Alternative Signature of TeV Strings: Reduction in QCD Jet Production

K. Oda – Theory Group, KEK
N. Okada—University of Maryland

Citation of published version:

Oda, K., Okada, N. (2002): Alternative Signature of TeV Strings: Reduction in QCD Jet Production. *Physical Review D,* 66(9). DOI: https://doi.org/10.1103/PhysRevD.66.095005

Alternative signature of TeV strings: Reduction in QCD jet production

Kin-ya Oda*
Theory Group, KEK, Tsukuba, Ibaraki 305-0801, Japan

Nobuchika Okada†
Department of Physics, University of Maryland, College Park, Maryland 20742

(Received 8 May 2002; published 11 November 2002)

In string theory, it is well known that any hard scattering amplitude inevitably suffers exponential suppression. We demonstrate that, if the string scale is $M_s < 2$ TeV, this intrinsically stringy behavior leads to a dramatic reduction in the QCD jet production rate with very high transverse momenta $p_T \geqslant 2$ TeV at CERN LHC. This suppression is sufficient to be observed in the first year of low-luminosity running. Our prediction is based on the universal behavior of string theory, and therefore is qualitatively model independent. This signature is alternative and complementary to conventional ones such as Regge resonance (or string ball or black hole) production.

DOI: 10.1103/PhysRevD.66.095005 PACS number(s): 11.25.Mj, 11.10.Kk, 13.85.Ni

I. INTRODUCTION

The discovery of D-branes in string theory [1] has greatly enhanced the phenomenological investigation of the interesting possibility that we are living in a four-dimensional subspace ("3-brane") of a larger D-dimensional world ("bulk"), which reveals itself through the gravitational interaction above the TeV scale [2].[1] The hierarchy between the electroweak scale and the four-dimensional Planck scale may be solved by arranging the volume of the extra $D-4$ dimensions so as to lower the fundamental gravitational scale of the D-dimensional theory down to $\mathcal{O}(1\ \text{TeV})$ [2]. This idea can be realized as a class of models in string theory (see Ref. [4] and citations therein), which possess the following characteristics in general. The fundamental string scale M_s is around TeV, where M_s is defined as $M_s = \alpha'^{-1/2}$ with α' being the string tension. The graviton corresponds to the massless mode of the closed string in the bulk, while the standard model fields are realized as the massless modes of open strings whose end points are on the D3-brane. (In general, the "extra" $D-4$ dimensions are compactified with a "large" radius which is $\sim$ fm–mm depending on D, and the remaining $10-D$ dimensions are compactified with string length.) In this paper we present a model-independent prediction derived from the universal nature of TeV scale string theory.

When the string scale is around TeV, we may expect signals at the CERN Large Hadron Collider (LHC), whose center-of-mass energy is designed to be $\sqrt{s} = 14$ TeV. The most direct signature of the TeV strings would be the production of excited string states observed as a number of Regge resonances [5–7]. We note, however, that there exist other new physics such as technicolor theories [8] or preon models [9] which also provide similar resonance events such as "technihadrons" or excited states of the known fermions. It is not easy, in a hadron collider, to exclude other possible particle theories and confirm the TeV scale string theory only from the observation of the resonance states.

In addition, there are somewhat perplexing consequences of semiclassical TeV scale gravity: a black hole might be formed and hide all the interactions for the process at $\sqrt{\hat{s}} \geqslant M_D$ [10,11], where $\sqrt{\hat{s}}$ and M_D are the center-of-mass energy of the scattering partons and the D-dimensional Planck scale, respectively.[2] It has been suggested that the black hole production cross section is simply given by the geometrical cross section $\sigma = \pi R_S^2$, where R_S is the Schwarzschild radius of the black hole whose mass is equal to $\hat{s}$. This claim is based on the assumption that the classical hoop conjecture is valid for the quantum process of the partons and thus the black hole formation occurs whenever the impact parameter is smaller than R_S. However, this proposition based on semiclassical quantum gravity is still being debated [12–14].

