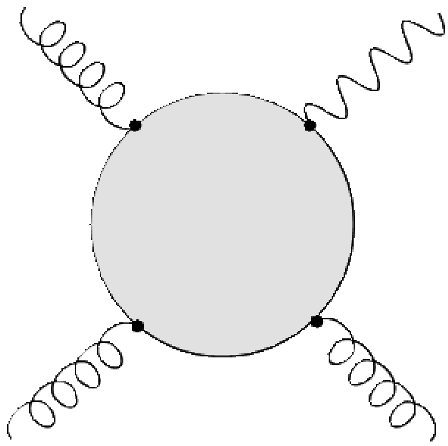
String Phenomenology at the LHC: $gg \to \gamma g$ as a signal for low mass strings



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Abstract

This is essentially a review paper on some of the potential string-induced phenomenology that might be observed at the LHC. I briefly recap the motivations behind low mass string theories and then discuss a number of potential experimental consequences, finally focusing on a particularly promising signal for low mass strings : The $gg \to \gamma g$ process which only occurs at loop level in the SM, but which is present at the string disk level (i.e. at tree level) in low mass string theories. Having worked out the relevant scattering amplitudes I estimate the discovery reach for this type of signal at the LHC.

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1 Introduction : TeV scale strings as a solution to the hierarchy problem

The question posed by the problem of mass hierarchy in fundamental particle physics is essentially why there is such an enormous difference between the scales associated with gravitation - the Planck mass $M_P \sim 10^{19} GeV$ - and with electroweak and strong nuclear forces - the electroweak scale $m_{EW} \sim 10^2 GeV$ and the strong scale $\Lambda_{QCD} \sim 100 MeV$. However, in a quantum theory of particles along the lines of the standard model, radiative corrections to the masses of all particles will be generated by the vacuum expectation value of the Higgs. These corrections are proportional to the ultraviolet cutoff of the theory in question, which is given by M_P in the presence of gravity, and as such one should expect particle masses to lie in the region of the Planck scale. To prevent this from happening and restore the actually observed particle masses a severe fine-tuning in the fundamental parameters of the theory is required. Even the smallest divergences (the fine-tuning would have to hold to more than 30 places) would result in drastically different masses [1].

A number of different theories have been proposed to give a more natural explanation to the observed mass hierarchy than coincidental fine-tuning. However, before having a closer look at them, it is worth pointing out that, from a methodological point of view, most, if not all of them share the common strategy of "installing" a mechanism for cutting off the divergences in the TeV region. As such they proceed in analogy to the electroweak case in the standard model, where the radiative stability of the weak scale is ensured by taking m_{EW} as the ultraviolet cutoff of the theory. In other words, they effectively introduce a new scale at TeV energies thus stabilizing particle masses at their experimental values - as opposed to attempting to find a mechanism directly providing the apparent fine-tuning.

In low-energy supersymmetry scenarios with superparticle masses in the TeV region the quadratic divergences in the Higgs-induced mass corrections are cut off by the supersymmetry breaking mass splittings (in the softly broken case - in the exact supersymmetry limit they cancel). The hierarchy problem is consequently solved by fiat and furthermore such theories come equipped with a number of well-known other benefits such as the potential role of dark matter candidates which the new superparticles can play. However, from a string theoretical point of view, a serious problem remains. When embedding a low energy supersymmetric theory into string theory the Planck scale becomes proportional to the string scale. Unfortunately the embedding also gives rise to a separation of the grand unification and string scales by almost two orders of magnitude. Thus it appears that a new hierarchy problem emerges which will need to be explained by the introduction of yet again new parameters or scales.[1] Whilst this does not take away from the benefits offered by low-energy supersymmetric theories, it raises the question whether we can do without the introduction of these new parameters or scales.

Split supersymmetry is such a candidate. Here scalar masses are heavy, whereas fermions such as gauginos and higgsinos are light. Gauge coupling unification is preserved, yet by introducing an explicit difference between the masses for supersymmetric fermions and scalar particles the cancelation of divergent mass-correcting loop effects ceases to work as before and one still either needs fine-tuning in the low energy regime or another mechanism accounting for the mass hierarchy.

The emergence of new strong dynamics as provided by e.g. technicolor and little Higgs models provides another option. Here a cutoff in the TeV range is introduced by rendering the low(er) energy dynamics invalid at such scales. The generation of quadratic divergences via Higgs-induced effects is avoided, since a different dynamics takes over at scales $\approx$ TeV.

