes from the QCD heavy-flavor production of

$$\gamma\gamma \to b\bar{b}/c\bar{c} + 2 \text{ jets}, \quad (20)$$

where each jet can come from a gluon or a light quark. Here the $c\bar{c}$ pair can also fake the $B$-tag though with a much lower probability than the $b$ quark. We calculate the QCD $b\bar{b}/c\bar{c} + 2$ jets background by a parton-level calculation, in which the subprocesses are generated by MADGRAPH [21]. Another possible source of background is

$$\gamma\gamma \to W^+W^-, \quad (21)$$

followed by the hadronic decay of the $WW$ pair. The decay of the $W$ boson into a $b$ quark is severely suppressed by $(V_{cb} + V_{ub})^2 \simeq (0.05)^2$. The chance of seeing two $b$ quarks in the $WW$ decay is very rare, of the order of $6 \times 10^{-6}$. On the other hand, $WW$ production is still a possible background because of the $W$ boson decay into a $c$ quark. The charm quark may be tagged with a displaced vertex with a small mistag probability, thus may be misinterpreted as a $b$ quark.

We assume a 50% $B$-tagging efficiency and a chance of 5% mistag (a charm quark misinterpreted as a $b$ quark) in our study. Typical cuts on detecting the $b$-jets and light jets are applied:

$$p_T(b) > 15 \text{ GeV}, \quad p_T(j) > 15 \text{ GeV},$$
\[ |\cos \theta_b| < 0.9, \quad |\cos \theta_j| < 0.9, \quad \cos \theta_{b,j}, \cos \theta_{j,j}, \cos \theta_{b,j} < 0.9, \]

where \(\cos \theta_b\) and \(\cos \theta_j\) denote the cosine of the angle of the outgoing \(b\) quark and the jet, respectively, and \(\cos \theta_{b,j}, \cos \theta_{j,j}, \cos \theta_{b,j}\) denote the cosine of the angle between the two \(b\) quarks, between the two jets, and between the \(b\) quark and the jet, respectively. The angular cuts are mainly for the detection purpose and for removing the collinear divergence in the calculation. We have applied a gaussian smearing \(\Delta E/E = 0.5/\sqrt{E}\), where \(E\) is in GeV, to the final-state \(b\)-jets and light jets to simulate the detector resolution. Since the Higgs boson is produced together with a radion mainly via an intermediate graviton KK state, the Higgs boson tends to have a large \(p_T \sim m_G^{(1)}/2\). Therefore, a transverse momentum cut on the \(b\bar{b}\) pair is very efficient against the QCD background while only hurts the signal marginally. We apply a cut

\[ p_T(b\bar{b}) > 100 \text{ GeV} , \quad (22) \]

to reduce the background. Finally, we apply the invariant mass constraint of the \(b\bar{b}\) pair to be near the Higgs boson mass and that of the jet pair to be near the radion mass:

\[ |m_{b\bar{b}} - m_H| < 10 \text{ GeV} , \quad |m_{jj} - m_\phi| < 10 \text{ GeV} . \quad (23) \]

We summarize the cross sections for the signal and various backgrounds under successive cuts in Table I. The final signal-to-background ratio is quite promising. For \(m_\phi = 30 \text{ GeV}\) and \(\xi = 1.4\) we obtain a signal-to-background ratio about 1.3 : 1. For \(m_\phi = 70 \text{ GeV}\) with \(\xi = 0.6\) a ratio of 1.4 : 1 can be obtained. A ratio of 0.37 : 1 is obtained for the case of \(m_\phi = 150 \text{ GeV}\) and \(\xi = 0.25\). Note that the signal cross section scales as \(\xi^2\) as long as the positive mass-square constraints for the scalar bosons are satisfied.

In Table II we show the results at \(\sqrt{s_{ee}} = 0.8 \text{ TeV}\), at which the effect of the first graviton resonance is large, such that the signal cross section is substantially larger than the background after the cuts. This is because the majority of the photon collisions are at \(\sqrt{s} = \sqrt{x_1x_2s_{ee}} \sim (0.7 - 0.8)(800 \text{GeV})\), which is very close to the first KK graviton resonance.