On the other hand, it has been claimed that black hole formation occurs in TeV scale string theory [15] due to the correspondence principle for black holes and strings [16,17]. The picture is as follows. As one raises the energy of the string scattering, highly excited string states are produced and tangled to form "string balls." If one raises the energy further, the amplitude of the string ball production would be smoothly connected with that of the semiclassical black hole production at the "correspondent point." Here, we note that while the correspondence principle is safe for the static situation (entropy), there is no evidence of it for the dynamical process (S-matrix) involving non-Bogomol'nyi-Prasad-Sommerfield (BPS) black holes.[3] The problem here is that the energy range which can be explored at LHC is not high enough [11] so that the semiclassical treatment of the produced black hole is less reliable. Therefore, even if the black hole is really produced, a deeper understanding of the correspondence principle is necessary to analyze the physical process in detail.

*Email address: kin-ya.oda@kek.jp
†Email address: okadan@physics.umd.edu

[1]The idea of the extra dimensions above the TeV scale was proposed earlier [3].

[2]The D-dimensional Planck scale is given by $M_D \sim g_s^{-2/(D-2)} M_s$, where g_s is the string coupling constant.

[3]We thank M. Natsuume for bringing this point to our attention.

In order to complement the conventional signature of resonance mode (or string ball or black hole) production, which involves the complicated arguments above, it is important to find an alternative. In this paper we employ the hard-scattering behavior of the string amplitude. It is well known that any string amplitudes are exponentially suppressed in the hard scattering limit [18,19]. In order to demonstrate this stringy suppression, we examine the QCD jet production rate with very high transverse momentum. We show that, depending on the string scale, the rate is dramatically reduced from the standard model prediction. This result is the reflection of the intrinsically stringy structure and is distinct from any other new (particle) physics.

II. HARD SCATTERING LIMIT OF STRING AMPLITUDES

Gross and Mende have shown that any closed string amplitudes exhibit a universal behavior in the hard scattering limit, i.e., $s\to\infty$ with fixed t/s [18]. Since these amplitudes are dominated by the tachyonic form factors, this behavior is independent of the type of string theory (bosonic string, superstrings, heterotic strings, etc.), the perturbative string vacuum, and external states of the scattering [19]. This argument has also been extended to open strings [20]. The G-loop amplitude was found to be exponentially suppressed[4] as

$$\mathcal{A}_G(s,t)\sim e^{-\alpha' s f(\theta)/(G+1)}, \tag{1}$$

where the positive function $f(\theta)$ is given in terms of $\lambda=-t/s=\sin^2(\theta/2)$ by $f(\theta)=-\lambda\log\lambda-(1-\lambda)\log(1-\lambda)$. This amplitude reproduces the hard scattering limits of both the Veneziano (open string) and Virasoro-Shapiro (closed string) amplitudes at the tree level, $G=0$.

Although the perturbation series of Eq. (1) is badly divergent, it is still possible to evaluate the limit of the amplitude with its leading terms resummed to all orders by utilizing Borel transform techniques [22]. The resultant suppression is found to be milder, but is still significant:

$$\mathcal{A}_{\mathrm{MO}}(s,t)\sim e^{-c\sqrt{\alpha' s f(\theta)}}, \tag{2}$$

where c is a factor having little angle dependence and is in principle calculable.[5] For simplicity, we set $c=1$ in the following analysis.

In the toy model of the TeV scale string theory, it has been shown at the tree level [7] that every standard model amplitude (corresponding to the open string scattering) is multiplied by a *common* stringy form factor

$$\mathcal{S}(s,t)=\frac{\Gamma(1-\alpha' s)\Gamma(1-\alpha' t)}{\Gamma(1-\alpha' s-\alpha' t)}, \tag{3}$$

