Dijet Signals for Low Mass Strings at the Large Hadron Collider

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Assuming that the fundamental string mass scale is in the TeV range and the theory is weakly coupled, we discuss possible signals of string physics at the Large Hadron Collider (LHC). In D-brane constructions, the dominant contributions to full-fledged string amplitudes for all the common QCD parton subprocesses leading to dijets are completely independent of the details of compactification, and can be evaluated in a parameter-free manner. We make use of these amplitudes evaluated near the first resonant pole to determine the discovery potential of LHC for the first Regge excitations of the quark and gluon. Remarkably, the reach of LHC after a few years of running can be as high as 6.8 TeV. Even after the first 100 pb^{-1} of integrated luminosity, string scales as high as 4.0 TeV can be discovered. Data on $pp \to$ direct$\gamma+$ jet can provide corroboration for string physics at scales as high as 5 TeV.

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Experiments at the Large Hadron Collider (LHC) will search for the Higgs boson—the last missing piece of the standard model (SM). At the same time, searches will be conducted for signals of new physics beyond it. Although many extensions of the SM have been proposed over the last 30 years, none of them has been favored so far by the experimental data. There are differences, however, between various models with respect to their testability in the energy range accessible at the LHC. It has recently become clear [1–4] that superstring theory offers one of the most robust scenarios beyond the SM, provided that its fundamental mass scale is sufficiently "low", i.e., of order TeV. The purpose of this Letter is to discuss dijet signals of low mass string theory at the LHC and to determine the discovery reach based on the measurements of the corresponding cross sections.

The mass scale M_s of fundamental strings can be as low as few TeV provided that spacetime extends into large extra dimensions, allowing a novel solution for the hierarchy problem [5,6]. This mass determines the center of mass energy threshold $\sqrt{s} \geq M_s$ for the production of Regge resonances in parton collisions, thus for the onset of string effects at the LHC [4]. (Mandelstam variables s, t, u used in this Letter refer to parton subprocesses.) We consider the extensions of the SM based on open strings ending on D-branes, with gauge bosons due to strings attached to stacks of D-branes and chiral matter due to strings stretching between intersecting D-branes [7]. Only one assumption is necessary in order to set up a solid framework: the string coupling must be small in order to rely on perturbation theory in the computations of scattering amplitudes. In this case, black hole production and other strong gravity effects occur at energies above the string scale; therefore at least a few lowest Regge recurrences are available for examination, free from interference with some complex quantum gravitational phenomena. Starting from a small string coupling, the values of the SM coupling constants are determined by D-brane configurations and the properties of extra dimensions; hence, that part of superstring theory requires intricate model building; however, as argued in Refs. [1–4], some basic properties of Regge resonances like their production rates and decay widths are completely model-independent. The resonant character of parton cross sections should be easy to detect at the LHC if the string mass scale is not too high. Direct photon production channels discussed in [1,2] are very interesting because the signal is due to processes that are absent in the SM at the leading (tree level) order of perturbation theory. On the other hand, dijet production is a standard tree-level QCD process. However, as argued in this Letter, the presence of string resonances may lead to a dramatic enhancement of the production rates and allow access to higher string mass scales. In addition, dijet calculations do not depend on the unknown mixing parameter characterizing the embedding of hypercharge in the extended D-brane gauge group. Moreover, as discussed later in this work, future measurements of dijet angular distributions can provide a potent method for distinguishing between various compactification scenarios.