To disentangle possible conflict between the electroweak precision data and the direct search bound of the SM Higgs, we studied the possibility that the SM Higgs mixes with the radion of the RS model such that the Higgs branching ratio into \(b\bar{b}\) becomes smaller and thus escapes the limit of direct search. We have explored the direct search of the radion-Higgs...
TABLE I: Cross sections in fb under successive application of the cuts mentioned in the text at $\sqrt{s_{ee}} = 0.5$ TeV. Unpolarized beams are chosen. The results for $m_\phi = 30$ ($\xi = 1.4$), 70 ($\xi = 0.6$) and 150 ($\xi = 0.25$) GeV are shown in the corresponding rows.

| Cuts | $h\phi$ signal | $b\bar{b}jj$ | $c\bar{c}jj$ | $W^+W^- \rightarrow c\bar{c}jj$ |
|------|---------------|--------------|--------------|-------------------------------|
| $p_T(b,j) > 15$ GeV | 0.14 | (0.16) | 39 | 346 | 287 |
| $|\cos\theta_{b\bar{b},j}| < 0.9$ | (0.16) | 39 | 346 | 287 |
| $\cos\theta_{b\bar{b},j} < 0.9$ | (0.018) | 39 | 346 | 287 |
| additional | 0.14 | (0.092) | 1.32 | 0.0009 |
| $p_T(b\bar{b}) > 100$ GeV | (0.15) | 4.5 | 38 | 26 |
| additional | 0.14 | (0.091) | 1.0 | (0.242) |
| $|M_{b\bar{b}} - m_H| < 10$ GeV | (0.15) | (0.091) | (1.0) | (0.242) |
| $|M_{jj} - m_\phi| < 10$ GeV | (0.014) | (0.027) | (0.215) | (0.82) |
| with B-tag $\epsilon_b = 0.5$ | 0.034 | 0.023 | 0.0033 | $2 \times 10^{-6}$ |
| or C-mistag $\epsilon_c = 0.05$ | (0.037) | (0.023) | (0.0025) | $(6 \times 10^{-4})$ |
| | (0.0035) | (0.007) | $(5 \times 10^{-4})$ | $(2 \times 10^{-3})$ |

associated production at photon colliders, and shown that the photon colliders achieved by the laser backscattering technique are very special in probing such a process. As is well known, the advantage of photon colliders is the capability of adjusting the polarization of incoming photons such that the signal of the spin-2 graviton exchange can be largely enhanced compared to the unpolarized collision. We have found the enhancement factor is around five, unless $\xi$ is quite small. This can be attained by aligning the polarization of the incoming photons such that the nature of the spin-2 graviton exchange is fully enhanced. We have also studied the corresponding background to $h\phi$ pair productions at photon colliders step-by-step in detail. Specifically, we have considered the major backgrounds from the QCD heavy flavor production of $\gamma\gamma \rightarrow b\bar{b}/c\bar{c} + 2$ jets and $\gamma\gamma \rightarrow W^+W^- \rightarrow c\bar{c} + 2$ jets. By imposing the various cuts, we have shown that the associated production signal of $h\phi$ can be comparable or even much larger than those backgrounds at photon colliders.
TABLE II: Cross sections in fb under successive application of the cuts mentioned in the text at √s_{ee} = 0.8 TeV. Unpolarized beams are chosen. The mixing parameter ξ = 1.4 and m_φ = 30 GeV are chosen.

| Cuts                      | hφ signal | b̄bjj | cĉjj | W^+W^- → cĉjj |
|---------------------------|-----------|-------|------|---------------|
| p_T(b, j) > 15 GeV       |           |       |      |               |
| | cos θ_{b̄bjj} | 6.25     | 31    | 280  | 156           |
| | cos θ_{b̄b,j}  |           |       |      |               |
| additional               |           |       |      |               |
| p_T(bb) > 200 GeV        |           |       |      |               |
| | additional         |           |       |      |               |
| | M_φ - m_H | 5.95     | 0.0015| 0.031| ~ 0           |
| | M_\bar{H} - m_Φ |           |       |      |               |
| with B-tag ϵ_b = 0.5     |           |       |      |               |
| or C-mistag ϵ_c = 0.05   | 1.49      | 0.0004| 8 × 10^{-5}| ~ 0 |

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REFERENCES

[1] See, e.g., J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, *The Higgs Hunter’s Guide* (Addison-Wesley, Reading, MA, 1990).