Finally models with TeV scale strings make use of the extra dimensions appearing in string theory in order to lower the string scale M_S so that it will act as a cutoff scale for the divergences. Low energy supersymmetry is not needed in these frameworks. However, the fundamental idea is in fact independent of the nature of the ultimate theory of gravity and hence not irrevocably tied to string theory. Instead it is the existence of extra dimensions that serves as the basis for such arguments. Suppose we have an additional n compactified spatial dimensions of radius R, where we leave the exact topology associated with that compactification (e.g. a torus) unspecified. Hence a 4+n dimensional theory is yielded. Now consider two test masses m_1 and m_2 placed at a distance $r \ll R$ from each other. Then we can neglect the fact that some dimensions are compactified with radius R (whilst others are fully extended) and work out the gravitational potential for 4+n dimensions via Gauss' law to get

$$V(r) \sim \frac{m_1 m_2}{M_{P(4+n)}^{n+2}} \frac{1}{r^{n+1}} \qquad , \qquad r \ll R, \tag{1}$$

where $M_{P(4+n)}$ is the "Planck scale" of the 4+n dimensional theory. Contrast this case with the one where the test masses are placed at a distance $r \gg R$. Here gravitational flux lines cannot penetrate all the way into the extra dimensions and the familiar $\frac{1}{r}$ potential (as in the 4-dimensional case) is yielded

$$V(r) \sim \frac{m_1 m_2}{M_{P(4+n)}^{n+2} R^n} \frac{1}{r} \qquad , \qquad r \gg R, \tag{2}$$

and consequently, and this is the crux of the argument, the effective 4-dimensional Planck scale M_P is

$$M_P^2 \sim M_{P(4+n)}^{n+2} R^n. \tag{3}$$

The Planck scale of the underlying 4+n dimensional theory can therefore, in principle, be chosen to be significantly smaller than the effective 4-dimensional scale at the expense of introducing an additional n dimensions with a "large" radius R. This is exciting news, because it opens up the possibility of solving the hierarchy problem by nullification. If $M_{P(4+n)} \sim O(TeV)$, then the weakness of gravity in 4 dimensions as compared to the other gauge interactions is a consequence of the large compactification radii R compared to the fundamental length scale m_{EW}^{-1}. As such gravitational forces do indeed become comparable to the other gauge forces at the weak scale. Hence there is only one fundamental scale in the theory. If the abolition of M_P as a fundamental scale in nature seems too much of a blasphemy, one should keep in mind that, whilst electroweak interactions have been tested at length scales $\sim m_{EW}^{-1}$, gravitational interactions have only been probed to about $\sim 1cm$ and, from an experimental point of view, the belief that M_P stays a fundamental scale when probing gravity down to scales of about $\sim M_P^{-1} = 10^{-33}cm$ is certainly quite a leap of faith.

Of course there are constraints. The fact alluded to above, that gauge forces within the standard model have been probed to scales of $\sim m_{EW}^{-1}$, also means that SM particles must be localized on a 4-dimensional sub-manifold of the full 4+n dimensional space. Only the graviton could propagate freely into the extra dimensions. Furthermore, as I will show in the following section, there are constraints on the number n of extra dimensions via the size of R.

Up to now the argument has been independent of any specific realization of a theory of quantum gravity. When embedding the above into string theory, which already comes equipped with extra dimensions, the string scale M_S takes over the role of the 4+n dimensional Planck scale. For example, when considering a type I string theory with D-branes and 9 spatial dimensions, one can create a so-called braneworld description of our universe, i.e. one where our universe is localized on a p-dimensional hypersurface (where $p \geqslant 3$). Now, assuming an isotropic transverse space of $n = 9 - p$ compactified dimensions with a common radius R, we get the analogue of (3) [1]

$$M_P^2 = \frac{1}{g^2} M_S^{2+n} R_\perp^n \qquad , \qquad g_S \simeq g^2, \tag{4}$$

where g_S is the string coupling. This approach therefore radically breaks with the traditional line of thought that string theory only becomes relevant at distances comparable to the Planck length $\sim 10^{-33}cm$ and utilizes the extra dimensions intrinsic to string theory to lower M_S into the TeV region.

