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Direct bound on the minimal universal extra dimension model from the $t\bar{t}$ resonance search at the Tevatron

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In the minimal universal extra dimension model, the second Kaluza-Klein gluon $g^{(2)}$ has loop-induced vertices with the standard model quarks, mediated by the first Kaluza-Klein modes of the quark and the gluon/Higgs boson. With a top quark pair, this vertex is enhanced by the cooperation of strong coupling of a gluon and a large Yukawa coupling of a top quark, leading to substantial branching ratio of $\mathrm{BR}(g^{(2)} \to t\bar{t}) \approx$ 7%–8%. As the $g^{(2)}$ coupling with two gluons appears via dimension-6 operators, the process $q\bar{q} \to g^{(2)} \to t\bar{t}$ is very efficient to probe the minimal universal extra dimension model. The recent Tevatron data on the search for the $t\bar{t}$ resonance at $\sqrt{s} = 1.96$ TeV with an integrated luminosity of 8.7 fb^{-1} are shown to give the direct bound on the $g^{(2)}$ mass above 800 GeV. The implication and future prospect at the LHC are discussed also.

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I. INTRODUCTION

As the heaviest known fundamental particle in the standard model (SM) [1], the top quark is phenomenologically more attractive because it can be tagged. If produced near the threshold, the top quark is identified with a b-tagged jet and a W boson, or three jets of which the invariant mass is near the top quark mass. If highly boosted, it is tagged by the substructure of collimated jets [2]. The top tagging opens new channels to new physics beyond the SM, especially through resonant $t\bar{t}$ production. The $t\bar{t}$ resonances have been searched at the Tevatron [3–6] and at the LHC [7,8]. The most recent results have been reported in the 36th International Conference on High Energy Physics (ICHEP 2012) by the ATLAS [9], CMS [10], and Tevatron experiments [11]. No experiment has found any significant evidence, which constrains various new physics models with $t\bar{t}$ resonances [12].

One interesting candidate for $t\bar{t}$ resonance is the second Kaluza-Klein (KK) gluon in the minimal universal extra dimension (mUED) model [13]. This model has an additional single flat extra dimension of size R, compactified on an S_1/Z_2 orbifold. All of the SM fields propagate freely in the whole five-dimensional (5D) spacetime, each of which has a tower of KK modes. A lot of interest has been drawn as it provides solutions for proton decay [14], the number of fermion generations [15], and supersymmetry breaking [16]. Most of all, the conservation of KK parity makes this model more appealing: the compactification scale R^{-1} can be as low as about 300 GeV; the lightest KK particle is a good candidate for the cold dark matter [17].

There are various *indirect* constraints on the lower bound on $R^{-1} \gtrsim 300$ GeV from the ρ parameter [18], the electroweak precision tests [19], the muon $(g-2)$ measurement [20], the flavor changing neutral currents [21], and the recent measurement of the Higgs boson mass 125 GeV [22]. An upper limit on $R^{-1} \lesssim 1.4$ TeV is from dark matter constraints to avoid overclosing the Universe [17,23]. Despite the intensive studies on collider signatures [24], direct search for the mUED model has generic difficulties because of nearly degenerate KK mass spectrum. The most accessible new particles, the first KK modes, end up with two missing the lightest KK particles and very soft SM particles, which is very challenging to measure at hadron colliders. If the mUED model is extended including nonvanishing fermion bulk mass μ, called the split mUED model [25], the second KK gauge boson has tree level vertices with the SM fermions. The split mUED model has been strongly constrained by the CMS experiment: $R^{-1} \gtrsim 800$ GeV for $\mu > 100$ GeV [26].

As a smoking gun signature of the mUED model at hadron colliders, we focus on the *second* KK gluon. The major decay modes are KK number conserving into $q^{(2)}q$ and $q^{(1)}\bar{q}^{(1)}$, but they are kinematically suppressed by nearly degenerate KK mass spectrum. Loop-induced decays into SM particles become important, which are mediated by the first KK modes of the quark and the gluon/Higgs boson [27,28]. In addition, large Yukawa coupling of the top quark enhances branching ratio of $g^{(2)} \to t\bar{t}$. This $g^{(2)}$ is a very good candidate for the $t\bar{t}$ resonance. As shall be shown, the recent Tevatron data on the resonant $t\bar{t}$ search have set a significant direct bound on the mUED model. This is our main result.