The physical processes underlying dijet production at the LHC are the collisions of two partons, producing two final partons that fragment into hadronic jets. The corresponding $2 \to 2$ scattering amplitudes, computed at the leading order in string perturbation theory, are collected in Ref. [4]. The amplitudes involving four gluons as well as those with two gluons plus two quarks do not depend on the compactification details and are completely model-independent. All string effects are encapsulated in these amplitudes in one "form factor" function of Mandelstam variables s, t, u (constrained by $s + t + u = 0$):

$$V(s,t,u) = \frac{su}{tM_s^2} B(-s/M_s^2, -u/M_s^2) = \frac{\Gamma(1-s/M_s^2)\Gamma(1-u/M_s^2)}{\Gamma(1+t/M_s^2)}. \quad (1)$$

The physical content of the form factor becomes clear after using the well-known expansion in terms of s-channel resonances [8]:

$$B(-s/M_s^2, -u/M_s^2) = -\sum_{n=0}^{\infty} \frac{M_s^{2-2n}}{n!} \frac{1}{s - nM_s^2} \times \left[\prod_{J=1}^{n} (u + M_s^2 J)\right], \quad (2)$$

which exhibits s-channel poles associated to the propagation of virtual Regge excitations with masses $\sqrt{n}M_s$. Thus near the nth level pole ($s \to nM_s^2$):

$$V(s,t,u) \approx \frac{1}{s - nM_s^2} \frac{M_s^{2-2n}}{(n-1)!} \prod_{J=0}^{n-1} (u + M_s^2 J). \quad (3)$$

In specific amplitudes, the residues combine with the remaining kinematic factors, reflecting the spin content of particles exchanged in the s channel, ranging from $J = 0$ to $J = n + 1$.

The amplitudes for the four-fermion processes like quark-antiquark scattering are more complicated because the respective form factors describe not only the exchanges of Regge states but also of heavy Kaluza-Klein and winding states with a model-dependent spectrum determined by the geometry of extra dimensions. Fortunately, they are suppressed, for two reasons. First, the QCD $SU(3)$ color group factors favor gluons over quarks in the initial state. Second, the parton luminosities in proton-proton collisions at the LHC, at the parton center of mass energies above 1 TeV, are significantly lower for quark-antiquark subprocesses than for gluon-gluon and gluon-quark [2]. The collisions of valence quarks occur at higher luminosity; however, there are no Regge recurrences appearing in the s-channel of quark-quark scattering [4].

Before proceeding, we pause to present our notation. The first Regge excitations of the gluon (g) and quarks (q) will be denoted by g^*, q^*, respectively. In the D-brane models under consideration, the ordinary $SU(3)$ color gauge symmetry is extended to $U(3)$, so that the open strings terminating on the stack of "color" branes contain an additional $U(1)$ gauge boson C and its excitations to accompany the gluon and its excitations. The first excitation of the C will be denoted by C^*.

In the following we isolate the contribution to the partonic cross section from the first resonant state. Note that far below the string threshold, at partonic center of mass energies $\sqrt{s} \ll M_s$, the form factor $V(s,t,u) \approx 1 - \frac{\pi^2}{6} su/M_s^4$ [4] and therefore the contributions of Regge excitations are strongly suppressed. The s-channel pole terms of the average square amplitudes contributing to dijet production at the LHC can be obtained from the general formulae given in Ref. [4], using Eq. (3). However, for phenomenological purposes, the poles need to be softened to a Breit-Wigner form by obtaining and utilizing the correct *total* widths of the resonances [3]. After this is done, the contributions of the various channels are as follows:

$$|\mathcal{M}(gg \to gg)|^2 = \frac{19}{12}\frac{g^4}{M_s^4}\left\{W_{g^*}^{gg\to gg}\left[\frac{M_s^8}{(s-M_s^2)^2 + (\Gamma_{g^*}^{J=0} M_s)^2} + \frac{t^4+u^4}{(s-M_s^2)^2 + (\Gamma_{g^*}^{J=2} M_s)^2}\right] + W_{C^*}^{gg\to gg}\left[\frac{M_s^8}{(s-M_s^2)^2 + (\Gamma_{C^*}^{J=0} M_s)^2} + \frac{t^4+u^4}{(s-M_s^2)^2 + (\Gamma_{C^*}^{J=2} M_s)^2}\right]\right\}, \quad (4)$$