It is remarkable that all of the above proposals for "solving" the hierarchy problem will be experimentally testable at the LHC and one can therefore be hopeful that there will be additional guidance from experiments on how to best approach the hierarchy problem in the near future. In the remainder of this paper, however, I will focus on theories with TeV scale strings and the possible phenomenology they might induce. By making use of the intrinsic resources of string theory (extra dimensions) and since they do not need the introduction of other additional parameters and scales, they do appear to be a very natural answer to the hierarchy problem from inside string theory - a theory which, after all, has long been viewed as the primary candidate for a unification of Planck scale quantum gravity and (sub-)TeV scale standard model physics.

2 Potential low mass string signals at the LHC

In this section I will try to give a very broad overview of some of the possible string phenomenology at the LHC and other experiments, before focussing on the particular example of the $gg \to \gamma g$ interaction in section 3. So what are the experimental signatures one can expect from TeV scale strings and extra dimensions?

1) Production of Regge excitations. The existence of such excitations is due to having TeV scale strings and does not itself require the existence of extra dimensions. For M_S in the TeV region a tower of infinite string excitations opens up above the string mass threshold, with new particles following the Regge trajectories for vibrating strings [2]

$$j = j_0 + \alpha' M^2, \tag{5}$$

where j is the spin and α' the Regge slope parameter, which is related to the string mass scale by

$$M_S^2 = (\alpha')^{-1}. \tag{6}$$

The string states associated to these excitations will yield new contributions to standard model processes such as the scattering of quarks and gluons. At tree-level the contributions are due to the exchange of massive string excitations which encompass all Regge recurrences. For some processes, e.g. n-gluon scattering in QCD, these contributions are independent of the details of compactification and therefore provide a model-independent signature for TeV scale strings.

2) New exotic particles may appear around M_S. As an example new massive Z' gauge bosons are predicted by many string models. These new particles are connected to additional U(1) gauge symmetries, which are a consequence of the fact that the gauge group for open strings terminating on a stack of N identical D-branes (see section 3.1. for details on why this is relevant) is U(N) and not SU(N), where $N \geqslant 2$.

3) Kaluza-Klein and winding excitations for all standard model particles along the extra dimensions parallel to the p-brane (i.e. the p-dimensional hypersurface our universe is localized on). Their spectrum depends on the topology of the extra-dimensions and hence on the details of compactification for those dimensions. As an example, the masses of these excitations for one extra parallel spatial dimension are [1]

$$M_m^2 = M_0^2 + \frac{m^2}{R_\parallel^2} \qquad , \qquad m = 0, \pm 1, \pm 2, ... \tag{7}$$

where $R_\parallel$ is the compactification radius of the parallel dimension. The virtual exchange of such excitations will then contribute to standard model processes and lead to deviating cross-sections, which will in turn also establish experimental bounds on $R_\parallel$. At high enough energies a direct production of these excitations is possible as well.

4) In typical type I string theories, SM particles are open strings completely confined to a specific D-brane, but gravitons are free to propagate into the extra dimensions. At energies above M_S a large number of events with missing energy carried away by the gravitons will be observable. In hadron colliders, processes with jets + missing energy provide a good signature for this kind of effect. Furthermore, gravitational interactions just prior to collisions will lead to an abrupt decrease in the beam energy.

In other possible realizations of theories with extra-dimensions (see section 1) which do not completely confine standard model particles to our 4-dimensional world, but only localize them up to some energy scale allowing escape of such particles above this scale, even more spectacular "escape" phenomenology could be yielded, with high energy SM particles disappearing and appearing again seemingly out of nowhere. For details, see [3].

5) As equation (1) might already have suggested, radical modifications to gravitational forces on scales below 1 cm are possible. When putting $M_{P(4+n)} \sim m_{EW}$ in (3) and demanding that M_P takes its familiar value we obtain

$$R \sim 10^{\frac{30}{n}-17} cm \times \left(\frac{1TeV}{m_{EW}}\right)^{1+\frac{2}{n}}. \tag{8}$$

Hence, for $n = 1$ we get $R \sim 10^{13} cm$ which would lead to modifications to gravity on solar system scales and which is consequently excluded empirically. For $n = 2$ however one only gets $R \sim 1mm$. Table-top experiments which have already tested Newtonian gravitation at short distances (e.g. involving a torsion pendulum) should be able to test this case soon. As can be seen from (1) and (2) one would expect a change in Newton's law from a $\frac{1}{r}$ to a $\frac{1}{r^{n+2}}$ behavior for $n = 2$ at these scales. For higher values of n the associated compactification radius R becomes yet smaller and modifications to gravitational force-laws at small scales will unlikely be observed in the near future. For other potential new sub-millimeter forces such as new scalar forces related to supersymmetry breaking, see [1].