II. THE LOOP-INDUCED VERTEX OF $g^{(2)}$ IN THE MUED MODEL

The mUED model has one additional flat extra dimension with size R. The word *universal* is from the setup that

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the whole 5D spacetime is accessible to all of the SM fields. Each SM particle has an infinite tower of KK modes, effectively described by the KK expansion [28,29]. Chiral SM fermions from vector-like 5D fermions are achieved by the compactification of the extra dimension on an S^1/Z_2 orbifold: the zero mode fermion with wrong chirality is removed by imposing odd parity under the orbifold projection $y \to -y$, called the Z_2 parity.

The KK mass is of geometrical origin, which is at tree level

$$M_{KK}^{(n)} = \begin{cases} \sqrt{M_n^2 + m_0^2} & \text{(Boson);} \\ M_n + m_0 & \text{(Fermion),} \end{cases} \tag{1}$$

where $M_n = nR^{-1}$, n is the KK number, and m_0 is the corresponding SM particle mass. We have nearly degenerate KK mass spectrum with the given n. The radiative corrections to the KK masses [30] play a crucial role in the phenomenologies, determining whether a specific decay mode is kinematically allowed or not. In the mUED model where boundary kinetic terms vanish at the cutoff scale Λ, the radiative corrections to the KK masses are well defined and finite.

A KK mode decays. At tree level, the decay respects the conservation of KK number. However, the high degeneracy in the KK mass spectrum suppresses these decay modes because of small kinematic space. Kaluza-Klein number violating decays, which occur at one-loop level, can be significant. Note that KK parity $(-1)^n$ is still conserved at loop level.

For the phenomenological signatures of the second KK quark and gluon, it is important to note that some KK-parity conserving but KK-number violating couplings are negligible:

(1) The vertex of $q^{(2)}$-q-g is negligible. The four dimensional operator $\bar{q}^{(2)}\gamma^\mu \frac{\lambda^a}{2} q g_\mu^a$ vanishes because of the gauge invariance. The next higher dimensional operator $\bar{q}^{(2)}\sigma^{\mu\nu}\frac{\lambda^a}{2} q F_{\mu\nu}^a$ is suppressed by $1/\Lambda$.

(2) The vertex g-g-$g^{(2)}$ is suppressed by $1/\Lambda^2$ since it appears through dimension-6 operators [31,32].[1] In addition, the couplings do not have the logarithmic enhancement factor which is about 4.6–6.4 for $\Lambda R = 20$, 50. Dimension-4 operators that would generate the vertex vanish by the unbroken 4D gauge invariance and the absence of the kinetic and mass mixing between g and $g^{(2)}$ [32].

In what follows, therefore, we ignore these two kinds of vertices.

The loop-induced vertices of $g^{(2)}$ are only with the SM quarks, given by

$$\begin{aligned} &-i\frac{g_s}{\sqrt{2}}\sum_{X=L,R}\left(\frac{\bar{\delta} m_{g_2}^2}{M_2^2} - 2\frac{\bar{\delta} m_{q_{X2}}}{M_2}\right)\bar{f}_X\gamma^\mu\frac{\lambda^a}{2}P_X f_X g_\mu^{a(2)} \\ &\equiv -i\frac{g_s}{\sqrt{2}}\left(\frac{1}{16\pi^2}\ln\frac{\Lambda^2}{Q^2}\right)\bar{f}\gamma^\mu\frac{\lambda^a}{2}\{\hat{g}_{fL}P_L + \hat{g}_{fR}P_R\} f g_\mu^{a(2)}, \end{aligned} \tag{2}$$

where $P_{R,L} = (1 \pm \gamma^5)/2$, Q is the regularization scale, and $\bar{\delta} m$ is the boundary mass correction [30]. For the $g^{(2)}$ production, we adopt $Q = 2R^{-1}$. The effective couplings of $g^{(2)}$ with $t_{L,R}$ are

$$\begin{aligned} \hat{g}_{tL} &= \frac{1}{8}[44g_3^2 - 27g_2^2 - g_1^2 + 12h_t^2] \approx 6.2, \\ \hat{g}_{tR} &= \frac{11}{2}g_3^2 - 2g_1^2 + 3h_t^2 \approx 8.9. \end{aligned} \tag{3}$$