$$|\mathcal{M}(gg \to q\bar{q})|^2 = \frac{21}{12}\frac{g^4}{M_s^4}\left[W_{g^*}^{gg\to q\bar{q}} \frac{ut(u^2+t^2)}{(s-M_s^2)^2 + (\Gamma_{g^*}^{J=2} M_s)^2} + W_{C^*}^{gg\to q\bar{q}} \frac{ut(u^2+t^2)}{(s-M_s^2)^2 + (\Gamma_{C^*}^{J=2} M_s)^2}\right], \quad (5)$$

$$|\mathcal{M}(q\bar{q} \to gg)|^2 = \frac{56}{27}\frac{g^4}{M_s^4}\left[W_{g^*}^{q\bar{q}\to gg} \frac{ut(u^2+t^2)}{(s-M_s^2)^2 + (\Gamma_{g^*}^{J=2} M_s)^2} + W_{C^*}^{q\bar{q}\to gg} \frac{ut(u^2+t^2)}{(s-M_s^2)^2 + (\Gamma_{C^*}^{J=2} M_s)^2}\right], \quad (6)$$

$$|\mathcal{M}(qg \to qg)|^2 = -\frac{4}{9}\frac{g^4}{M_s^2}\left[\frac{M_s^4 u}{(s-M_s^2)^2 + (\Gamma_{q^*}^{J=1/2} M_s)^2} + \frac{u^3}{(s-M_s^2)^2 + (\Gamma_{q^*}^{J=3/2} M_s)^2}\right], \quad (7)$$

where g is the QCD coupling constant ($\alpha_{\rm QCD} = \frac{g^2}{4\pi} \approx 0.1$) and $\Gamma^J_{(g^*,q^*,C^*)}$ are total decay widths for intermediate states g^*, C^*, and q^* (with angular momentum J) [3]. The associated weights of these intermediate states $W^{ij\to jk}_{(g^*,C^*,q^*)}$ are computed using the probabilities for the various entrance and exit channels [3].

The resonance would be visible in data binned according to the invariant mass M of the dijet, after setting cuts on the different jet rapidities, $|y_1|, |y_2| \leq 1$ [9] and transverse momenta $p_T^{1,2} > 50$ GeV. In Fig. 1 we show a representative plot of the invariant mass spectrum, for $M_s = 2$ TeV, detailing the contribution of each subprocess. The QCD background has been calculated at the partonic level from the same processes as designated for the signal, with the

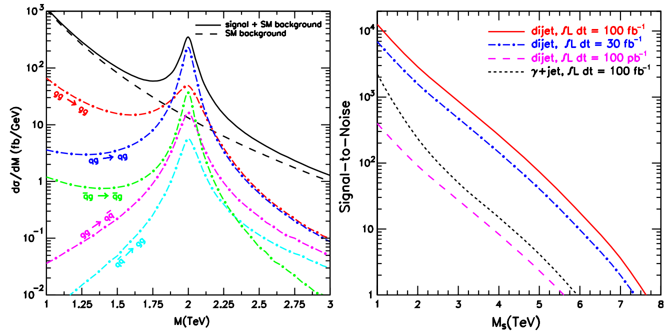

FIG. 1 (color online). Left: $d\sigma/dM$ (units of fb/GeV) vs M (TeV) is plotted for the case of SM QCD background (dashed line) and (first resonance) string signal + background (solid line). The dot-dashed lines indicate the different contributions to the string signal. Right: $pp \to$ dijet signal-to-noise ratio for three integrated luminosities. For comparison, we also show the signal-to-noise of $pp \to \gamma +$ jet, for $\kappa^2 \simeq 0.02$; see Ref. [1].

addition of $qq \to qq$ and $q\bar{q} \to q\bar{q}$. Our calculation, making use of the CTEQ6D parton distribution functions [10] agrees with that presented in [9].

