MANIFESTATIONS OF STRING THEORY
IN ASTROPHYSICAL DATA AND AT THE LHC

By

Satoshi Nawata

A DISSERTATION SUBMITTED IN
PARTIAL FULFILLMENT OF THE
REQUIREMENTS FOR THE DEGREE OF

DOCTOR OF PHILOSOPHY

IN
PHYSICS

at
The University of Wisconsin–Milwaukee
August 2008
With the advent of the LHC and the continuing influx of cosmological data, phenomenological aspects of string theory have received renewed attention in recent years and many problems have been properly incorporated in this framework. For instance, recent theoretical considerations in string theory have applied a statistical approach to the enormous landscape of metastable vacua. The large number of vacua may shed some light on the cosmological constant problem. In addition, in string theory, attempts have been made to address the hierarchy problem within the context of the existence of large or warped internal dimensions transverse to a braneworld where we are confined, which lowers the effective scale of gravity to the TeV region. If this were the case, unseen dimensions of the space-time can be at the border of the energy domain within reach of the next generation of particle accelerators.
Although the picture of the landscape may be the key to the cosmological constant problem, it is well-known that the compactification of a string background to a four dimensional solution undergoing accelerating expansion is difficult, which is described by the no-go theorem of Maldacena-Nuñez. In the first part of this Dissertation, we investigate the cosmological content of the Salam-Sezgin supergravity which circumvents one of the hypotheses of the no-go theorem of Maldacena-Nuñez and consequently can support a de Sitter phase when lifted to string theory. We find a solution to the field equations in qualitative agreement with the observed dark energy density. The carrier of the acceleration in the present de Sitter epoch is a quintessence field slowly rolling down its exponential potential. Intrinsically to this model, there is a second modulus, which is automatically stabilized and acts as a source of cold dark matter with a mass proportional to an exponential function of the quintessence field.

In the second part, we explore a "new physics" signal at LHC, in the processes $pp \rightarrow \gamma + \text{jet}$, $pp \rightarrow \gamma \gamma$, and $pp \rightarrow \text{dijet}$. In D-brane quivers, there are one or more additional $U(1)$ gauge symmetries, beyond the $U(1)_Y$ of the standard model, which 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 $N > 2$. Because of this, the photon will participate with the $SU(N)$ gauge boson in string tree level scattering processes which in the standard model occur only at one-loop level. In order to evaluate this stringy correction, we considered the processes $gg \rightarrow g\gamma$ and $gg \rightarrow \gamma\gamma$, and found that cross section measurements of the process $pp \rightarrow \gamma + \text{jet}$ at the LHC will attain $5\sigma$ discovery reach on low scale string models for $M_{\text{string}}$ as large as 4 TeV. We have also considered the processes $gg \rightarrow gg$, $gg \rightarrow q\bar{q}$, $q\bar{q} \rightarrow gg$, $gg \rightarrow gg$ and $q\bar{g} \rightarrow q\bar{q}$, and found that, for the $pp \rightarrow \text{dijet}$ channel, the LHC discovery reach will extend up to $M_{\text{string}} \sim 6.8$ TeV.
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Preface

This Dissertation is based on four papers that explore the rich interdisciplinary boundaries of string theory, particle physics, and cosmology.

The cosmological set-up presented in Chapter 2 is based on:

L. Anchordoqui, H. Goldberg, S. Nawata and C. Nuñez, Cosmology from String Theory, Phys. Rev. D 76, 126005 (2007).

Chapters 3 and 4 are based on material from:

L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Jet signals for low mass strings at the Large Hadron Collider, Phys. Rev. Lett. 100, 171603 (2008).

and

L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Direct photons as probes of low mass strings at the CERN LHC, Phys. Rev. D 78, 016005 (2008).

The ideas discussed in Chapter 5 are based on:

L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Nawata, S. Stieberger, and T. R. Taylor, Dijet signals of low mass strings at the LHC, arXiv:0808.0497 [hep-ph].
Acknowledgments

First and foremost, I would like to express my deep gratitude to my adviser, Professor Luis Anchordoqui. His patience, educational guidance, time and knowledge were paramount to my work within the last 2 years and this Dissertation would certainly not have been possible without him. Further, his willingness to introduce me to so many distinguished colleagues in the field of particle physics, notably Haim Goldberg, Dieter Lüst, Carlos Nuñez, Stephan Stieberger and Tomasz Taylor is also very much appreciated. The papers we have collaborated on have already allowed me to hone my research and analytical skills, thereby opening other avenues to future research and projects.

I am also thankful for the education and guidance from Distinguished Professor John Friedman and Distinguished Professor Leonard Parker, as well as the opportunity to work with everyone in the Center for Gravitation and Cosmology.

I have also benefited greatly from the time and guidance provided to me by Professor Richard Sorbello, in his role as the Graduate Program Advisor, and from Ms. Kate Valerius, Graduate Program Assistant, whose level of patience and willingness to help in so many of the non-academic areas (keeping me organized, teaching communication and writing skills, etc.) has truly been appreciated and will be remembered for a long time to come.
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Chapter 1

Introduction

Elementary Particle Physics seeks to understand, at the deepest level, the structure of matter and the forces by which it interacts. The experimental success of the standard model (SM) of weak, electromagnetic, and strong interactions can be considered as the triumph of the gauge symmetry principle to describe all physical phenomena up to energies $\sim 500$ GeV [1]. However, the SM remains unsatisfactory in some of its theoretical aspects. The major one concerns quantum gravity effects: the renormalization procedure that allows one to extract finite predictions for processes involving the three other fundamental forces fails when gravitational interactions are taken into account. String theory stands here as the only known consistent framework to incorporate these effects, replacing the elementary point particles (which form matter and mediate interactions) with a single extended object of vanishing width [2, 3]. The known fundamental particles appear “point-like” because the experimental energies probed thus far by colliders are too small to excite the string oscillation modes, so only the center of mass motion is perceived. In addition to these heavy oscillation modes, strings have new degrees of freedom that often take the classical geometry description of propagation into “hidden” compact dimensions, recovering the old ideas of Kaluza and Klein [4, 5].

The distance at which quantum gravity comes into play is unknown. Though this distance can be very small, $O(10^{-35} \text{ m})$, a particularly interesting possibility arises if it is $O(10^{-19} \text{ m})$, the distance at which the electromagnetic and weak forces are known to unify to form the electroweak force. Lowering the scale of quantum gravity into the TeV region provides a framework for solving the mass hierarchy problem and unifying all interactions. The apparent weakness of gravity can be accounted for by the existence of large [6] or warped [7] internal dimensions transverse to a braneworld where we are confined. If this were the case, spacetime’s unseen dimensions could be at the border of the energy domain...
within reach of the next generation of particle accelerators, beginning this year with the Large Hadron Collider (LHC) at CERN. In particular, the mass scale $M_s$ of fundamental strings would be as low as few TeV \[^{[8]}\]. This mass determines the center of mass energy threshold $\sqrt{\hat{s}} \geq M_s$ for the production of Regge resonances in parton collisions, thus for the onset of string effects at the LHC \[^{[9]}\]. In this Dissertation 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 \[^{[10]}\]. 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 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 we show in this Dissertation, 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.