which is essentially the Veneziano amplitude. A similar result with different factor $\mathcal{F}(s,t,u,m^2)$ is found for the emission of the graviton (or the production of Kaluza-Klein excitations) which corresponds to the tree level closed string emission [5,7]. Here, m is the mass of the Kaluza-Klein mode which is negligible when $D>6$ [4]. It can be confirmed that both the hard scattering limits of $\mathcal{S}(s,t)$ and $\mathcal{F}(s,t,u,0)$ reproduce the universal behavior of Eq. (1) with $G=0$ [7]. Therefore, it is reasonable to expect that the hard scattering amplitude (of the standard model particles) in TeV scale string theory is described by the standard model amplitude multiplied by the resummed stringy factor $\mathcal{A}(s,t)=e^{-\sqrt{\alpha' s f(\theta)}}$ such that

$$\frac{d\sigma}{dt}=\left.\frac{d\sigma}{dt}\right|_{\mathrm{SM}}|\mathcal{A}(s,t)|^2. \tag{4}$$

This is the basic formula in our analysis.[6]

In what follows, we give several comments to clarify our arguments. Firstly, if the known particles have further substructures as in preon models, the hard scattering amplitude will, in general, become soft at high energies. An interesting feature of string theory is that the stringy structure makes the hard scattering process much softer than any local field theory can do. The history of hadronic interactions tells us that this point played a crucial role in excluding the dual resonance (string) model, while giving support for the QCD parton model (see, e.g. Ref. [23]). Thus our result is distinct from any other new particle physics.

Secondly, we note that a string amplitude such as $\mathcal{S}(s,t)$ is, in general, a rapidly varying function with many zeros and poles. The poles correspond to the existence of the resonance states. Although the factor $e^{-\alpha' s f(\theta)/(G+1)}$ in the hard scattering limit is evaluated off the poles, we can regard it as the amplitude suitably averaged over all zeros and poles (see, e.g.Ref. [23]). Indeed, we can find the same suppression factor even on the pole, if we appropriately take the width into account and remove the singularity.

Finally, there is another limit of the scattering amplitude in the string theory called the Regge limit, i.e., $s\to\infty$ with t fixed [24]. In this limit, the amplitude behaves as $\sim(\alpha' s)^{\alpha' t+\alpha(0)}$, where a constant $\alpha(0)$ is the intercept.[7] In the argument of Ref. [15], the production cross section of the string balls is a monotonically increasing function of s, and is

[4]When initial momenta are transverse to the D-brane, the hard scattering behavior could be modified to the power-law suppression [21]. In our case, initial memonta are parallel to the D-brane ("our" 3-brane), and we may expect the same exponentially suppressed amplitude.

[5]For example, we obtain $c\simeq 1.2$ for the tachyon scattering in the bosonic closed string theory when g_s is set equal to the QCD coupling constant.

[6]We can explicitly show that $\mathcal{S}(s,t)\simeq\sqrt{2\pi\alpha' s}\tan(\theta/2)e^{-\alpha' s f(\theta)}$ utilizing Stirling's formula. The prefactor $\sqrt{2\pi\alpha' s}\tan(\theta/2)$ leads to the subleading contributions in Eq. (2) after the Borel resummation. Even if we naively multiply the stringy factor in Eq. (4) by this prefactor, our final prediction (namely, the deficit in the QCD-jet production rate) is modified only by a factor $<\mathcal{O}(10)$, and hence our conclusion is not changed.

[7]The intercept is positive for the realistic theory having massless modes in the mass spectrum. In particular $\alpha(0)=0$ for $\mathcal{S}(s,t)$.

connected to that of the black hole at the correspondent point. This seems to suggest that, in the string picture, black hole production corresponds to the contributions from $-\alpha(0)/\alpha' \lesssim t \lesssim 0$ in the Regge region since the fixed angle (hard scattering) amplitude is exponentially suppressed.[8]

In the black hole picture, all the substructures smaller than the Schwarzschild radius would be hidden and most particles would be emitted by evaporation through Hawking radiation. Consequently the rate of the emission of the hard quanta higher than the Hawking temperature is exponentially suppressed due to the Planck distribution of Hawking radiation such that $\sim e^{-E/T}$. From this point of view it is interesting that, in the string picture, the perturbative behavior $\sim e^{-E^2}$ in Eq. (1) is modified to $\sim e^{-E}$ as in Eq. (2) when summed up to all orders.