The coupling of $g^{(2)}$ with the left-handed (right-handed) light quark is the same as $\hat{g}_{tL}$ ($\hat{g}_{tR}$) except for the top Yukawa coupling, which is numerically 4.7 (6.1). As explicitly shown in Eq. (3), strong coupling of the gluon and large Yukawa coupling of the top quark play in the same direction to increase the branching ratio of the $g^{(2)}$ decay into $t\bar{t}$. We have BR($g^{(2)} \to t\bar{t}$) $\approx$ 7%–8%, depending on the model parameters. Note that the enhanced coupling of the KK gluon with top quarks occurs in the bulk Randall-Sundrum also, which constrains the model by the measured top quark cross section at the Tevaton [33].

III. DIRECT BOUNDS FROM THE $t\bar{t}$ RESONANCE SEARCH

The production of the second KK gluon is through $q\bar{q}$ annihilation. Even without large Yukawa couplings, light quarks also have sizable couplings with $g^{(2)}$. The best channel to probe second KK modes at a hadron collider is the $q\bar{q}$ annihilation production of $g^{(2)}$, followed by the decay $g^{(2)} \to t\bar{t}$. Although other second KK gauge bosons like $Z^{(2)}$ and $\gamma^{(2)}$ also produce $t\bar{t}$ resonance signal, their electroweak production is much smaller than the $g^{(2)}$ production by an order of magnitude [27].

The production of $g^{(2)}$ has additional production channels associated with soft jets. The heavy mass of $g^{(2)}$ and the steeply falling parton luminosities yield $g^{(2)}$ production near the threshold. The accompanied jet in the processes of $q\bar{q} \to gg^{(2)}$, $gq \to qg^{(2)}$, and $gg \to q\bar{q}g^{(2)}$ are soft. At a hadron collider, the number of jets are measured in terms of jet multiplicity. If jets are soft, however, they are very likely to be missed in the jet multiplicity. Soft jets tend to spread out. In addition, if the transverse momentum of soft jets are too low like below 20 GeV, soft jets cannot excite showers in the hadron calorimeter of the detector. Finally, the jets with $|\eta_j| > 2.5$, going out of the barrel and end caps of the hadron calorimeter, are also missed. Therefore,

[1] We thank Ayres Freitas for pointing this out.

we include the $g^{(2)}$ production with soft jets defined by $p_T^j < 20$ GeV or $|\eta_j| > 2.5$.

We first present the expected signal and the observed upper bound at the Tevatron. In Fig. 1, we show the production cross section of $g^{(2)}$ multiplied by $\mathrm{BR}(g^{(2)} \to t\bar{t})$, as a function of $g^{(2)}$ mass for $\Lambda R = 20, 30, 50$. For the signal event generator, we have used CalcHEP [34] with the CTEQ6 parton distribution function [35]. We set the renormalization μ_R and factorization μ_F scales to be equal as $\mu_R = \mu_F = M_{t\bar{t}}$. In order to estimate high order QCD corrections, we examine the scale variation of the cross section. When varying two scales μ_R and μ_F in the range between $M_{t\bar{t}}/2$ and $2M_{t\bar{t}}$, we obtain about 15% uncertainty in the production cross section of $p\bar{p} \to g^{(2)} \to t\bar{t}$ at the Tevatron. At the LHC, the magnitude of QCD scale uncertainties is similar. As shall be shown, this uncertainty is less than the variation of the cross section by changing the model parameter ΛR. Therefore, we consider only the median case of $\mu_R = \mu_F = M_{t\bar{t}}$ in what follows.

Larger ΛR yields larger signal, since the loop-induced vertices increase logarithmically with ΛR; see Eq. (2). The Tevatron search for $t\bar{t}$ resonance based on the 4.8 fb^{-1} data [6] has already set the lower bounds on $M_{g^{(2)}} \gtrsim 535$ GeV for $\Lambda R = 20$, $M_{g^{(2)}} \gtrsim 590$ GeV for $\Lambda R = 30$, and $M_{g^{(2)}} \gtrsim 710$ GeV for $\Lambda R = 50$. More recent Tevatron data, presented in ICHEP 2012, with total luminosity of 8.7 fb^{-1} [11] constrain further to be $M_{g^{(2)}} \gtrsim 800$ GeV. This bound applies for $\Lambda R = 20$–50. The Tevatron group presented their analysis only up to the $t\bar{t}$ invariant mass of 800 GeV. If naively extrapolating the observation, we have $M_{g^{(2)}} \gtrsim 820$ GeV for $\Lambda R = 20$, $M_{g^{(2)}} \gtrsim 870$ GeV for $\Lambda R = 30$, and $M_{g^{(2)}} \gtrsim 920$ GeV for $\Lambda R = 50$. These are very significant direct bounds on $M_{g^{(2)}}$ and thus the mUED model.