On a separate track, the remarkable accuracy of the Wilkinson Microwave Anisotropy Probe (WMAP) five-year observations has catapulted us into a new era in cosmology \[^{[11]}\]. This information, taken together with Big Bang nucleosynthesis (BBN) abundances \[^{[12]}\], astronomical observations charting the large scale distribution of galaxies \[^{[13, 14]}\], and luminosity distance measurements of Type Ia supernovae \[^{[15, 16, 17]}\], seems to ensure the existence of some unknown form of “dark” energy density that dominates the recent gravitational dynamics of the universe and yields a stage of cosmic acceleration. Moreover, these impressive experiments have weighed the universe and determined that the known particles make up only 5% of its mass, providing overwhelming evidence for new particles and fundamental laws of nature. The discovery of the two unknown components of the universe, the so-called “dark matter” and dark energy, pose an important challenge for particle physics that would be met in a few years by the LHC \[^{[18]}\]. Dark matter appears to consist of non-relativistic particles that only interact gravitationally and perhaps by weak interactions. The nature of the second unknown component is even less clear. The simplest candidate for such a missing energy is a positive cosmological constant $\Lambda$. Such an identification, however, unavoidably raises a series of questions: (a) Why is $\Lambda$ so small in particle physics units? – the so-called fine-tuning problem for $\Lambda$; (b) Why is $\Lambda \sim$
the present value (in Planck units) of the (dark) matter density? – the so-called “coincidence problem.” At present the most promising scenarios for answering (at least part of) these questions associate dark energy with a dynamical scalar field \[19\], generally called “quintessence” \[20\], whose potential goes to zero asymptotically (leaving therefore just the usual puzzle of why the “true” cosmological constant vanishes). The scalar field slowly rolls down such a potential reaching infinity (and zero energy) only after an infinite (or very long) time. While doing so quintessence produces an effective, time-dependent, cosmic energy density accompanied by a sufficiently negative pressure, i.e., an effective cosmological constant. As far as identifying quintessence is concerned several possibilities have been considered (see e.g., \[21\][22][23][24][25][26][27][28]; in particular, those motivated by string theory include the dilaton and time-varying moduli fields \[29\]. In addition to analyzing for new signals at the LHC, this Dissertation is aimed at investigating a number of topics that exploit the recent astrophysical observations mentioned above. Our overall goal is of furthering the understanding of cosmology while simultaneously investigating new areas of fundamental physics.

The Dissertation is organized as follows. In the next chapter, we study the cosmological content of Salam-Sezgin \[30\] six dimensional supergravity, which circumvents the no-go theorem of Maldacena-Nuñez when the six dimensional space time is embedded in Type I or Heterotic supergravity, in both cases with non-compact extra-dimensions: \( \mathcal{H}^{2,2} \times S^1 \). We solve the field equations matching the existing experimental data, and we find candidates for the quintessence field and cold dark matter, which can be written as linear combinations of the \( S^2 \) moduli field and the six dimensional dilaton field. In Chapters 3, 4, and 5 we discuss phenomenological aspects of low mass string theory related to experimental searched for physics beyond the SM at the LHC. Using a generic property of D-brane quiver models, we search for Regge recurrences at parton collision energies \( \sqrt{s} \sim M_s \). In these models, the photon mixes with the hypercharge and the gauge field of baryon number; hence some processes like \( gg \rightarrow g\gamma \) or \( gg \rightarrow \gamma\gamma \) show up at the string disk level. In Chapter 3 we analyze \( pp \rightarrow \gamma + \text{jet} \) and \( pp \rightarrow \gamma\gamma \) channels. By setting a high energy cut (300 GeV) on the transverse momentum of the photon, we find a signal at the LHC which could probe deviation from the SM at a 5\( \sigma \) significance for \( M_s \) as large as 2.3 TeV. After that, in Chapter 4 we study additional LHC observables and discuss potential methods to discriminate string resonances from other sources of new physics. We first study the invariant mass spectrum of string resonances in the \( pp \rightarrow \gamma + \text{jet} \) channel and find that bump searches could help to increase the LHC discovery reach up to \( M_s \sim 4 \) TeV. Then, we compare \( \gamma \) and \( Z \) production via Hawking evaporation of TeV-scale back holes and string excitations of D-brane models.
In Chapter 5, we extend our search of string signals at the LHC by analyzing the process $pp \rightarrow \text{dijet}$. We find that this channel exhibits a resonant behavior at a $5\sigma$ significance for $M_s$ as large as 6.8 TeV. Chapter 6 contains the main conclusions of this Dissertation. In Appendix A, we present the proof of the no-go theorem of Maldacena-Núñez. In Appendix B we discuss the the connection between Salam-Sezgin six dimensional supergravity and string theory. In Appendices C and D, we collect calculations that are somewhat technical for the main text.
Chapter 2

String Cosmology

The mechanism involved in generating a very small cosmological constant that satisfies ’t Hooft naturalness is one of the most pressing questions in contemporary physics. Recent observations of distant Type Ia supernovae [15, 16, 17] strongly indicate that the universe is expanding in an accelerating phase, with an effective de-Sitter (dS) constant $H$ that nearly saturates the upper bound given by the present-day value of the Hubble constant, i.e., $H \leq H_0 \sim 10^{-33}$ eV. According to the Einstein field equations, $H$ provides a measure of the scalar curvature of the space and is related to the vacuum energy density $\rho_{\text{vac}}$ through Friedmann’s equation, $3 M_{\text{Pl}}^2 H^2 \sim \rho_{\text{vac}}$, where $M_{\text{Pl}} \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass. However, the “natural” value of $\rho_{\text{vac}}$ coming from the zero-point energies of known elementary particles is found to be at least $\rho_{\text{vac}} \sim \text{TeV}^4$. Substitution of this value of $\rho_{\text{vac}}$ into Friedmann’s equation yields $H \geq 10^{-3}$ eV, grossly inconsistent with the set of supernova (SN) observations. The absence of a mechanism in agreement with ’t Hooft naturalness criteria then centers on the following question: why is the vacuum energy needed by the Einstein field equations 120 orders of magnitude smaller than any “natural” cut-off scale in effective field theory of particle interactions, but not zero?

Today, the most popular framework that can address aspects of this question is the anthropic approach, in which the fundamental constants are not determined through fundamental reasons, but rather because such values are necessary for life (and hence intelligent observers to measure the constants) [31]. Of course, in order to implement this idea in a concrete physical theory, it is necessary to postulate a multiverse in which fundamental physical parameters can take different values. Recent investigations in string theory have applied a statistical approach to the enormous “landscape” of metastable vacua present in the theory [32, 33, 34, 35]. A vast ensemble of metastable vacua with a small positive effective cosmological constant that can accommodate the low energy effective field theory of
the SM has been found. Therefore, the idea of a string landscape has been used to propose a concrete implementation of the anthropic principle.

Nevertheless, the compactification of a string/M-theory background to a four dimensional solution undergoing accelerating expansion has proved to be exceedingly difficult. The obstruction to finding dS solutions in the low energy equations of string/M theory is well known and summarized in the no-go theorem of [36, 37]. This theorem states that in a $D$-dimensional theory of gravity, in which (a) the action is linear in the Ricci scalar curvature, (b) the potential for the matter fields is non-positive, and (c) the massless fields have positive defined kinetic terms, there are no (dynamical) compactifications of the form: $ds^2_D = \Omega^2(y)(dx_d^2 + \hat{g}_{mn}dy^m dy^n)$, if the $d$ dimensional space has Minkowski $SO(1, d-1)$ or dS $SO(1, d)$ isometries and its $d$ dimensional gravitational constant is finite (i.e., the internal space has finite volume). Further details are given in the Appendix A. The conclusions of the theorem can be circumvented if some of its hypotheses are not satisfied. Examples where the hypotheses can be relaxed exist: (i) one can find solutions in which not all of the internal dimensions are compact [38]; (ii) one may try to find a solution breaking Minkowski or de Sitter invariance [39]; (iii) one may try to add negative tension matter (e.g., in the form of orientifold planes) [40]; (iv) one can even appeal to some intricate string dynamics [41].