6) Finally, at energies above the string scale, quantum gravity effects related to string physics such as mini black-hole production are possible. For further details, see e.g. [4].

The particular signature $gg \rightarrow \gamma g$ we will consider in the next section essentially falls into categories 1) and 2) above.

3 The case of $gg \to \gamma g$

3.1 Regge recurrences and an additional U(1) symmetry

In a type I string theory gauge interactions emerge as excitations of open strings with endpoints confined on D-branes. In a stack of N identical D-branes, each D-brane "gives rise" to a U(1) gauge symmetry, so to speak, so that the whole stack generates a U(N) gauge theory with an associated U(N) gauge group. Strings whose endpoints are fixed to the same stack are associated with gauge bosons, whereas strings stretching between two different stacks give rise to chiral matter such as fermions. Here we will consider processes taking place on a U(3) color stack of D-branes. Open strings with endpoints confined to this stack include the SU(3) gluons and an extra U(1) boson C_μ. The familiar electroweak hypercharge boson Y_μ is a linear combination of C_μ and other possible U(1) bosons terminating on different branes. The crucial point here is that the photon A_μ, via its dependence on Y_μ, will contribute to the tree-level gluon scattering on the U(3) color brane [5].

The $gg \to \gamma g$ process includes the SU(3) gluons g and a photon γ and is now possible at tree-level due to the new U(1) symmetry arising in the above construction. Hence this process has no equivalent in the standard model, where this symmetry is missing. Its observation would therefore amount to direct evidence for a low string mass Regge recurrence and will be observable as a contribution to $pp \to \gamma + jet$. Moreover, since string disk amplitudes (which govern processes involving four gluons and those involving two gluons and two quarks) are independent of the details of compactification, $gg \to \gamma g$ provides a completely model-independent signature for low mass strings.

3.2 Amplitudes

In order to calculate the amplitude for $gg \to \gamma g$ we will make use of so-called partial maximally helicity violating (MHV) amplitudes. These take a particularly simple form for the scattering of n gauge bosons and represent the amplitudes for the scattering of particles in specific helicity eigenstates. However, it is possible to "sew" such individual MHV amplitudes together in order to build arbitrarily complex tree diagrams, so that one is able to compute the full scattering amplitude by summing over the appropriate partial amplitudes.

As conjectured in [6] and later on derived in [7] the standard model n-gluon scattering helicity amplitudes $\mathcal{M}^{(n)}(h_1, h_2..., h_n)$ for gluons $1, ..., n$ with momenta $k_1....k_n$ and helicities $h_1....h_n$ are given by

$$\left|\mathcal{M}^{(n)}(+++++..)\right|^2 = c_n(g,N)[0 + O(g^4)] \tag{9}$$

$$\left|\mathcal{M}^{(n)}(-++++..)\right|^2 = c_n(g,N)[0 + O(g^4)] \tag{10}$$

$$\left|\mathcal{M}^{(n)}(--+++..)\right|^2 = c_n(g,N)\left[(k_1\cdot k_2)^4\sum_p[(k_1\cdot k_2)(k_2\cdot k_3)(k_3\cdot k_4)..(k_n\cdot k_1)]^{-1}+O(N^{-2})+O(g^2)\right], \tag{11}$$

where N is the number of colors, $c_n(g,N) = g^{2n-4}N^{n-2}(N^2-1)/2^{n-4}n$ and we are summing over all permutations p of $1,.....,n$. Clearly the first two amplitudes, i.e. the ones that most violate helicity conservation, are zero at tree-level. Furthermore, due to the fact that the gluons in question need to obey Bose-Einstein statistics, any permutation of the gluons leaves the expressions for the amplitude unchanged, i.e. for example $\mathcal{M}^{(n)}(-+++)$ is of the same form as $\mathcal{M}^{(n)}(+-++)$. Also the sign of the helicities involved does not change the form of the expression so that $\mathcal{M}^{(n)}(-+++) = \mathcal{M}^{(n)}(+---)$. For later reference let us note that for a 4-gluon process it therefore follows that the only amplitude which is non-zero at tree-level is (11), as here equations (9),(10) and (11) cover all possible cases.

