Jet Signals for Low Mass Strings at the Large Hadron Collider

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The mass scale M_s of superstring theory is an arbitrary parameter that can be as low as few TeVs if the Universe contains large extra dimensions. We propose a search for the effects of Regge excitations of fundamental strings at the CERN Large Hadron Collider (LHC), in the process $pp \to \gamma + \text{jet}$. The underlying parton process is dominantly the single photon production in gluon fusion, $gg \to \gamma g$, with open string states propagating in intermediate channels. If the photon mixes with the gauge boson of the baryon number, which is a common feature of D-brane quivers, the amplitude appears already at the string disk level. It is completely determined by the mixing parameter—and it is otherwise model (compactification) independent. Even for relatively small mixing, 100 fb^{-1} of LHC data could probe deviations from standard model physics, at a 5σ significance, for M_s as large as 3.3 TeV.

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At the time of its formulation and for years thereafter, superstring theory was regarded as a unifying framework for Planck-scale quantum gravity and TeV-scale standard model (SM) physics. Important advances were fueled by the realization of the vital role played by D-branes [1] in connecting string theory to phenomenology [2]. This has permitted the formulation of string theories with compositeness setting in at TeV scales [3] and large extra dimensions. There are two paramount phenomenological consequences for TeV-scale D-brane string physics: the emergence of Regge recurrences at parton collision energies $\sqrt{\hat{s}} \sim$ string scale $\equiv M_s$; and the presence of one or more additional $U(1)$ gauge symmetries, beyond the $U(1)_Y$ of the SM. The latter follows from the property that the gauge group for open strings terminating on a stack of N identical D-branes is $U(N)$ rather than $SU(N)$ for $N > 2$. [For $N = 2$ the gauge group can be $Sp(1)$ rather than $U(2)$.] In this Letter we exploit both these properties in order to obtain a "new physics" signal at LHC which, if traced to low-scale string theory, could with 100 fb^{-1} of integrated luminosity probe deviations from SM physics at a 5σ significance for M_s as large as 3.3 TeV.

To develop our program in the simplest way, we will work within the construct of a minimal model in which we consider scattering processes which take place on the (color) $U(3)$ stack of D-branes. In the bosonic sector, the open strings terminating on this stack contain, in addition to the $SU(3)$ octet of gluons, an extra $U(1)$ boson (C_μ, in the notation of [4]), most simply the manifestation of a gauged baryon number symmetry. The $U(1)_Y$ boson Y_μ, which gauges the usual electroweak hypercharge symmetry, is a linear combination of C_μ, the $U(1)$ boson B_μ terminating on a separate $U(1)$ brane, and perhaps a third additional $U(1)$ (say W_μ) sharing a $U(2)$ brane which is also a terminus for the $SU(2)_L$ electroweak gauge bosons W^a_μ. Thus, critically for our purposes, the photon A_μ, which is a linear combination of Y_μ and W^3_μ will participate with the gluon octet in (string) tree level scattering processes on the color brane, processes which in the SM occur only at one-loop level. Such a mixing between hypercharge and baryon number is a generic property of D-brane quivers; see, e.g., Refs. [4–6].

The process we consider (at the parton level) is $gg \to g\gamma$, where g is an $SU(3)$ gluon and γ is the photon. As explicitly calculated below, this will occur at string disk (tree) level, and will be manifest at LHC as a non-SM contribution to $pp \to \gamma + \text{jet}$. A very important property of string disk amplitudes is that they are completely model independent; thus the results presented below are robust, because they hold for arbitrary compactifications of superstring theory from ten to four dimensions, including those that break supersymmetry. The SM background for this signal originates in the parton tree level processes $gq \to \gamma q$, $g\bar{q} \to \gamma\bar{q}$, and $q\bar{q} \to \gamma g$. Of course, the SM processes will also receive stringy corrections which should be added to the pure bosonic contribution as part of the signal [7–10]. We leave this evaluation to a subsequent publication [11]; thus, the contribution from the bosonic process calculated here is to be regarded as a lower bound to the stringy signal. It should also be stated that, in what follows, we do not include effects of Kaluza-Klein recurrences due to compactification. We assume that all such effects are in the gravitational sector, and hence occur at higher order in string coupling [7].