The Salam-Sezgin six dimensional supergravity model [30] provides a specific example where the no-go theorem is not at work, because when their model is lifted to M theory the internal space is found to be non-compact [42] (See Appendix B). The lower dimensional perspective of this, is that in six dimensions the potential can be positive. This model has perhaps attracted the most attention because of the wide range of its phenomenological applications (see e.g., [43, 44, 45, 46, 47]). In this chapter we examine the cosmological implications of such a supergravity model during the epochs subsequent to primordial nucleosynthesis. We derive a solution of Einstein field equations that is in qualitative agreement with luminosity distance measurements of Type Ia supernovae [15, 16, 17], primordial nucleosynthesis abundances [12, 48], data from the Sloan Digital Sky Survey (SDSS) [49], and the most recent measurements from the Wilkinson Microwave Anisotropy Probe (WMAP) satellite [11]. The observed acceleration of the universe is driven by the “dark energy” associated to a scalar field slowly rolling down its exponential potential (i.e., kinetic energy density $< \text{potential energy density} \equiv \text{negative pressure}$). Very interestingly, the resulting cosmological model also predicts a cold dark matter (CDM) candidate. In analogy with the phenomenological proposal of [50, 51], such a nonbaryonic matter interacts with the dark energy field and therefore the mass of the CDM particles evolves with the exponential dark
energy potential. However, an attempt to saturate the present CDM component in this manner leads to gross deviations from present cosmological data. We will show that this type of CDM can account for up to about 7% of the total CDM budget. Generalizations of our scenario (using supergravities with more fields) might account for the rest.

### 2.1 Salam-Sezgin Cosmology

We begin with the action of Salam-Sezgin six dimensional supergravity [30], setting to zero the fermionic terms in the background (of course fermionic excitations will arise from fluctuations),

$$S = \frac{1}{4\kappa^2} \int d^6x \sqrt{g_6} \left[ R - \kappa^2 (\partial M \sigma)^2 - \kappa^2 e^{\kappa \sigma} F_{MN}^2 - \frac{2g^2}{\kappa^2} e^{-\kappa \sigma} - \frac{\kappa^2}{3} e^{2\kappa \sigma} G_{MNP}^2 \right]. \quad (2.1)$$

Here, $g_6 = \det g_{MN}$, $R$ is the Ricci scalar of $g_{MN}$, $F_{MN} = \partial [M A_N]$, $G_{MNP} = \partial [M B_{NP}] + \kappa A_{[M} F_{NP]}$, and capital Latin indices run from 0 to 5. A re-scaling of the constants: $G_6 \equiv 2\kappa^2$, $\phi \equiv -\kappa \sigma$ and $\xi \equiv 4g^2$ leads to

$$S = \frac{1}{2G_6} \int d^6x \sqrt{g_6} \left[ R - \partial (\partial \phi)^2 - \frac{\xi}{G_6} e^{\phi} - \frac{\xi}{2} e^{\phi} F_{MN}^2 - \frac{\xi}{6} e^{-2\phi} G_{MNP}^2 \right]. \quad (2.2)$$

The length dimensions of the fields are: $[G_6] = L^4$, $[\xi] = L^2$, $[\phi] = [g_{MN}] = 1$, $[A_M^2] = L^{-4}$, and $[F_{MN}^2] = [G_{MNP}^2] = L^{-6}$.

Now, we consider a spontaneous compactification from six dimension to four dimension. To this end, we take the six dimensional manifold $M$ to be a direct product of 4 Minkowski directions (hereafter denoted by $N_1$) and a compact orientable two dimensional manifold $N_2$ with constant curvature. Without loss of generality, we can set $N_2$ to be a sphere $S^2$, or a $\Sigma_2$ hyperbolic manifold with arbitrary genus. The metric on $M$ locally takes the form

$$ds_6^2 = ds_4(t, \vec{x})^2 + e^{2f(t, \vec{x})} d\sigma^2, \quad d\sigma^2 = \begin{cases} \frac{r^2}{e^f} (d\vartheta^2 + \sin^2 \vartheta d\varphi^2) & \text{for } S^2 \\ \frac{r^2}{e^f} (d\vartheta^2 + \sinh^2 \vartheta d\varphi^2) & \text{for } \Sigma_2 \end{cases}, \quad (2.3)$$

where $(t, \vec{x})$ denotes a local coordinate system in $N_1$, $r_c$ is the compactification radius of $N_2$. We assume that the scalar field $\phi$ is only dependent on the point of $N_1$, i.e., $\phi = \phi(t, \vec{x})$. We further assume that the gauge field $A_M$ is excited on $N_2$ and is of the form

$$A_\varphi = \begin{cases} b \cos \vartheta & (S^2) \\ b \cosh \vartheta & (\Sigma_2) \end{cases}. \quad (2.4)$$

This is the monopole configuration detailed by Salam-Sezgin [30]. Since we set the Kalb-Ramond field $B_{NP} = 0$ and the term $A_{[M} F_{NP]}$ vanishes on $N_2$, $G_{MNP} = 0$. The field strength becomes

$$F_{MN}^2 = 2b^2 e^{-4f}/r_c^4. \quad (2.5)$$
Taking the variation of the gauge field $A_M$ in Eq. (2.2) we obtain the Maxwell equation

$$\partial_M \left[ \sqrt{g_4 \sqrt{g_6} e^{2f - \phi} F^{MN}} \right] = 0.$$  

(2.6)

It is easily seen that the field strengths in Eq. (2.5) satisfy Eq. (2.6).

With this in mind, the Ricci scalar reduces to

$$R[M] = R[N_1] + e^{-2f} R[N_2] - 4\Box f - 6(\partial_\mu f)^2,$$  

(2.7)

where $R[M]$, $R[N_1]$, and $R[N_2]$ denote the Ricci scalars of the manifolds $M$, $N_1$, and $N_2$; respectively. (Greek indices run from 0 to 3). The Ricci scalar of $N_2$ reads

$$R[N_2] = \begin{cases} 
+2/r_c^2 & (S^2) \\
-2/r_c^2 & (\Sigma_2).
\end{cases}$$  

(2.8)

To simplify the notation, from now on, $R_1$ and $R_2$ indicate $R[N_1]$ and $R[N_2]$, respectively.

The determinant of the metric can be written as

$$\sqrt{g_4} = e^{2f} \sqrt{g_6 \sqrt{g_4}} = e^{2f} \sqrt{g_6}.$$

Hence, by using the field configuration given in Eq. (2.4), we can re-write the action in Eq. (2.2) as follows

$$S = \frac{1}{G_4} \int d^4x \sqrt{g_4} \left\{ e^{2f} \left[ R_1 + e^{-2f} R_2 + 2(\partial_\mu f)^2 - (\partial_\mu \phi)^2 \right] - \frac{\xi}{G_6} e^{2f + \phi} - \frac{G_6 b^2}{r_c^4} e^{-2f - \phi} \right\}.$$  

(2.10)

Let us consider now a rescaling of the metric of $N_1$: $\hat{g}_{\mu\nu} \equiv e^{2f} g_{\mu\nu}$ and $\sqrt{\hat{g}_4} = e^{4f} \sqrt{g_4}$. Such a transformation brings the theory into the Einstein conformal frame where the action given in Eq. (2.10) takes the form

$$S = \frac{1}{G_4} \int d^4x \sqrt{\hat{g}_4} \left[ R[\hat{g}_4] - 4(\partial_\mu f)^2 - (\partial_\mu \phi)^2 - \frac{\xi}{G_6} e^{2f + \phi} - \frac{G_6 b^2}{r_c^4} e^{-6f - \phi} + e^{-4f} R_2 \right].$$  