We now estimate (at the parton level) the LHC discovery reach. Standard bump-hunting methods, such as obtaining cumulative cross sections, $\sigma(M_0) = \int_{M_0}^{\infty} \frac{d\sigma}{dM} dM$, from the data and searching for regions with significant deviations from the QCD background, may reveal an interval of M suspected of containing a bump. With the establishment of such a region, one may calculate a signal-to-noise ratio, with the signal rate estimated in the invariant mass window $[M_s - 2\Gamma, M_s + 2\Gamma]$. The noise is defined as the square root of the number of background events in the same dijet mass interval for the same integrated luminosity.

The top two and bottom curves in Fig. 1 show the behavior of the signal-to-noise (S/N) ratio as a function of M_s for three integrated luminosities (100 fb^{-1}, 30 fb^{-1} and 100 pb^{-1}) at the LHC. It is remarkable that within 1–2 years of data collection, string scales as large as 6.8 TeV are open to discovery at the $\geq 5\sigma$ level. For 30 fb^{-1}, the presence of a resonant state with mass as large as 5.7 TeV can provide a signal of convincing significance ($S/N = 592/36 > 13$). The bottom curve, corresponding to data collected in a very early run of 100 pb^{-1}, shows that a resonant mass as large as 4.0 TeV can be observed with 10σ significance. Once more, we stress that these results contain no unknown parameters. They depend only on the D-brane construct for the SM, and are independent of compactification details.

For comparison with our previous analysis, we also show in Fig. 1 a fourth curve, for the process $pp \to \gamma +$ jet. (In what follows, γ refers to an isolated gamma ray.) In Ref. [2] a cut ($p_T^\gamma > 300$ GeV) was selected for discovery of new physics. As far as the signal is concerned, this cut is largely equivalent to selecting on γ-jet invariant masses in the 2–5 TeV range, with cuts on photon and jet rapidities $|y_1|, |y_2| < 2.4$ [11]. However, for $M_s > 2$ TeV the background is greatly reduced with the dijet mass method used here, resulting in an extension of the discovery reach, up to about 5 TeV [12]. The signal used to obtain the results displayed in Fig. 1 includes the parton subprocesses $gg \to g\gamma$ (which does not exist at tree level in QCD, and which was the only subprocess evaluated in [1,2]), $qg \to q\gamma$, $\bar{q}g \to \bar{q}\gamma$, and $q\bar{q} \to g\gamma$. All except the first have been calculated in QCD and constitute the SM background. The projection of the photon onto the C gauge boson was also effected in the last-cited references. Although the discovery reach is not as high as that for dijets, the measurement of $pp \to \gamma +$ jet can potentially provide an interesting corroboration for the stringy origin for new physics manifest as a resonant structure in LHC data.

We now turn to the analysis of the angular distributions. QCD parton-parton cross sections are dominated by t-channel exchanges that produce dijet angular distributions which peak at small center of mass scattering angles. In contrast, nonstandard contact interactions or excitations of resonances result in a more isotropic distribution. In terms of rapidity variables for standard transverse momentum cuts, dijets resulting from QCD processes will preferentially populate the large rapidity region, while the new processes generate events more uniformly distributed in the entire rapidity region. To analyze the details of the rapidity space the D0 Collaboration [13] introduced a new parameter R, the ratio of the number of events, in a given dijet mass bin, for both rapidities $|y_1|, |y_2| < 0.5$ and both rapidities $0.5 < |y_1|, |y_2| < 1.0$ [14]. In Fig. 2 we compare the results from a full CMS detector simulation of the ratio R [15], with predictions from LO QCD and model-independent contributions to the q^*, g^*, and C^* excitations. For an integrated luminosity of 10 fb^{-1} the LO QCD contributions with $\alpha_{\rm QCD} = 0.1$ (corresponding to running scale $\mu \approx M_s$) are within statistical fluctuations of the full