But for now consider the general case of tree-level scattering of n gauge bosons in Yang-Mills theory, i.e. in the Standard Model. We take the n gauge boson momenta $k_1, k_2...\ k_n$ to all be directed inward. In order to get the MHV amplitude we assume that the two bosons with momenta k_1 and k_2 have negative helicities, whilst the remaining gauge bosons have positive helicities. Furthermore the n gauge bosons are in U(N) gauge group states corresponding to the generators $T^{a_1}, T^{a_2}...T^{a_n}$ respectively, which completely determine their color charges. Then the explicit amplitude (which follows from equation (11) above) for such a maximally helicity violating process is given by [8]:

$$\mathcal{M}^{(n)}_{YM} = ig^{n-2}Tr(T^{a_1}T^{a_2}...T^{a_n})\frac{\langle 12\rangle^4}{\langle 12\rangle\langle 23\rangle...\langle n1\rangle}, \tag{12}$$

where g is the U(N) coupling constant. The spinor product $\langle k_1k_2\rangle$ is defined to be

$$\langle k_1k_2\rangle\ = \overline{u}_-(k_1)u_+(k_2), \tag{13}$$

and the following properties hold

$$\langle k_1k_2\rangle = -\langle k_2k_1\rangle \quad and \quad \langle k_1k_2\rangle\langle k_1k_2\rangle^* = 2k_1k_2\ . \tag{14}$$

Now, taking into account string effects and restricting ourselves to 4 gauge bosons, the amplitude is modified and we get [9]

$$\mathcal{M}^{(4)} = V(k_1,k_2,k_3,k_4)\mathcal{M}^{(4)}_{YM}, \tag{15}$$

where the stringy effect on the amplitude is completely contained in the Veneziano form-factor $V(k_1,k_2,k_3,k_4)$:

$$V(k_1, k_2, k_3, k_4) = \frac{\Gamma(1-s)\Gamma(1-u)}{\Gamma(1+t)}, \tag{16}$$

where s, t and u are the familiar Mandelstam variables normalized to string units

$$s = \frac{2k_1k_2}{M_S^2}, t = \frac{2k_1k_3}{M_S^2}, u = \frac{2k_1k_4}{M_S^2}, \text{and } s+t+u = 0, \tag{17}$$

and the Gamma function is defined to be $\Gamma(z) = \int_0^\infty t^{z-1}e^{-t}dt$, i.e. an extension of the factorial function for an arbitrary complex number z with positive real part, as usual.

Hence we can now write down the full partial MHV amplitude for the scattering of four gauge bosons as appropriate when considering $gg \to \gamma g$. N here equals 4 and we have two gauge bosons (1 and 2) with negative helicities and two (3 and 4) with positive helicities - the only amplitude we need to compute for 4-gluon scattering as shown above. The full partial MHV amplitude therefore is [10]

$$\mathcal{M}_{part}(1^-, 2^-, 3^+, 4^+) = 4g^2 Tr(T^{a_1}T^{a_2}T^{a_3}T^{a_4})\frac{\langle 12\rangle^4}{\langle 12\rangle\langle 23\rangle\langle 34\rangle\langle 41\rangle}V(k_1, k_2, k_3, k_4). \tag{18}$$

Consider the generators $T^{a_1}, T^{a_2}...T^{a_4}$ in this expression. For three SU(N) gluons g_1, g_2, g_3 and one U(1) gauge boson γ_4 they take the following form in the fundamental representation

$$T^{a_1} = T^a \ , \ \ T^{a_2} = T^b \ , \ \ T^{a_3} = T^c \ , \ \ T^{a_4} = QI, \tag{19}$$

where Q is the associated U(1) charge and I is the identity matrix. Using the normalization condition for the U(N) generators

$$Tr(T^aT^b) = \frac{1}{2}\delta^{ab}, \tag{20}$$

the color factor in the expression for the amplitude (18) is

$$Tr(T^{a_1}T^{a_2}T^{a_3}T^{a_4}) = Q(d^{abc} + \frac{i}{4}f^{abc}), \tag{21}$$

where d^{abc} is the symmetrized trace and hence totally symmetric, whereas f^{abc} is the totally antisymmetric structure constant.