(2.11)

The four dimensional Lagrangian is then

$$L = \sqrt{g} \frac{G_4}{G_6} \left[ R - 4(\partial_\mu f)^2 - (\partial_\mu \phi)^2 - V(f, \phi) \right],$$  

(2.12)

with

$$V(f, \phi) \equiv \frac{\xi}{G_6} e^{2f + \phi} + \frac{G_6 b^2}{r_c^4} e^{-6f - \phi} - e^{-4f} R_2,$$  

(2.13)

where to simplify the notation we have defined: $g \equiv \hat{g}_4$ and $R \equiv R[\hat{g}_4]$. 
Let us now define a new orthogonal basis, \( X \equiv (\phi + 2f)/\sqrt{G_4} \) and \( Y \equiv (\phi - 2f)/\sqrt{G_4} \), so that the kinetic energy terms in the Lagrangian are both canonical, i.e.,

\[
L = \sqrt{g} \left[ \frac{R}{G_4} - \frac{1}{2} (\partial X)^2 - \frac{1}{2} (\partial Y)^2 - \tilde{V}(X,Y) \right],
\]

where the potential \( \tilde{V}(X,Y) \equiv V(f,\phi)/G_4 \) can be re-written (after some elementary algebra) as \[53\]

\[
\tilde{V}(X,Y) = e^{\sqrt{G_4}X} \left[ \frac{G_6 h^2}{r^4_c} e^{-2\sqrt{G_4}X} - R_2 e^{-\sqrt{G_4}X} + \frac{\xi}{G_6} \right].
\]

The field equations are

\[
R_{\mu \nu} - \frac{1}{2} g_{\mu \nu} R = \frac{G_4}{2} \left[ \left( \partial_{\mu} X \partial_{\nu} X - \frac{g_{\mu \nu}}{2} \partial_{\eta} X \partial_{\eta} X \right) + \left( \partial_{\mu} Y \partial_{\nu} Y - \frac{g_{\mu \nu}}{2} \partial_{\eta} Y \partial_{\eta} Y \right) - g_{\mu \nu} \tilde{V}(X,Y) \right],
\]

(2.16)

\( \Box X = \partial X \tilde{V} \), and \( \Box Y = \partial Y \tilde{V} \). In order to allow for a dS era, we assume that the metric takes the form

\[
ds^2 = -dt^2 + e^{2h(t)} dx^2,
\]

(2.17)

and that \( X \) and \( Y \) depend only on the time coordinate, i.e., \( X = X(t) \) and \( Y = Y(t) \). Then, the equations of motion for \( X \) and \( Y \) can be written as

\[
\ddot{X} + 3h \dot{X} = -\partial X \tilde{V}
\]

(2.18)

and

\[
\ddot{Y} + 3h \dot{Y} = -\partial Y \tilde{V},
\]

(2.19)

whereas the only two independent components of Eq. (2.16) are

\[
\dot{h}^2 = \frac{G_4}{6} \left[ \frac{1}{2} (\dot{X}^2 + \dot{Y}^2) + \tilde{V}(X,Y) \right]
\]

(2.20)

and

\[
2\ddot{h} + 3\dot{h}^2 = \frac{G_4}{2} \left[ -\frac{1}{2} (\dot{X}^2 + \dot{Y}^2) + \tilde{V}(X,Y) \right].
\]

(2.21)

The terms in the square brackets in Eq. (2.15) take the form of a quadratic function of \( e^{-\sqrt{G_4}X} \). This function has a global minimum at \( e^{-\sqrt{G_4}X_0} = R_2 r^4_c / (2 G_6 b^2) \). Indeed, the necessary and sufficient condition for a minimum is that \( R_2 > 0 \), so hereafter we only consider the spherical compactification, where \( e^{-\sqrt{G_4}X_0} = M_{Pl}^2 / (4 \pi b^2) \). The condition for the potential to show a dS rather than an AdS or Minkowski phase is \( \xi b^2 > 1 \). Now, we expand Eq. (2.15) around the minimum,

\[
\tilde{V}(X,Y) = \frac{e^{\sqrt{G_4}Y}}{G_4} \left[ \mathcal{K} + \frac{M_{Pl}^2}{2} (X - X_0)^2 + \mathcal{O} \left( (X - X_0)^3 \right) \right],
\]

(2.22)
where

$$M_X \equiv \frac{1}{\sqrt{\pi}} \frac{b r_c}{b r_c}$$

(2.23)

and

$$K \equiv \frac{M_{Pl}^2}{4\pi^2 b^2}(b^2 \xi - 1).$$

(2.24)

As shown by Salam-Sezgin [30], the requirements for preserving a fraction of supersymmetry (SUSY) in spherical compactifications to four dimensions imply \(b^2 \xi = 1\), corresponding to winding number \(n = \pm 1\) for the monopole configuration. Consequently, a (\(Y\)-dependent) dS background can be obtained only through SUSY breaking. For now, we will leave open the symmetry breaking mechanism and come back to this point after our phenomenological discussion. The \(Y\)-dependent physical mass of the \(X\)-particles at any time is

$$M_X(Y) = e^{\sqrt{G_4} Y/2} \frac{M_X}{\sqrt{G_4}} ,$$

(2.25)

which makes this a varying mass particle (VAMP) model [50, 51], although, in this case, the dependence on the quintessence field is fixed by the theory. The dS (vacuum) potential energy density is

$$V_Y = e^{\sqrt{G_4} Y} \frac{G_4}{G_4} K .$$

(2.26)

In general, classical oscillations for the \(X\) particle will occur for

$$M_X > H = \sqrt{\frac{G_4 \rho_{tot}}{3}} ,$$

(2.27)

where \(\rho_{tot}\) is the total energy density. (This condition is well known from axion cosmology [54]). A necessary condition for this to hold can be obtained by saturating \(\rho\) with \(V_Y\) from Eq. (2.26) and making use of Eqs. (2.23) to (2.27), which leads to \(\xi b^2 < 7\). Of course, as we stray from the present into an era where the dS energy is not dominant, we must check at every step whether the inequality (2.27) holds. If the inequality is violated, the \(X\)-particle ceases to behave like CDM.

In what follows, some combination of the parameters of the model will be determined by fitting present cosmological data. To this end we assume that SM fields are confined to \(N_1\) and we denote with \(\rho_{rad}\) the radiation energy, with \(\rho_X\) the matter energy associated with the \(X\)-particles, and with \(\rho_{mat}\) the remaining matter density. With this in mind, Eq. (2.19) can be re-written as

$$\ddot{Y} + 3H \dot{Y} = -\frac{\partial V_{eff}}{\partial Y} ,$$

(2.28)

where \(V_{eff} \equiv V_Y + \rho_X\) and \(H\) is defined by the Friedmann equation

$$H^2 \equiv \dot{h}^2 = \frac{1}{3M_{Pl}^2} \left[ \frac{1}{2} \dot{Y}^2 + V_{eff} + \rho_{rad} + \rho_{mat} \right] ,$$

(2.29)
(Note that the matter energy associated to the \(X\) particles is contained in \(V_{\text{eff}}\).)