The most direct way to compute the amplitude for the scattering of four gauge bosons is to consider the case of polarized particles because all nonvanishing contributions can be then generated from a single, maximally helicity violating (MHV), amplitude—the so-called partial MHV amplitude [12]. Assume that two vector bosons, with the momenta k_1 and k_2, in the $U(N)$ gauge group states corresponding to the generators T^{a_1} and T^{a_2} (here in the fundamental representation), carry negative helicities while the other two, with the momenta k_3 and k_4 and gauge group states T^{a_3} and T^{a_4}, respectively, carry positive helicities. Then the partial amplitude for such an MHV configuration is given by [13,14]

$$A(1^-,2^-,3^+,4^+) = 4g^2\mathrm{Tr}(T^{a_1}T^{a_2}T^{a_3}T^{a_4}) \times \frac{\langle 12\rangle^4}{\langle 12\rangle\langle 23\rangle\langle 34\rangle\langle 41\rangle}V(k_1,k_2,k_3,k_4), \tag{1}$$

where g is the $U(N)$ coupling constant, $\langle ij\rangle$ are the standard spinor products written in the notation of Refs. [15,16], and the Veneziano form factor,

$$V(k_1,k_2,k_3,k_4) = V(s,t,u) = \frac{\Gamma(1-s)\Gamma(1-u)}{\Gamma(1+t)}, \tag{2}$$

is the function of Mandelstam variables, here normalized in the string units

$$s = \frac{2k_1k_2}{M_s^2}, \qquad t = \frac{2k_1k_3}{M_s^2}, \qquad u = \frac{2k_1k_4}{M_s^2}: \qquad s+t+u=0. \tag{3}$$

(For simplicity we drop carets for the parton subprocess.) Its low-energy expansion reads

$$V(s,t,u) \approx 1 - \frac{\pi^2}{6}su - \zeta(3)stu + \ldots. \tag{4}$$

We are interested in the amplitude involving three $SU(N)$ gluons g_1, g_2, g_3, and one $U(1)$ gauge boson γ_4 associated to the same $U(N)$ quiver:

$$T^{a_1} = T^a, \qquad T^{a_2} = T^b, \qquad T^{a_3} = T^c, \qquad T^{a_4} = QI, \tag{5}$$

where I is the $N\times N$ identity matrix and Q is the $U(1)$ charge of the fundamental representation. The $U(N)$ generators are normalized according to

$$\mathrm{Tr}\,(T^aT^b) = \tfrac{1}{2}\delta^{ab}. \tag{6}$$

Then the color factor

$$\mathrm{Tr}\,(T^{a_1}T^{a_2}T^{a_3}T^{a_4}) = Q(d^{abc} + \tfrac{i}{4}f^{abc}), \tag{7}$$

where the totally symmetric symbol d^{abc} is the symmetrized trace while f^{abc} is the totally antisymmetric structure constant.

The full MHV amplitude can be obtained [13,14] by summing the partial amplitudes (1) with the indices permuted in the following way:

$$\mathcal{M}(g_1^-,g_2^-,g_3^+,\gamma_4^+) = 4g^2\langle 12\rangle^4\sum_\sigma \frac{\mathrm{Tr}(T^{a_{1\sigma}}T^{a_{2\sigma}}T^{a_{3\sigma}}T^{a_4})V(k_{1_\sigma},k_{2_\sigma},k_{3_\sigma},k_4)}{\langle 1_\sigma 2_\sigma\rangle\langle 2_\sigma 3_\sigma\rangle\langle 3_\sigma 4\rangle\langle 41_\sigma\rangle}, \tag{8}$$

where the sum runs over all 6 permutations σ of $\{1,2,3\}$ and $i_\sigma \equiv \sigma(i)$. As a result, the antisymmetric part of the color factor (7) cancels and one obtains