Figure 2 shows the expected signal in the mUED model and the observed 95% C.L. upper bound at the LHC with $\sqrt{s} = 7$ TeV. We present the 200 pb^{-1} [7] and 2.05 fb^{-1} data [9] by the ATLAS experiment, and 4.6 fb^{-1} [8] and 5.0 fb^{-1} data [10] by the CMS experiment. Since we have included only the $q\bar{q}$ annihilation production of $g^{(2)}$, this is a conservative limit. The gluon fusion production, though negligible from dimension-6 operators, would increase the signal. The signal is still quite below the upper bound set by the ATLAS and CMS experiments.

Finally, we suggest another efficient observable based on a unique feature of $g^{(2)}$. The dependence of the model parameters $(R^{-1}, \Lambda R)$ on the $g^{(2)}$-q-$\bar{q}$ vertex in Eq. (2) is common for all SM quarks. The ratio of the $g^{(2)}$ decay rate into $t\bar{t}$ to that into $q\bar{q}$ is almost fixed by the SM parameters. Minor dependence exists through the mass of $g^{(2)}$, which affects kinematics.

In Fig. 3, we show the ratio $\mathrm{BR}(g^{(2)} \to jj)/\mathrm{BR}(g^{(2)} \to t\bar{t})$ as a function of $M_{g^{(2)}}$. Here the dijet signal includes all 5 light quarks. Two lines for $\Lambda R = 20, 50$ are almost identical as expected. In addition, the value of this ratio is large about 3.5. This is attributed to the number of flavors although light quarks have smaller effective couplings with $g^{(2)}$ than the top quark. If $g^{(2)}$ is observed as a $t\bar{t}$ resonance, we should see the same resonance in the dijet channel with about 3.5 times larger rate. This is one of the

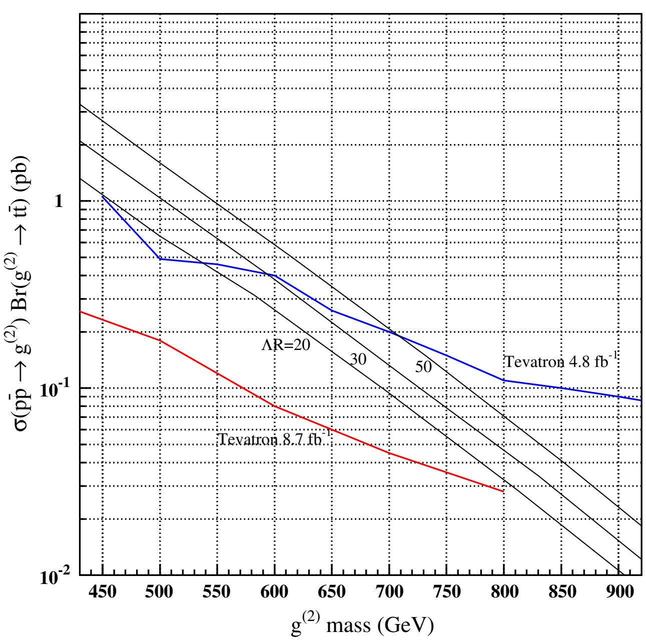

FIG. 1 (color online). Expected and observed 95% C.L. upper limit on $\sigma(p\bar{p} \to g^{(2)} \to t\bar{t})$ as a function of $t\bar{t}$ invariant mass at the Tevatron with $\sqrt{s} = 1.96$ TeV.