It is more convenient to consider the evolution in \(u \equiv -\ln(1+z)\), where \(z\) is the redshift parameter. As long as the oscillation condition is fulfilled, the VAMP CDM energy density is given in terms of the \(X\)-particle number density \(n_X\)

\[
\rho_X(Y, u) = M_X(Y) \ n_X(u) = C \ e^{\sqrt{G} Y / 2} \ e^{-3u},
\]

where \(C\) is a constant to be determined by fitting to data. Along with Eq. \(2.26\), these define for us the effective \((u\)-dependent\) VAMP potential

\[
V_{\text{eff}}(Y, u) \equiv V_Y + \rho_X = A \ e^{\sqrt{G} Y} + C \ e^{\sqrt{G} Y / 2} \ e^{-3u},
\]

where a \(A\) is just a constant given in terms of model parameters through Eqs. \(2.22\) and \(2.24\).

Hereafter we adopt natural units, \(M_{\text{Pl}} = 1\). Denoting by a prime derivatives with respect to \(u\), the equation of motion for \(Y\) becomes

\[
\frac{Y''}{1 - Y'^2 / 6} + 3 Y' + \frac{\partial_u \rho}{\rho} Y' / 2 + 3 \partial_Y V_{\text{eff}} = 0,
\]

where \(\rho = V_{\text{eff}} + \rho_{\text{rad}} + \rho_{\text{mat}}\). Quantities of importance are the dark energy density

\[
\rho_Y = \frac{1}{2} H^2 Y'^2 + V_Y,
\]

generally expressed in units of the critical density \((\Omega \equiv \rho / \rho_c)\)

\[
\Omega_Y = \frac{\rho_Y}{3H^2},
\]

and the Hubble parameter

\[
H^2 = \frac{\rho}{3 - Y'^2 / 2}.
\]

The equation of state is

\[
w_Y = \left[ \frac{H^2}{2} \frac{Y'^2}{2} - V_Y \right] \left[ \frac{H^2}{2} \frac{Y'^2}{2} + V_Y \right]^{-1}.
\]

We pause to note that the exponential potential \(V_Y \sim e^{\lambda Y / M_{\text{Pl}}}\), with \(\lambda = \sqrt{2}\). Asymptotically, this represents the crossover situation with \(w_Y = -1/3\) \(^{[23]}\), implying expansion at constant velocity. Nevertheless, we will find that there is a brief period encompassing the recent past \((z \leq 6)\) where there has been significant acceleration.

Returning now to the quantitative analysis, we take \(\rho_{\text{mat}} = B e^{-3u}\) and \(\rho_{\text{rad}} = 10^{-4} \rho_{\text{mat}} e^{-u} f(u)\) where \(B\) is a constant and \(f(u)\) parameterizes the \(u\)-dependent number of radiation degrees of freedom.\(^1\) In order to interpolate the various thresholds appearing prior to

\(^1\)This assumption will be justified \textit{a posteriori} when we find that \(\rho_X \ll \rho_{\text{mat}}\).
recombination (among others, QCD and electroweak), we adopt a convenient phenomeno-
logical form $f(u) = \exp(-u/15)$ \[56\]. We note at this point that solutions of Eq. (2.32) are
independent by an overall normalization for the energy density. This is also true for the
dimensionless quantities of interest $\Omega_{Y}$ and $w_{Y}$.

With these forms for the energy densities, Eq. (2.32) can be integrated for various choices
of $A$, $B$, and $C$, and initial conditions at $u = -30$. We take as initial condition $Y(-30) = 0$.
Because of the slow variation of $Y$ over the range of $u$, changes in $Y(-30)$ are equivalent
to altering the quantities $A$ and $C$ \[57\]. In accordance to equipartition arguments \[57, 58\]
we take $Y'(-30) = 0.08$. Because the $Y$ evolution equation depends only on energy density
ratios, and hence only on the ratios $A : B : C$ of the previously introduced constants, we
may, for the purposes of integration and without loss of generality, arbitrarily fix $B$ and
then scan the $A$ and $C$ parameter space for applicable solutions. In Fig. 1 we show a sample
qualitative fit to the data. It has the property of allowing the maximum value of $X$-CDM
(a) about 7% of the total dark matter component) before the fits deviate unacceptably from
data.

It is worth pausing at this juncture to examine the consequences of this model for
variation in the fine structure constant and long range forces. Specifically, excitations of
the electromagnetic field on $N_1$ will, through the presence of the dilaton factor in Eq. (2.2),
seemingly induce variation in the electromagnetic fine structure constant $\alpha_{\text{em}} = e^2/4\pi$, as
well as a violation of the equivalence principle through a long range coupling of the dilaton
to the electromagnetic component of the stress tensor. We now show that these effects are
extremely negligible in the present model. First, it is easily seen using Eqs. (2.2) and (2.3)
together with Eqs. (2.8)-(2.15), that the electromagnetic piece of the lagrangian as viewed
from $N_1$ is

$$L_{\text{em}} = -\frac{2\pi}{4}e^{-\sqrt{G_4}X} \tilde{f}_{\mu\nu}^2,$$ (2.37)

where $\tilde{f}_{\mu\nu}$ denotes a quantum fluctuation of the electromagnetic $U(1)$ field. (Fluctuations
of the $U(1)$ background field are studied in the next section.) At the equilibrium value
$X = X_0$, the exponential factor is

$$e^{-\sqrt{G_4}X_0} = \frac{M_{\text{Pl}}^2}{4\pi b^2},$$ (2.38)

so that we can identify the electromagnetic coupling $(1/e^2) \simeq M_{\text{Pl}}^2/b^2$. This shows that
$b \sim M_{\text{Pl}}$. We can then expand about the equilibrium point, and obtain an additional factor
of $(X - X_0)/M_{\text{Pl}}$. This will do two things \[60\]: (a) At the classical level, it will induce a
variation of the electromagnetic coupling as $X$ varies, with $\Delta \alpha_{\text{em}}/\alpha_{\text{em}} \simeq (X - X_0)/M_{\text{Pl}}$; and
Figure 1: The upper panel shows the evolution of $Y$ as a function of $u$. Today corresponds to $z = 0$ and for primordial nucleosynthesis $z \approx 10^{10}$. We set the initial conditions $Y(-30) = 0$ and $Y'(-30) = 0.08$; we take $A : B : C = 11 : 0.3 : 0.1$. The second panel shows the evolution of $\Omega_Y$ (solid line), $\Omega_{\text{mat}}$ (dot-dashed line), and $\Omega_{\text{rad}}$ (dashed line) superposed over experimental best fits from SDSS and WMAP observations [49, 11]. The curves are not actual fits to the experimental data but are based on the particular choice of the $Y$ evolution shown in the upper panel, which provides eyeball agreement with existing astrophysical observations. The lower panel shows the evolution of the equation of state $w_Y$ superposed over the best fits to WMAP + SDSS data sets and WMAP + SNGold [11]. The solution of the field equations is consistent with the requirement from primordial nucleosynthesis, $\Omega_Y < 0.045$ (90%CL) [12, 48]; it also shows the established radiation and matter dominated epochs, and at the end shows an accelerated dS era [59].

(b) at the quantum level, exchange of $X$ quanta will induce a new force through coupling to the electromagnetic component of matter.

Item (b) is dangerous if the mass of the exchanged quanta are small, so that the force is long range. This is not the case in the present model: from Eq. (2.22) the $X$ quanta have
mass of $O(\mathcal{M}_X \Lambda_{Pl}) \sim M_{Pl}/(r_c b)$, so that if $r_c$ is much less than $O(\text{cm})$, the forces will play no role in the laboratory or cosmologically.