While all our calculations in this paper are performed in the string picture, the exponential suppression of the hard scatterings might be considered as a consequence of black hole formation, if the correspondence principle is valid also for the *S*-matrix and black holes are really formed as is claimed in Ref. [15].

III. HIGH p_T JET PRODUCTION RATE

In order to demonstrate the stringy effect in the hard scattering limit, we study the QCD jet production rate with very high transverse momentum. At LHC, high p_T events can be triggered using reasonable $p_T^{\min} = \mathcal{O}(100\text{ GeV})$ [26] which is far below the scale $\mathcal{O}(1\text{ TeV})$ considered here. We take $p_T^{\min} = 2$ TeV, for which trigger efficiency would be practically 100%. The jet resolution is expected to be of order 10 GeV, hence negligible for our purposes. The jet production cross section (which is dominated by QCD scatterings among quarks and gluons) for a given interval of the transverse momentum $[p_T, p_T + \Delta p_T]$ is evaluated by

$$\Delta\sigma_{[p_T, p_T+\Delta p_T]}(s) = \int_0^1 dx_1 \int_0^1 dx_2 \int d\hat{t} \times \sum_{ijkl} \frac{1}{1+\delta_{kl}} f_i(x_1, Q) f_j(x_2, Q) \times \left. \frac{d\hat{\sigma}_{ij\to kl}}{d\hat{t}} \right|_{\mathrm{SM}} |\mathcal{A}(\hat{s}, \hat{t})|^2, \tag{5}$$

where $\hat{s}$, $\hat{t}$, and $\hat{u}$ are the Mandelstam variables of the parton-parton scattering obeying $\hat{s} = x_1 x_2 s$ and $\hat{s} + \hat{t} + \hat{u} = 0$; the region of integration for $\hat{t}$ is determined by the conditions $p_T^2 < \hat{t}\hat{u}/\hat{s} < (p_T + \Delta p_T)^2$ and $\hat{s} + \hat{t} > 0$ for each set of x_1 and x_2; f_i are parton distribution functions, i,j,k,l are summed over gluon and quark flavors, and Q is the typical energy of the parton scattering defined as $Q^2 = 2\hat{s}\hat{t}\hat{u}/(\hat{s}^2 + \hat{t}^2 + \hat{u}^2)$. $d\hat{\sigma}_{ij\to kl}/d\hat{t}|_{\mathrm{SM}}$ are the QCD parton-level cross sections of the standard model (summarized, e.g., in Ref. [27]). In our analysis, we employ CTEQ5M1 [28] for the parton distribution functions.

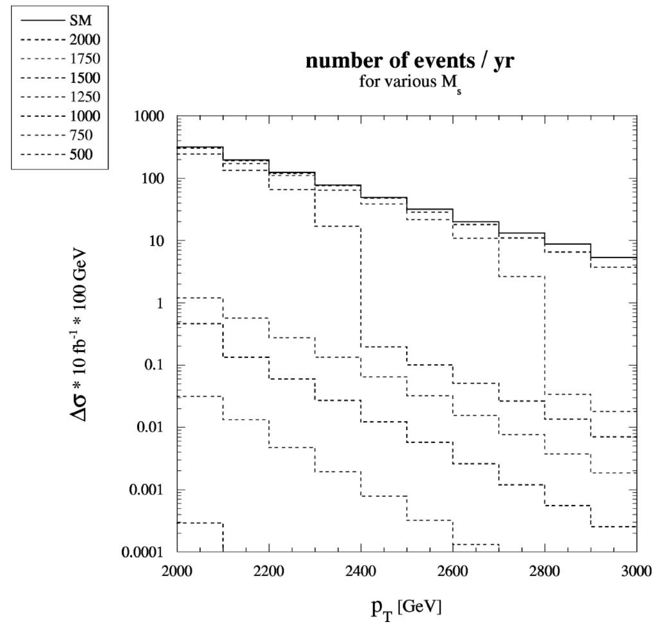

FIG. 1. Number of jet production events per year in the 100 GeV bin, expected to be observed in the LHC with the first low-luminosity running (10 fb^{-1}/yr). The solid line is the standard model prediction and the lines below are the corresponding results in TeV scale string theory with the string scale varied from $M_s = 2$ TeV to 0.5 TeV.