$$\mathcal{M}(g_1^-,g_2^-,g_3^+,\gamma_4^+) = 8Qd^{abc}g^2\langle 12\rangle^4\left(\frac{\mu(s,t,u)}{\langle 12\rangle\langle 23\rangle\langle 34\rangle\langle 41\rangle} + \frac{\mu(s,u,t)}{\langle 12\rangle\langle 24\rangle\langle 13\rangle\langle 34\rangle}\right), \tag{9}$$

where

$$\mu(s,t,u) = \Gamma(1-u)\left(\frac{\Gamma(1-s)}{\Gamma(1+t)} - \frac{\Gamma(1-t)}{\Gamma(1+s)}\right). \tag{10}$$

All nonvanishing amplitudes can be obtained in a similar way. In particular,

$$\mathcal{M}(g_1^-,g_2^+,g_3^-,\gamma_4^+) = 8Qd^{abc}g^2\langle 13\rangle^4\left(\frac{\mu(t,s,u)}{\langle 13\rangle\langle 24\rangle\langle 14\rangle\langle 23\rangle} + \frac{\mu(t,u,s)}{\langle 13\rangle\langle 24\rangle\langle 12\rangle\langle 34\rangle}\right), \tag{11}$$

and the remaining ones can be obtained either by appropriate permutations or by complex conjugation.

In order to obtain the cross section for the (unpolarized) partonic subprocess $gg\to g\gamma$, we take the squared moduli of individual amplitudes, sum over final polarizations and colors, and average over initial polarizations and colors. As an example, the modulus square of the amplitude (8) is

$$|\mathcal{M}(g_1^-,g_2^-,g_3^+,\gamma_4^+)|^2 = 64Q^2d^{abc}d^{abc}g^4 \times \left|\frac{s\mu(s,t,u)}{u} + \frac{s\mu(s,u,t)}{t}\right|^2. \tag{12}$$

Taking into account all $4(N^2-1)^2$ possible initial polarization or color configurations and the formula [17]

$$\sum_{a,b,c} d^{abc}d^{abc} = \frac{(N^2-1)(N^2-4)}{16N}, \tag{13}$$

we obtain the average squared amplitude

$$|\mathcal{M}(gg\to g\gamma)|^2 = g^4Q^2C(N)\left\{\left|\frac{s\mu(s,t,u)}{u} + \frac{s\mu(s,u,t)}{t}\right|^2 + (s\leftrightarrow t) + (s\leftrightarrow u)\right\}, \tag{14}$$

where

$$C(N) = \frac{2(N^2-4)}{N(N^2-1)}. \tag{15}$$

The two most interesting energy regimes of $gg \to g\gamma$ scattering are far below the string mass scale M_s and near the threshold for the production of massive string excitations. At low energies, Eq. (14) becomes

$$|\mathcal{M}(gg \to g\gamma)|^2 \approx g^4 Q^2 C(N) \frac{\pi^4}{4}(s^4 + t^4 + u^4) \quad (s, t, u \ll 1). \tag{16}$$

The absence of massless poles, at $s = 0$ etc., translated into the terms of effective field theory, confirms that there are no exchanges of massless particles contributing to this process. On the other hand, near the string threshold $s \approx M_s^2$ (where we now restore the string scale)

$$|\mathcal{M}(gg \to g\gamma)|^2 \approx 4g^4 Q^2 C(N) \frac{M_s^8 + t^4 + u^4}{M_s^4[(s - M_s^2)^2 + (\Gamma M_s)^2]} \quad (s \approx M_s^2), \tag{17}$$