As far as the variation of $\alpha_{em}$ is concerned, we find that $\rho_X/\rho_{\text{mat}} = (C/B)e^Y/e^\sqrt{2}$, so that

$$\rho_X \simeq 3 \times 10^{-120} e^{-3u} M_{Pl}^4 e^{Y/e^\sqrt{2}} = \frac{1}{4} \mathcal{M}_X (X - X_0)^2 e^{Y/e^\sqrt{2}} M_{Pl}^2 \ .$$

This then gives,

$$\sqrt{\langle (X - X_0)^2 \rangle} \equiv \Delta X_{\text{rms}} \approx 10^{-60} e^{-3u/2} M_{Pl} e^{Y/(2 e^\sqrt{2})} / \mathcal{M}_X \ .$$

During the radiation era, $Y \simeq \text{const} \simeq 0$ (see Fig. 1), so that during nucleosynthesis ($u \simeq -23$) $\Delta X_{\text{rms}}/M_{Pl} \simeq 10^{-45}/\mathcal{M}_X$, certainly no threat. It is interesting that such a small value can be understood as a result of inflation: from the equation of motion for the $X$ field, it is simple to see that during a dS era with Hubble constant $H$, the amplitude $\Delta X_{\text{rms}}$ is damped as $e^{-3Ht/2}$. For 50 e-foldings, this represents a damping of $10^{32}$. In order to make the numbers match (assuming a pre-inflation value $\Delta X_{\text{rms}}/M_{Pl} \sim 1$), an additional damping of $\sim 10^{13}$ is required from reheat temperature to primordial nucleosynthesis. With the $e^{-3u/2}$ behavior, this implies a low reheat temperature, about $10^6$ GeV. Otherwise, one may just assume an additional fine-tuning of the initial condition on $X$.

As mentioned previously, the solutions of Eq. (2.32), as well as the quantities we are fitting to ($\Omega_Y$ and $w_Y$), depend only on the ratios of the energy densities. From the eyeball fit in Fig. 1 we have, up to a common constant, $\rho_{\text{ordinary matter}} \equiv \rho_{\text{mat}} \propto 0.3 \ e^{-3u}$ and $V_Y \propto 11 e^{\sqrt{2} Y}$. We can deduce from these relations that

$$\frac{V_Y(\text{now})}{\rho_{\text{mat}}(\text{now})} = \frac{11}{0.3} e^{\sqrt{2} Y(\text{now})} \simeq 36 \ e^{\sqrt{2} Y(\text{now})} \ .$$

Besides, we know that $\rho_{\text{mat}}(\text{now}) \simeq 0.3 \rho_c(\text{now}) \simeq 10^{-120} M_{Pl}^4$. Now, Eqs. (2.22) and (2.24) lead to

$$V_Y(\text{now}) = e^{\sqrt{2} Y(\text{now})} \frac{M_{Pl}^4}{8\pi r_c^2 b^2} (b^2 \xi - 1)$$

so that from Eqs. (2.41) and (2.42) we obtain

$$\frac{1}{8\pi r_c^2 b^2} (b^2 \xi - 1) \simeq 10^{-119} \ .$$

It is apparent that this condition cannot be naturally accomplished by choosing large values of $r_c$ and/or $b$. There remains the possibility that SUSY breaking [45] or non-perturbative effects lead to an exponentially small deviation of $b^2 \xi$ from unity, such that $b^2 \xi = 1 + \ldots$
\[O(10^{-119})\]^2 Since a deviation of \(b^2 \xi\) from unity involves a breaking of supersymmetry, a small value for this dimensionless parameter, perhaps \((1 \text{ TeV}/M_{\text{Pl}})^2 \sim 10^{-31}\), can be expected on the basis of ’t Hooft naturalness. It is the extent of the smallness, of course, which remains to be explained.

### 2.2 Fluctuations in the Background Configuration

In this section we study the quantum fluctuations of the \(U(1)\) field associated to the background configuration. We start by considering fluctuations of the background field \(A_M^0\) in the 4 dimensional space, i.e,

\[A_M \rightarrow A_M^0 + \epsilon a_M,\]  

where \(A_M^0 = 0\) if \(M \neq \varphi\) and \(a_M = 0\) if \(M = \vartheta, \varphi\). The fluctuations on \(A_M^0\) lead to

\[F_{MN} \rightarrow F_{MN}^0 + \epsilon f_{MN}.\]  

Then,

\[F_{MN}F^{MN} = g^{ML} g^{NP}[F_{MN}^0 F_{LP}^0 + \epsilon F_{MN}^0 f_{LP} + \epsilon^2 f_{MN} f_{LP}].\]  

The second term vanishes and the first and third terms are nonzero because \(F_{MN}^0 \neq 0\) in the compact space and \(f_{MN} \neq 0\) in the 4 dimensional space. If the Kalb-Ramond potential \(B_{NM} = 0\), then the 3-form field strength can be written as

\[G_{MNP} = \kappa A_{[M} F_{NP]} = \frac{\kappa}{3!} [A_M F_{NP} + A_P F_{MN} - A_N F_{MP}].\]  

Now we introduce notation of differential forms, in which the usual Maxwell field and field strength read

\[A_1 = A_M dx^M\] and \[F_2 = F_{MN} dx^M \wedge dx^N;\]  

\(^2\) Before proceeding, we remind the reader that the requirements for preserving a fraction of SUSY in spherical compactifications to four dimensions imply \(b^2 \xi = 1\), corresponding to the winding number \(n = \pm 1\) for the monopole configuration. In terms of the Bohm-Aharonov argument on phases, this is consistent with the usual requirement of quantization of the monopole. The SUSY breaking has associated a non-quantized flux of the field supporting the two sphere. In other words, if we perform a Bohm-Aharonov-like interference experiment, some phase change will be detected by a \(U(1)\) charged particle that circulates around the associated Dirac string. The quantization of fluxes implied the unobservability of such a phase, and so in our cosmological set-up, the parallel transport of a fermion will be slightly path dependent. One possibility is that the non-compact \(\rho\) coordinate (in the uplift to ten dimensions, see Appendix B) is the direction in which the Dirac string exists. Then the cutoff necessary on the physics at large \(\rho\) will introduce a slight (time-dependent) perturbation on the flux quantization condition.
respectively. (Note that $dx^M \wedge dx^N$ is antisymmetrized by definition.) With this in mind, the 3-form reads
\[ G_3 = \kappa A_1 \wedge F_2 = \kappa A_M F_{NP} \ dx^M \wedge dx^N \wedge dx^P. \] (2.49)

Substituting Eqs. (2.44) and (2.45) into Eq. (2.49) we obtain
\[ G_3 = \kappa \left[ (A_0^M + \epsilon a_M) (F_0^N + \epsilon f_{NP}) \ dx^M \wedge dx^N \wedge dx^P \right]. \] (2.50)

The background fields read
\[ A_0^1 = b \cos \vartheta \ d\varphi, \quad F_0^2 = -b \sin \vartheta \ d\vartheta \wedge d\varphi, \] (2.51)
and the fluctuations on the probe brane become
\[ a_1 = a_\mu dx^\mu, \quad f_2 = f dx^\mu \wedge dx^\nu, \quad \text{with} \quad f = \partial_\mu a_\nu - \partial_\nu a_\mu. \] (2.52)

All in all,
\[ \frac{G_3}{\kappa} = A_0^0 F_0^0 \ d\varphi \wedge d\vartheta \wedge d\varphi + \epsilon A_0^0 f_{\mu \nu} \ d\varphi \wedge dx^\mu \wedge dx^\nu + \epsilon F_0^0 a_\mu \ d\vartheta \wedge d\varphi \wedge dx^\mu + \epsilon^2 a_\mu f_{\zeta \nu} dx^\mu \wedge dx^\zeta \wedge dx^\nu. \] (2.53)