The result is shown in Fig. 1. Here, the center-of-mass energy $\sqrt{s}$ is taken at $\sqrt{s} = 14$ TeV as is designed for LHC. We have plotted the number of events per year for the QCD jet production expected for the low-luminosity running (10 fb^{-1}/yr) within 100 GeV bin from $p_T = 2$ TeV to 3 TeV. The solid line denotes the standard model prediction and the lines below correspond to the predictions of TeV scale string theory with the string scale from $M_s = 2$ TeV to 0.5 TeV.

In this analysis we have applied the stringy factor $e^{-\sqrt{\alpha'\hat{s}}f(\theta)}$ only when $\hat{s} > 10M_s^2$ [i.e.we have conservatively applied $\mathcal{A}(\hat{s}, \hat{t}) = 1$ when $\hat{s} \leq 10M_s^2$].[9] The result barely depends on the value of this cutoff; the condition $\hat{s} > 10M_s^2$ is already satisfied in most regions of the integration that we perform at very high transverse momenta.[10]

[8] This correspondence in the Regge region was mentioned in the earlier work [25].

[9] It is estimated that the hard scattering limit is applicable at $N \gtrsim (g_s^2/4\pi)^{-1}$, where $N \sim \alpha'\hat{s}$ [29]. Note that we are treating QCD scatterings with the coupling constant $g \sim 1$ and also that the string coupling constant is expected to be of $\mathcal{O}(1)$ in the original consideration of the TeV scale string theory [4].

[10] More explicitly, this condition is satisfied in all the region of 2 TeV $\leq p_T \leq$ 3 TeV for $M_s < 1.5$ TeV. In Fig. 1, we can observe the influence of this cutoff at $p_T < 2.4$ TeV, 2.8 TeV and 2 TeV $\leq p_T \leq 3$ TeV for $M_s = 1.5$ TeV, 1.75 TeV and 2 TeV, respectively.

IV. SUMMARY

We have demonstrated that the exponential suppression of the hard scattering amplitude in TeV scale string theory leads to a dramatic reduction in the production rate of the QCD jets with very high transverse momentum $p_T \geqslant 2$ TeV at LHC when the string scale is at $M_s < 2$ TeV. This deficit is sufficient to be observed in the first year low-luminosity running. This signature is complementary to the conventional signature of Regge resonance (or string ball or black hole) production. We note that although we have presented the calculation only with the leading QCD scatterings among quarks and gluons, *any* other standard model cross sections will suffer the same suppression due to the universal nature of the high energy behavior of the string theory. It is worth investigating various hard scattering processes in the TeV scale string theory.

Note added

In our calculation, we have considered the string loop corrections. In the field theory, the infrared divergences from virtual soft gluon loops are canceled by the real soft gluon emissions from external lines and are absent when we treat a properly constructed jet cross section. This cancellation is not altered in the string theory. This can be explained in two ways. Firstly, string amplitude is calculated in such a way that external lines are inserted as vertex operators, which is already mode expanded and the ‘‘field-theoretical limit’’ is taken (by utilizing the modular invariance), and therefore treatment of this type of infrared divergences (i.e. the argument of attaching virtual and real soft gluon lines onto external lines) is not modified from the field-theoretical treatment. Secondly, the fact that the sum of virtual and real soft gluon contributions may be calculated by a properly constructed jet cross section follows from the unitarity of the theory. We note that the theory remains unitary when we go beyond the low energy effective theory to string theory.