Summing the partial amplitudes over all permutations σ of {1,2,3} we obtain the total MHV amplitude

$$\mathcal{M}(g_1^-, g_2^-, g_3^+, \gamma_4^+) = 4g^2\langle 12\rangle^4 \sum_\sigma \frac{Tr(T^{a_{1\sigma}}T^{a_{2\sigma}}T^{a_{3\sigma}}T^{a_4})V(k_{1\sigma}, k_{2\sigma}, k_{3\sigma}, k_4)}{\langle 1_\sigma 2_\sigma\rangle\langle 2_\sigma 3_\sigma\rangle\langle 3_\sigma 4\rangle\langle 41_\sigma\rangle}, \tag{22}$$

where the fact that the summation only runs over permutations in the denominator (except for gauge boson 4 which has to stay fixed since it is required to remain a U(1) gauge boson γ_4) is a consequence of the initial fixing of negative helicities for two gauge bosons and becomes clear when looking back at the general expression (12) for the MHV amplitude above. The antisymmetric part of the color factor $Q\frac{i}{4}f^{abc}$ cancels, since summing over all possible permutations amounts to a symmetrization of the amplitude, so that one eventually obtains

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

where we have defined what is essentially the difference between two Veneziano amplitudes

$$\mu(s,t,u) = \Gamma(1-u) \left(\frac{\Gamma(1-s)}{\Gamma(1+t)} - \frac{\Gamma(1-t)}{\Gamma(1+s)} \right) = V(s,t,u) - V(t,s,u). \tag{24}$$

All other non vanishing amplitudes, e.g. $\mathcal{M}(g_1^-, g_2^+, g_3^-, \gamma_4^+)$, can be calculated analogously.

In order to work out the total cross-section for the (unpolarized) process $gg \to \gamma g$ we take the squared moduli of all those individual non-vanishing (polarized) amplitudes, sum over final polarization and color states and average over all possible initial polarization and color states. The resulting total average squared amplitude is

$$|\mathcal{M}(gg \to g\gamma)|^2 = g^4Q^2 \frac{2(N^2-4)}{N(N^2-1)} \left\{ \left| \frac{s\mu(s,t,u)}{u} + \frac{s\mu(s,u,t)}{t} \right| + (s \leftrightarrow t) + (s \leftrightarrow u) \right\}. \tag{25}$$

Now the low-energy expansion for the Veneziano form-factor, i.e. for when $(s,t,u \ll 1)$, is [10]

$$V(s,t,u) \approx 1 - \frac{\pi^2}{6} su - \zeta(3) stu + ... \tag{26}$$

Hence the expression for the total average amplitude squared at low energies becomes

$$|\mathcal{M}(gg \to g\gamma)|^2 \approx g^4Q^2 \frac{2(N^2-4)}{N(N^2-1)} \frac{\pi^4}{4} (s^4 + t^4 + u^4), \tag{27}$$

where the fact that no singularities are present indicates that no exchange of massless particles contributes to this process. However, close to the string threshold, i.e. $s \approx 1$ in the normalized units we have been using, or $s \approx M_S^2$ when we restore the string scale, the amplitude expression becomes

$$|\mathcal{M}(gg \to g\gamma)|^2 \approx 4g^4Q^2 \frac{2(N^2-4)}{N(N^2-1)} \frac{M_S^8 + t^4 + u^4}{M_S^4\left[(s-M_S^2)^2 + (\Gamma M_S)^2\right]}, \tag{28}$$

where we now do have a singularity with the variable Γ parameterizing its smearing to a Breit-Wigner form due to the finite decay width resonances propagating in the s channel. In particular, the singularity itself results from the presence of a massive string mode propagating in the s channel.

In order to fix the value of Q in the above calculations we need to take into account the mixing effects between the different U(1) boson fields at play. If the process under consideration was $gg \to gg$ or $gg \to Cg$ with C a U(1) gauge field directly tied to the color U(3) brane, Q would be fixed at $\sqrt{1/6}$ due to the normalization condition on the trace of the generators (20). For $gg \to \gamma g$ we need to incorporate two further projections, however. Firstly from the "extra" U(1) boson C_μ to the familiar electroweak hypercharge boson Y_μ, resulting in an additional mixing factor κ. And secondly from Y_μ to the photon A_μ contributing a Weinberg factor $\cos\theta_W$. Thus the extra stringy contribution is the C-Y mixing term κ which, however, turns out to be model-dependent. In the minimal model we are considering here it takes a value of $\kappa \sim 0.12$, but this value can vary considerably due to e.g. extra U(1) gauge bosons partnering the $SU(2)_L$ electroweak gauge bosons W^a_μ on a separate U(2) brane. Here we will take $\kappa^2 = 0.02$ which results in the following value for Q

$$Q^2 = \frac{1}{6}\kappa^2 \cos^2\theta_W \simeq 2.55 \times 10^{-3} \quad (\kappa^2/0.02). \tag{29}$$