with the singularity (smeared with a width Γ) reflecting the presence of a massive string mode propagating in the s channel. In what follows we will take $N = 3$, set g equal to the QCD coupling constant ($g^2/4\pi = 0.1$), and $\Gamma \simeq (g^2/16\pi)(2j + 1)^{-1} M_s$, with $j = 2$ [18]. Before proceeding with numerical calculation, we need to make precise the value of Q. If we were considering the process $gg \to Cg$, where C is the $U(1)$ gauge field tied to the $U(3)$ brane, then $Q = \sqrt{1/6}$ due to the normalization condition (6). However, for $gg \to \gamma g$ there are two additional projections: from C_μ to the hypercharge boson Y_μ, giving a mixing factor κ; and from Y_μ onto a photon, providing an additional factor $\cos\theta_W$ ($\theta_W =$ Weinberg angle). The C-Y mixing coefficient is model dependent: in the minimal model [4] it is quite small, around $\kappa \simeq 0.12$ for couplings evaluated at the Z mass, which is modestly enhanced to $\kappa \simeq 0.14$ as a result of renormalization group running of the couplings up to 2.5 TeV. It should be noted that in models [5,6] possessing an additional $U(1)$ which partners $SU(2)_L$ on a $U(2)$ brane, the various assignment of the charges can result in values of κ which can differ considerably from 0.12. In what follows, we take as a fiducial value $\kappa^2 = 0.02$. Thus, if (17) is to describe $gg \to \gamma g$,

$$Q^2 = \tfrac{1}{6}\kappa^2 \cos^2\theta_W \simeq 2.55 \times 10^{-3} \quad (\kappa^2/0.02). \tag{18}$$

In order to assess the possibility of discovery of new physics above background at LHC, we adopt the kind of signal introduced in [19] to study detection of TeV-scale black holes at the LHC, namely, a high-$k_\perp$ isolated γ or Z. Thus, armed with parton distribution functions (CTEQ6D) [20] we have calculated integrated cross sections $\sigma(pp \to \gamma + \text{jet})|_{k_\perp(\gamma) > k_{\perp,\min}}$ for both the background QCD processes and for $gg \to \gamma g$, for an array of values for the string scale M_s. Our results are shown in Fig. 1. As can be seen in the left panel, the background is significantly reduced for large $k_{\perp,\min}$. At very large values of $k_{\perp,\min}$, however, event rates become problematic. In the right panel we show the cross section and number of events (before cuts) in a 100 fb^{-1} run at LHC for both SM processes (dashed line) and for the string amplitude (solid line), for $k_{\perp,\min} = 300$ GeV, as a function of the string scale M_s.

Our significant results are encapsuled in Fig. 2, where we show the signal-to-noise ratio (signal/$\sqrt{\text{SM background}}$) as a function of M_s for an integrated luminosity of 100 fb^{-1}. The solid line indicates an optimistic case with $\kappa^2 = 0.02$, and 100% detector efficiency with no additional cuts beyond $k_\perp(\gamma) > 300$ GeV. This allows 5σ discovery for M_s as large as 3.5 TeV. The dashed ($\kappa^2 = 0.01$) and dot-dashed ($\kappa^2 = 0.02$) lines indicate more conservative scenarios in which considerations of detector efficiency and γ isolation cuts reduce the total number of events by a factor of 2. In this case, for $\kappa^2 = 0.02$, discovery is now possible for M_s as large as 3.3 TeV. Even in the pessimistic case, for $\kappa^2 = 0.01$ and 50% detector efficiency, a string scale as large as 3.1 TeV can be discovered. The dotted curve allows an illustrative view of the LHC reach in a hypothetical non-