Using Eq. (2.51) and the antisymmetry of the wedge product, Eq. (2.53) can be re-written as
\[ \frac{G_3}{\kappa} = \epsilon \left[ b \cos \vartheta f_{\mu \nu} d\varphi \wedge dx^\mu \wedge dx^\nu - ba_\mu \sin \vartheta d\vartheta \wedge d\varphi \wedge dx^\mu + \epsilon a_\mu f_{\zeta \nu} dx^\mu \wedge dx^\zeta \wedge dx^\nu \right]. \] (2.54)

From the metric
\[ ds^2 = e^{2\alpha} dx_4^2 + e^{2\beta} (d\vartheta^2 + \sin \vartheta^2 d\varphi^2) \] (2.55)
we can write the vielbeins
\[ e^a = e^\alpha dx^a, \quad e^\vartheta = e^\beta d\vartheta, \quad e^\varphi = e^\beta \sin \vartheta d\varphi, \]
\[ dx^a = -e^{-\alpha} e^a, \quad d\vartheta = e^{-\beta} e^\vartheta, \quad d\varphi = \frac{e^{-\beta}}{\sin \vartheta} e^\varphi \] (2.56)
where $\beta \equiv f + \ln r_c$. (Lower latin indeces from the beginning of the alphabet indicate coordinates associated to the four dimensional Minkowski spacetime with metric $\eta_{ab}$.) Substituting into Eq. (2.53) we obtain
\[ \frac{G_3}{\kappa} = \epsilon \left[ b \cos \frac{\vartheta}{\sin \vartheta} e^{-2\alpha - 2\beta} f_{ab} e^a \wedge e^b + e^{-\alpha - 2\beta} a_\mu e^\mu \wedge e^a + \epsilon e^{-3\alpha} a_\mu f_{ab} e^a \wedge e^c \wedge e^b \right], \] (2.57)
where $f_{ab} = \partial_a a_b - \partial_b a_a$. Because the three terms are orthogonal to each other, a straightforward calculation leads to
\[ G_3^2 = \kappa^2 e^2 (b^2 \cot^2 \vartheta e^{-4\alpha - 2\beta} f_{ab}^2 + b^2 e^{-2\alpha - 2\beta} a_a^2) + \mathcal{O}(e^4). \] (2.58)
Then, the 5th term in Eq. (2.2) can be written as

\[
S_G^3 = -\frac{1}{2G_6} \int d^4x \frac{G_6}{6} e^{4\alpha + 2\beta} \sqrt{\eta_4} e^{-2\phi} \int d\vartheta d\varphi \sin \vartheta \left[ (\kappa^2 e^2 b^2 \cot^2 \vartheta e^{-4\alpha - 2\beta}) f_{ab}^2 + \left( \kappa^2 e^2 b^2 e^{-2\alpha - 4\beta} \right) a^2_\alpha \right],
\]

whereas the contribution from the 4th term in Eq. (2.2) can be computed from Eq. (2.46) yielding

\[
S_F^2 = -\frac{1}{2G_6} \int d^4x \sqrt{\eta_4} e^{2\beta - \phi} e^2 f_{ab}^2 = -\int d^4x \sqrt{\eta_4} e^{2\beta - \phi} r^2 e^2 f_{ab}^2.
\]

Thus,

\[
S_G^3 + S_F^2 = -\int d^4x \left[ \frac{1}{4g^2} f_{ab}^2 + m^2 \frac{a^2_\alpha}{2} \right],
\]

where the four dimensional effective coupling and the effective mass are of the form

\[
\frac{1}{g^2} = 4 \epsilon^2 \sqrt{\eta_4} \left[ \pi e^{2\beta - \phi} r^2_\epsilon + \frac{1}{12} \kappa^2 b^2 e^{-2\phi} \int d\vartheta d\varphi \sin \vartheta \cot^2 \vartheta \right] \to \infty
\]

and

\[
m^2 = \frac{2}{3} \pi \kappa^2 b^2 e^{2\alpha - 2\beta - 2\phi}.
\]

For the moment we let \( \int d\vartheta d\varphi \sin \vartheta \cot^2 \vartheta = N \), where eventually we set \( N \to \infty \). Now to make quantum particle identification and coupling, we carry out the transformation \( a_\alpha \to g \hat{a}_\alpha \). This implies that the second term in the right hand side of Eq. (2.61) vanishes, yielding

\[
f_{ab} = \partial_a (g \hat{a}_b) - \partial_b (g \hat{a}_a) = \partial_a g \hat{a}_b - \partial_b g \hat{a}_a + g \partial_a \hat{a}_b - g \partial_b \hat{a}_a = g \hat{f}_{ab} + \hat{a} \wedge dg
\]

and consequently to leading order in \( N \)

\[
\frac{1}{g^2} f_{ab}^2 = \frac{1}{g^2} [g^2 \hat{f}_{ab}^2 + (\hat{a} \wedge dg)^2 + 2 g \hat{a}_b \hat{f}_{ab} \partial_a g].
\]

If the coupling depends only on the time variable,

\[
\frac{1}{g^2} f_{ab}^2 \to \hat{f}_{ab}^2 + \left( \frac{\dot{g}}{g} \right)^2 \hat{a}_\alpha^2 + 2 \frac{\dot{g}}{g} \dot{\hat{a}}_i \hat{f}_{ti}^i.
\]
where \( \dot{g} = \partial_t g \) and lower latin indices from the middle of the alphabet refer to the brane space-like dimensions. If we choose a time-like gauge in which \( a_t = 0 \), then the term 

\[(\dot{g}/g) \dot{a}_i \dot{a}_i\]

can be written as \((1/2)(\dot{g}/g)(d/dt)(\dot{a}_i)^2\), which after an integration by parts gives 

\[-(1/2)[(d/dt)(\dot{g}/g)]\dot{a}_i^2;\]

with \( g \sim e^{-\phi} \), the factor in square brackets becomes \(-\ddot{\phi}\). Since \( \phi = \sqrt{G_4}(X + Y) \), the rapidly varying \( \dddot{X} \) will average to zero, and one is left just with the very small \( \dddot{Y} \), which is of order Hubble square. For the term \((\dot{g}/g)^2(a_i)^2\), the term \((\dddot{X})^2\) also averages to order Hubble square, implying that the induced mass term is of horizon size. These “paraphotons” carry new relativistic degrees of freedom, which could in turn modify the Hubble expansion rate during Big Bang nucleosynthesis (BBN). Note, however, that these extremely light gauge bosons are thought to be created through inflaton decay and their interactions are only relevant at Planck-type energies. Since the quantum gravity era, all the paraphotons have been redshifting down without being subject to reheating, and consequently at BBN they only count for a fraction of an extra neutrino species in agreement with observations.
Chapter 3

Photon Signals of Low Mass Strings at the LHC

The CERN’s LHC is the greatest basic science endeavor in history. Spectacular physics results are expected to follow in short order once it turns on this year. The LHC will push nucleon-nucleon center-of-mass energies up to $\sqrt{s} = 5.5$ TeV for Pb-Pb collisions, and $\sqrt{s} = 14$ TeV for $pp$ collisions. The ALICE detector will observe the very messy debris of heavy ion collisions, whereas the ATLAS and CMS detectors will observe the highest-energy particle collisions produced by the accelerator. The LHC will probe deeply into the sub-fermi distances, committing to careful searches for new particles and interactions at the TeV scale.