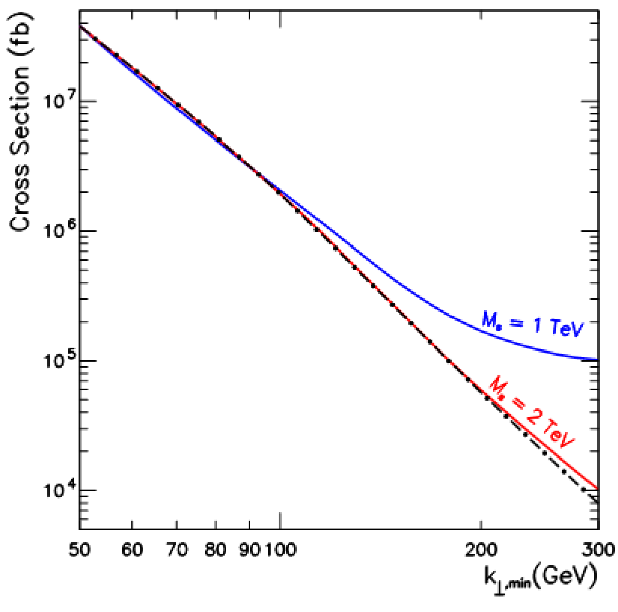

Figure 1: The QCD (dot dashed line) and QCD + string (solid lines) cross sections for two values of M_S

3.3 Discovery contours : Cross-sections and Signal-to-Noise ratios

Having calculated the relevant amplitudes and employing the already known parton distribution functions [11] it is now possible to extract integrated cross-sections for $gg \to \gamma g$ and for other QCD background processes which also give rise to $pp \to \gamma + jet$. Employing the minimal model sketched in 3.1, we will consider interactions on a color U(3) stack of D-branes and hence take $N = 3$. g is taken to be the QCD coupling constant and hence $g^2/4\pi = 0.1$. Furthermore, $\Gamma \simeq (g^2/16\pi)\,(2j+1)^{-1}\,M_S$ where $j = 2$.

A particularly well suited signal in looking for $gg \to \gamma g$ contributions to $pp \to \gamma + jet$ are high $k_\perp$ isolated γ or Z. One can therefore calculate the integrated cross sections $\sigma(pp \to \gamma + jet)|_{k_\perp(\gamma) > k_{\perp,\min}}$. The results for different string scales are shown in Fig 1. Clearly larger values of $k_{\perp,\min}$ give the best signal, as the background is reduced significantly in these regions. When imposing a 300 GeV cut on the transverse momentum the QCD background cross section is approx. $8 \times 10^3 fb$ which equates to about 8×10^5 events for a luminosity of $100 fb^{-1}$.

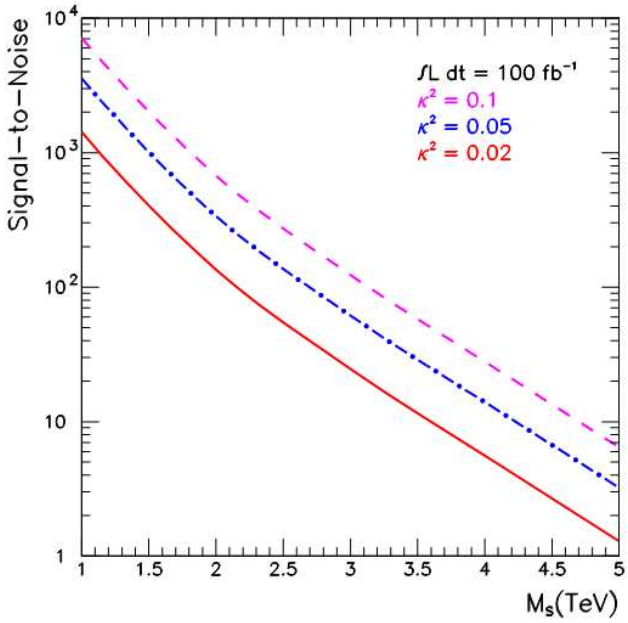

Figure 2: Signal-to-noise ratios for different values of the mixing parameter k and an integrated luminosity of $100 fb^{-1}$

Fig. 2 shows the signal-to-noise ratio $\left(signal/\sqrt{SM\ background}\right)$ for different values of κ and for an integrated luminosity of $100 fb^{-1}$. We can see that a 5σ discovery is possible in principle for $M_S \leqslant 4 TeV$. One should also note that isolation cuts need to be imposed on the photon to minimize misidentification with e.g. high-$k_\perp$ neutral pions.