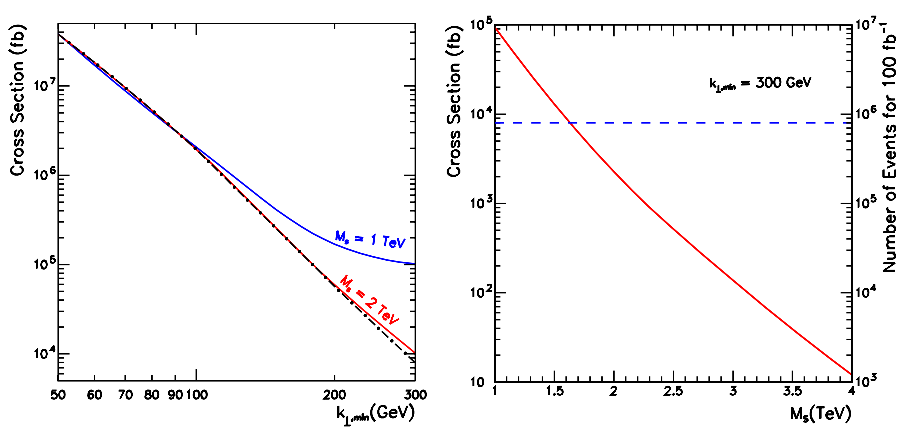

FIG. 1 (color online). In the left panel we display the behavior of the QCD cross section (dot-dashed line) and string + QCD cross section (solid line) for $pp \to \gamma + \text{jet}$ for two values of the string scale. In the right panel we show the cross section and the number of events for fixed $k_{\perp\min} = 300$ GeV and varying string scale. The horizontal dashed line represents the SM background.

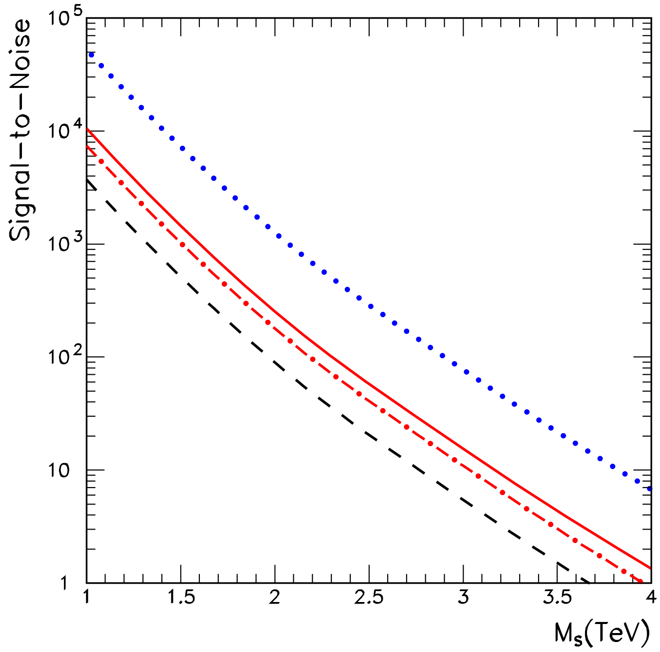

FIG. 2 (color online). Signal-to-noise ratio for an integrated luminosity of 100 fb^{-1}. The solid line indicates the optimistic case with 100% detector efficiency and $\kappa^2 = 0.02$. The dashed ($\kappa^2 = 0.01$) and dot-dashed ($\kappa^2 = 0.02$) lines indicate more conservative scenarios in which considerations of detector efficiency and selection cuts reduce the total number of events by a factor of 2. The dotted line corresponds to a very optimistic nonminimal case with $\kappa^2 = 0.1$ and 100% detector efficiency.

minimal (and optimistic) scenario in which the baryon $U(1)$ (C_μ) has a sizeable hypercharge component.

In summary, we have shown that cross section measurements of the process $pp \to \text{high} - k_\perp \gamma + \text{jet}$ at the LHC will attain 5σ discovery reach on low-scale string models for M_{string} as large as 3.3 TeV, even with detector efficiency of 50% [21]. In order to minimize misidentification with a high-$k_\perp \pi^0$, isolation cuts must be imposed on the photon, and to trigger on the desired channel, the hadronic jet must be identified [22]. We will leave the exact nature of these cuts for the experimental groups.

In closing, we would like to note that the results presented here are conservative, in the sense that we have not included in the signal the stringy contributions to the SM processes. These will be somewhat more model dependent since they require details of the fermion quiver assignments, but we expect that these contributions can potentially double the signal, significantly increasing the reach of LHC for low-scale string discovery. In addition, a similar treatment of $pp \to Z + \text{jet}$, $Z \to \ell^+ \ell^-$ could provide a potentially cleaner signal. The stringy calculation to include transverse Z's will be presented in a future work.

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