The precise value of the infrared cutoff is not important to our conclusion because we are estimating a kind of inclusive cross section of the high p_T jet event with any number of additional jets. Though no one has ever succeeded in explicitly showing that the $2 \to N$ process is always suppressed in string theory when $2 \to 2$ process is exponentially suppressed, this fact is widely believed and proving this is beyond the scope of the current paper. In this paper we simply assume that this is the case and present that dominant $2 \to 2$ contributions are suppressed.

ACKNOWLEDGMENTS

We are grateful to N. Ishibashi, K. Odagiri, J. Hisano, and Y. Kitazawa for useful comments, and A. Akeroyd and K. Odagiri for reading the manuscript. K. O. would like to thank I. Kishimoto for helpful conversations in the early stages of this work and especially T. Suyama and M. Natsuume for extensive discussions.

[1] J. Polchinski, Phys. Rev. Lett. **75**, 4724 (1995).
[2] N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. B **429**, 263 (1998).
[3] I. Antoniadis, Phys. Lett. B **246**, 377 (1990).
[4] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. B **436**, 257 (1998).
[5] E. Dudas and J. Mourad, Nucl. Phys. **B575**, 3 (2000).
[6] E. Accomando, I. Antoniadis, and K. Benakli, Nucl. Phys. **B579**, 3 (2000).
[7] S. Cullen, M. Perelstein, and M.E. Peskin, Phys. Rev. D **62**, 055012 (2000).
[8] E. Farhi and L. Susskind, Phys. Rep. **74**, 277 (1981).
[9] R.E. Marshak, Phys. Rep. **137**, 73 (1986).
[10] S.B. Giddings and S. Thomas, Phys. Rev. D **65**, 056010 (2002).
[11] S. Dimopoulos and G. Landsberg, Phys. Rev. Lett. **87**, 161602 (2001).
[12] M.B. Voloshin, Phys. Lett. B **518**, 137 (2001).
[13] S.B. Giddings, hep-ph/0110127.
[14] M.B. Voloshin, Phys. Lett. B **524**, 376 (2002).
[15] S. Dimopoulos and R. Emparan, Phys. Lett. B **526**, 393 (2002).
[16] G.T. Horowitz and J. Polchinski, Phys. Rev. D **55**, 6189 (1997).
[17] G.T. Horowitz and J. Polchinski, Phys. Rev. D **57**, 2557 (1998).
[18] D.J. Gross and P.F. Mende, Phys. Lett. B **197**, 129 (1987).
[19] D.J. Gross and P.F. Mende, Nucl. Phys. **B303**, 407 (1988).
[20] D.J. Gross and J.L. Manes, Nucl. Phys. **B326**, 73 (1989).
[21] J.L.F. Barbon, Phys. Lett. B **382**, 60 (1996).
[22] P.F. Mende and H. Ooguri, Nucl. Phys. **B339**, 641 (1990).
[23] M. B. Green, J. H. Schwarz, and E. Witten, *Superstring Theory. Vol. 1: Introduction*, Cambridge Monographs On Mathematical Physics (Cambridge University Press, Cambridge, U.K, 1987).
[24] D. Amati, M. Ciafaloni, and G. Veneziano, Phys. Lett. B **197**, 81 (1987).
[25] D. Amati, M. Ciafaloni, and G. Veneziano, Phys. Lett. B **216**, 41 (1989).
[26] CMS Technical Proposal, CERN/LHCC/94–38 (1994); ATLAS Technical Proposal, CERN/LHCC/94–43 (1994); ATLAS Technical Design Report, CERN/LHCC/99–15 (1999); http://atlasinfo.cern.ch/Atlas/.
[27] R.K. Ellis, W.J. Stirling, and B.R. Webber, *Cambridge Monographs on Particle Physics, Nuclear Physics, and Cosmology: 8* (Cambridge University Press, Cambridge, U.K., 1996).
[28] CTEQ Collaboration, H.L. Lai *et al.*, Eur. Phys. J. C **12**, 375 (2000).
[29] F. Cornet, J.I. Illana, and M. Masip, Phys. Rev. Lett. **86**, 4235 (2001).