Unfortunately this reach is lowered slightly when treating the resonance region more carefully than we have here e.g. by computing the decay widths for $J = 0$ and $J = 2$ Regge recurrences of the gluon octet [12]. The fact that these resonances are slightly wider than expected has a lowering effect on the cross section thus leading to a reduced discovery reach.

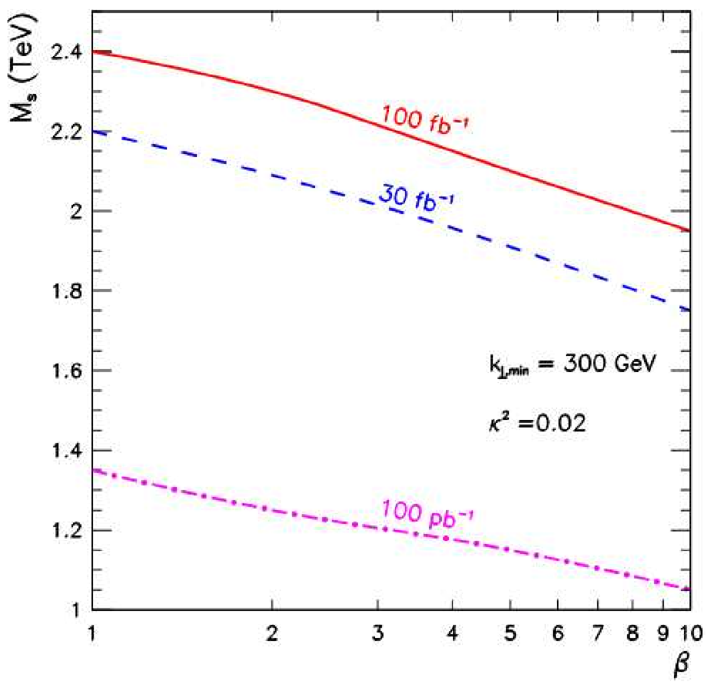

Figure 3: 5σ discovery contours for different integrated luminosities in the (M_S,β) plane

Let us now define the following parameter:

$$\beta = \frac{background\ due\ to\ misidentified\ \pi^0\ after\ isolation\ cuts}{QCD\ background\ from\ direct\ photon\ production} + 1. \tag{30}$$

The noise from direct photon QCD production is therefore increased by a factor of $\sqrt{\beta}$. Taking into account the refined cross-sections and different integrated luminosities we can now plot the 5σ discovery contours in the (string scale, β) plane. As can be seen in Fig. 3, the expected LHC discovery reach should now run up to 2.3 TeV. The $100pb^{-1}$ luminosity is plotted to compare the discovery reach in detectors like CMS and ATLAS to that of detectors observing e.g. Pb-Pb collisions like ALICE, where, however, the relevant luminosities are presently estimated to lie significantly below $100pb^{-1}$ [13].

4 Conclusion

The emergence of TeV scale string theories has given rise to the possibility of a number of string-related experimental signals at the LHC. The $gg \to \gamma g$ process, the dominant underlying parton process for the excitation of string resonances in $pp \to \gamma + jet$, provides a particularly promising signal in the search for Regge excitations of fundamental strings, as it is both not present at tree-level in the standard model and also model-independent except for the precise value of the mixing parameter κ. It consequently provides us with the opportunity to decisively test the validity of low mass string theories up to a certain scale. One should keep in mind, however, what it takes for these signals to be observable at the LHC. In short one needs a low string scale, large extra dimensions and a weak string coupling. If such searches for low mass string signatures do indeed turn out to be successful, analogous processes such as four-fermion interactions (e.g. quark-antiquark scattering) can also be used as precision tests on low mass string models, as such processes depend on compactification details via the exchange of heavy Kaluza-Klein and winding states. Finally one should note that the above estimates do therefore not include compactification-dependent effects of string-related corrections to SM processes. Nevertheless such contributions might in fact almost double the LHC's reach in searching for low mass strings and should therefore certainly not be neglected.

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The image on the cover-page and Fig. 3 are taken from [12]. Fig 1 is taken from [10]. Fig. 2 is taken from [5].

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