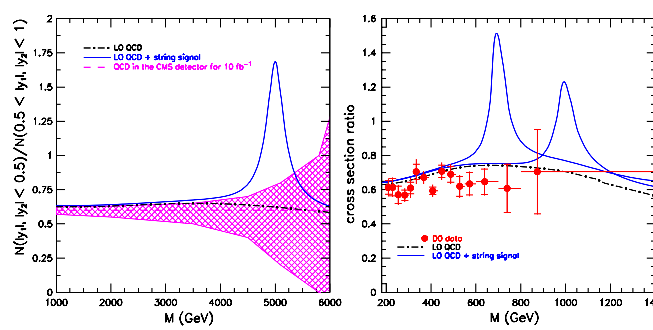

FIG. 2 (color online). Left: For a luminosity of 10 fb^{-1}, the expected value (dashed line) and statistical error (shaded region) of the dijet ratio of QCD in the CMS detector [15] is compared with LO QCD (dot-dashed line) and LO QCD plus lowest massive string excitation. Right: Ratio R of dijet invariant mass cross sections. The experimental points (solid circles) reported by the D0 Collaboration [13] are compared to a LO QCD calculation indicated by a dot-dashed line. The ratio R of the string + LO QCD invariant mass cross section is also shown as a solid line.

CMS detector simulation [16]. Since one of the purposes of utilizing NLO calculations is to fix the choice of the running coupling, we take this agreement as rationale to omit loops in QCD and in string theory. It is clear from Fig. 2 that incorporating NLO calculation of the background and the signal would not significantly change the large deviation of the string contribution from the QCD background. It is noteworthy that although not included in the present analysis, the signal due to $qq \to qq$, with Regge recurrences and Kaluza-Klein excitations exchanged in the t- and u-channels, could yield a small departure from the QCD value of R outside the resonant region. In an optimistic scenario, measurements of this modification could shed light on the D-brane structure of the compact space. It could also serve to differentiate between a stringy origin for the resonance as opposed to an isolated structure such as a Z', which would not modify R outside the resonant region [17].

Before closing, we discuss the impact of Tevatron data on generic D-brane constructions. To do so, we compute the cross section ratio, in a given dijet mass bin, for both rapidities satisfying $|y_1|, |y_2| < 0.5$ and $0.5 < |y_1|, |y_2| < 1.0$, respectively. We use these results to calculate the ratio R, which is shown in Fig. 2. We can essentially read off from the figure that values of $M_s < 700$ GeV are excluded at an extremely high C.L. Unfortunately, the dijet mass resolution is not enough to clear up the resonant structure for larger M_s. For example, at first sight the string prediction for $M_s = 1$ TeV seems to depart from the data in the last dijet mass bin. To ascertain any possible deviation, we integrated the cross section over the last inavriant mass bin [$\sigma(0 < |y_1|, |y_2| < 0.5) \simeq 125$ fb, $\sigma(0.5 < |y_1|, |y_2| < 1) \simeq 147$ fb] and estimated an average value of $R = 0.86$ that is in agreement with D0 data [13] at the 1σ level. One should keep in mind that at the Tevatron the valence quark contribution $q\bar{q} \to q\bar{q}$ is presumably an important component of the signal. However, in spite of its importance, we are unable to include this contribution without invoking a specific compactification scenario [18].

In D-brane constructions, the full-fledged string amplitudes supplying the dominant contributions to dijet cross sections are completely independent of the details of compactification. If the string scale is in the TeV range, such extensions of the SM can be of immediate phenomenological interest. In this Letter we have made use of the amplitudes evaluated near the first resonant pole to determine the discovery potential at LHC for the first Regge excitations of the quark and gluon. Remarkably, the reach of LHC after a few years of running can be as high as 6.8 TeV (with $S/N = 210/42$) [19]. Even after the first 100 pb^{-1} of integrated luminosity, string scales as high as 4.0 TeV can be discovered with $S/N = 55/6$.

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