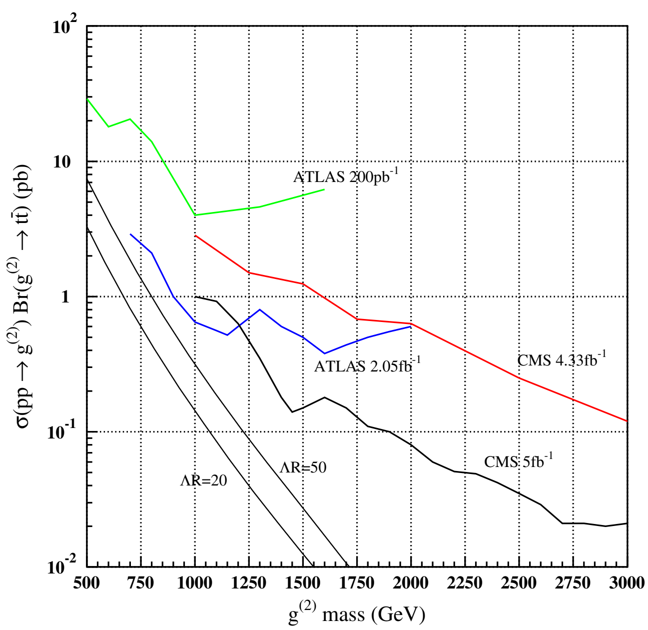

FIG. 2 (color online). Expected and observed 95% C.L. upper limit on $\sigma(pp \to g^{(2)} \to t\bar{t})$ as a function of $t\bar{t}$ invariant mass at the LHC with $\sqrt{s} = 7$ TeV.

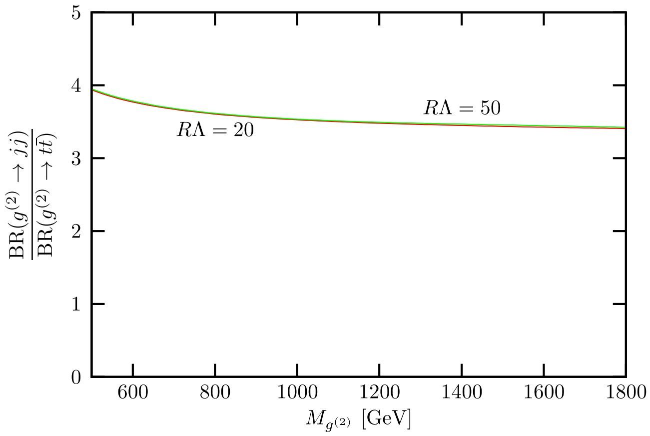

FIG. 3 (color online). The ratio $\mathrm{BR}(g^{(2)} \to jj)/\mathrm{BR}(g^{(2)} \to t\bar{t})$ as a function of the $g^{(2)}$ mass for $\Lambda R = 20, 50$. Here j denotes the light quark such that $\mathrm{BR}(g^{(2)} \to jj) = \sum_{q=u,d,c,s,b} \mathrm{BR}(g^{(2)} \to q\bar{q})$.

most powerful signals of the mUED model. Recently, the CMS and ATLAS experiments have reported their search for dijet resonance [36]. The current upper bounds are too weak to constrain the mUED model yet. With a large data set, the future prospect at the LHC through the correlation between the $t\bar{t}$ and dijet channels is very promising.

IV. CONCLUSIONS

Despite of its various theoretical virtues, the mUED model is one of the most elusive models for hadron colliders. The nearly degenerate KK mass spectrum makes the signals of the first KK modes difficult to probe, each of which decays into a very soft SM particle with missing transverse energy. Turning our attention to the second KK mode, we have found one efficient process, $p\bar{p} \to q\bar{q} \to g^{(2)} \to t\bar{t}$. Strong coupling of a gluon and large Yukawa coupling of a top quark cooperate to enhance the branching ratio of $g^{(2)} \to t\bar{t}$. The vertex g-g-$g^{(2)}$ appears from dimension-6 operators, which leads to the main $g^{(2)}$ production through $q\bar{q}$ annihilation.

We have shown that the recent Tevatron $t\bar{t}$ search with 8.7 fb^{-1} data has set very significant direct bound on the mUED model. The $g^{(2)}$ mass below 800 GeV is excluded for all model parameters. At the LHC, the suppression of $g^{(2)}$ gluon fusion production reduces the sensitivity. No direct bounds from the resonant $t\bar{t}$ data have been derived yet. However, the suggested correlation between the $t\bar{t}$ resonance and the dijet resonance is expected to provide a good probe of the model.

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