[2] G. Abbiendi *et al.*, The LEP Working Group for Higgs Boson Searches, Phys. Lett. **B565**, 61 (2003).

[3] “Electroweak Measurements”, talk by C. Diaconu at the Lepton-Photon 2005, June 30 – July 5, Uppsala, Sweden.
[4] L. Randall, R. Sundrum, Phys. Rev. Lett. 83, 3370 (1999); L. Randall, R. Sundrum, Phys. Rev. Lett. 83, 4690 (1999).

[5] W. D. Goldberger and M. B. Wise, Phys. Rev. Lett. 83, 4922 (1999); W. D. Goldberger and M. B. Wise, Phys. Lett. B475, 275 (2000).

[6] C. Csaki, M. Graesser, L. Randall and J. Terning, Phys. Rev. D62, 045015 (2000).

[7] S. B. Bae, P. Ko, H. S. Lee and J. Lee, Phys. Lett. B487, 299 (2000).

[8] G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B595, 250 (2001).

[9] C. Csaki, M. L. Graesser and G. D. Kribs, Phys. Rev. D63, 065002 (2001); C. S. Kim, J. D. Kim and Jeonghyeong Song, Phys. Lett. B511, 251 (2001); C. S. Kim, J. D. Kim and Jeong-hyeon Song, Phys. Rev. D67, 015001 (2003).

[10] U. Mahanta and S. Rakshit, Phys. Lett. B480, 176 (2000); K. Cheung, Phys. Rev. D63, 056007 (2001); U. Mahanta and A. Datta, Phys. Lett. B483, 196 (2000); S. C. Park, H. S. Song and J. Song, Phys. Rev. D65, 075008 (2002); S. C. Park, H. S. Song and J. Song, Phys. Rev. D63, 077701 (2001); C. S. Kim, Kang Yoon Lee and Jeonghyeong Song, Phys. Rev. D64, 015009 (2001).

[11] M. Chaichian, A. Datta, K. Huitu and Z. Yu, Phys. Lett. B524, 161 (2002).

[12] T. Han, G. D. Kribs and B. McElrath, Phys. Rev. D64, 076003 (2001).

[13] J. L. Hewett and T. G. Rizzo, JHEP 0308, 028 (2003) arXiv:hep-ph/0202155.

[14] D. Dominici, B. Grzadkowski, J. F. Gunion and M. Toharia, Nucl. Phys. B 671, 243 (2003) arXiv:hep-ph/0206192.

[15] J. F. Gunion, M. Toharia and J. D. Wells, Phys. Lett. B 585, 295 (2004) arXiv:hep-ph/0311219.

[16] K. Cheung, C.S. Kim, and J.-H. Song, Phys. Rev. D67, 075017 (2003).

[17] K. Cheung, C.S. Kim, and J.-H. Song, Phys. Rev. D69, 075011 (2004).
[18] I.F. Ginzburg, G.L. Kotkin, V.G. Serbo, and V.I. Telnov, Nucl. Instr. and Meth. 205, 47 (1983); I.F. Ginzburg, G.L. Kotkin, S.L. Panfil, V.G. Serbo, and V.I. Telnov, Nucl. Instr. and Meth. 219, 5 (1984); V.I. Telnov, hep-ex/9908005, hep-ex/9910010.

[19] H. Davoudiasl, J. L. Hewett and T. G. Rizzo, Phys. Lett. B473, 43 (2000).

[20] H. Davoudiasl, J. L. Hewett and T. G. Rizzo, Phys. Rev. D63, 075004 (2001).

[21] T. Stelzer and W.F. Long, Comput. Phys. Commun. 81, 357 (1994); F. Maltoni and T. Stelzer, JHEP 0302, 027 (2003).