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 SM physics. Important advances were fueled by the realization of the vital role played by D-branes [61] in connecting string theory to phenomenology [10]. This has permitted the formulation of string theories with compositeness setting in at TeV scales [8] 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 $\sqrt{s} \sim$ 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 chapter we exploit both these properties in order to obtain a “new physics” signal at the LHC which, if traced to low scale string theory, could with 100 fb$^{-1}$ of data probe deviations from the SM physics at a $5\sigma$ significance for $M_s$ as large as 2.3 TeV.
3.1 Perturbative D-brane Models

The concept of D-branes was introduced in the late 80’s [63]. They are described as a geometric locus where strings can end. In a quantum gravity theory, like string theory, any defect or extended object in spacetime can bend and it will consequently have excitations. For D-branes, the excitations are the open string attached to them. Furthermore, one can have various D-branes on top of one another. In these situations one needs to consider open strings with Chan-Paton indices [64] on them, and thus one has continuous gauge symmetries associated to the ends of the string. The allowed gauge groups in these D-brane constructions are those that can have a large $N$ limit: $U(N)$, $SO(N)$, $Sp(N)$. Besides, each end of the string carries a fundamental charge with respect to the stack of branes on which it ends. Hence, any open string will carry the quantum numbers associated to some type of bifundamental representation. These novel constructions provide a framework for particle physics on a brane, yielding a possible realization of the SM within string theory.

To describe the field theory degrees of freedom it is convenient to introduce a graphic notation, generally referred to as quivers or moose diagrams. The gauge degrees of freedom (brane stacks) are then described by nodes in a graph. The open string particles are given by edges connecting two vertices and arrows that dictate if the corresponding end of the string is fundamental or anti-fundamental. One should also label the edges according to the other quantum numbers that the particles carry. In the perturbative regime, where the low energy dynamics is given in terms of open strings alone (all other non-perturbative states are heavy), the interactions are generated by disc diagrams (a relevant example is pictured in Fig. 2). These are single traces of fields. If a vertex has $n + 2$ particles attached to it, it will appear with a coupling constant dependence of $g^n$, where $g$ is the open string coupling constant.

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 ($C_\mu$, in the notation of [65]), most simply the manifestation of a gauged baryon number symmetry. The $U(1)_Y$ boson $Y_\mu$, which gauges the usual electroweak hypercharge symmetry, is a linear combination of $C_\mu$, the $U(1)$ boson $B_\mu$ terminating on a separate $U(1)$ brane, and perhaps a third additional $U(1)$ (say $W_\mu$) sharing a $U(2)$ brane to which are also a terminus for the $SU(2)_L$ electroweak gauge bosons $W_\mu^a$. Thus, critically for our purposes, the photon $A_\mu$, which is a linear combination of $Y_\mu$ and $W_\mu^3$, 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. [65, 66, 67]. The vector boson $Z'_\mu$, orthogonal to the hypercharge, must grow a mass $M_{Z'}$ in order to avoid long range forces between baryons other than gravity and Coulomb forces. The anomalous mass growth allows the survival of global baryon number conservation, preventing fast proton decay [68].

The processes we consider (at the parton level) are $gg \to g\gamma$ and $gg \to \gamma\gamma$, where $g$ is an $SU(3)$ gluon and $\gamma$ is the photon. As explicitly calculated below, these will occur at string disk (tree) level, and will be manifest at the LHC as a non-SM contribution to $pp \to \gamma + \text{jet}$ and $pp \to \gamma\gamma$. 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 these signals originates in the parton tree level processes $gq \to \gamma q$, $g\bar{q} \to \gamma \bar{q}$, $q\bar{q} \to \gamma g$, and $q\bar{q} \to \gamma\gamma$. Of course, the SM processes will also receive stringy corrections which should be added to the pure bosonic contribution as part of the signal. We postpone their evaluation until Chapter 5. 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 [72].

---

$^1$Some qualitative and quantitative considerations of these processes have been discussed in [69, 70, 71].
### 3.2 The String Amplitude

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 non-vanishing contributions can then be generated from a single, maximally helicity violating (MHV), amplitude – the so-called partial MHV amplitude [73]. 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. (All momenta are incoming.) Then the partial amplitude for such an MHV configuration is given by [74, 75]

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

where \( g \) is the \( U(N) \) coupling constant, \( \langle ij \rangle \) are the standard spinor products written in the notation of Refs. [76, 77], and the Veneziano formfactor [78, 79],

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

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

\[
s = \frac{2k_1 k_2}{M^2}, \quad t = \frac{2k_1 k_3}{M^2}, \quad u = \frac{2k_1 k_4}{M^2} : \quad s + t + u = 0.
\]

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

\[
V(s, t, u) \approx 1 - \frac{\pi^2}{6} s u - \zeta(3) s t u + \ldots
\]

We first consider the amplitude involving three \( SU(N) \) gluons \( g_1, g_2, g_3 \) and one \( U(1) \) gauge boson \( \gamma_4 \) associated to the same \( U(N) \) quiver:

\[
T^{a_1} = T^a, \quad T^{a_2} = T^b, \quad T^{a_3} = T^c, \quad T^{a_4} = Q_c I,
\]

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

\[
\text{Tr}(T^a T^b) = \frac{1}{2} \delta^{ab}.
\]

Then the color factor

\[
\text{Tr}(T^{a_1} T^{a_2} T^{a_3} T^{a_4}) = Q_c (d^{abc} + \frac{i}{4} f^{abc}),
\]

where

\[
d^{abc} + \frac{i}{4} f^{abc}.
\]
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 \[74, 75\] by summing the partial amplitudes \[(3.1)\] with the indices permuted in the following way:

$$
\mathcal{M}(g_1^-, g_2^-, g_3^+, \gamma_4^+) = 4 g^2 (12)^4 \sum \sigma \frac{\text{Tr} \left( T^{a_1} T^{a_2} T^{a_3} T^{a_4} \right) V(k_{1\sigma}, k_{2\sigma}, k_{3\sigma}, k_{4\sigma})}{\langle 12 \rangle \langle 23 \rangle \langle 34 \rangle \langle 41 \rangle} + 1^{16}
\tag{3.8}
$$
where the sum runs over all 6 permutations $\sigma$ of $\{1, 2, 3\}$ and $i_\sigma \equiv \sigma(i)$. Note that in the effective field theory of gauge bosons there are no Yang-Mills interactions that could generate this scattering process at the tree level. Indeed, $V = 1$ at the leading order of Eq.\[(3.4)\] and the amplitude vanishes due to the following identity:

$$
\begin{align*}
1^{(12)(23)(34)(41)} + 1^{(23)(31)(14)(42)} + 1^{(31)(12)(24)(43)} & = 0.
\tag{3.9}
\end{align*}
$$

Similarly, the antisymmetric part of the color factor \[(3.7)\] cancels out in the full amplitude \[(3.8)\]. As a result, one obtains:

$$
\mathcal{M}(g_1^-, g_2^-, g_3^+, \gamma_4^+) = 8 Q_c d^{abc} d^{abc} g^2 (12)^4 \left( \frac{\mu(s, t, u)}{\langle 12 \rangle \langle 23 \rangle \langle 34 \rangle \langle 41 \rangle} + \frac{\mu(s, u, t)}{\langle 12 \rangle \langle 24 \rangle \langle 13 \rangle \langle 34 \rangle} \right),
\tag{3.10}
$$
where

$$
\mu(s, t, u) = \Gamma(1 - u) \left( \frac{\Gamma(1 - s)}{\Gamma(1 + t)} - \frac{\Gamma(1 - t)}{\Gamma(1 + s)} \right).
\tag{3.11}
$$

All non-vanishing amplitudes can be obtained in a similar way. In particular,

$$
\mathcal{M}(g_1^-, g_2^+, g_3^-, \gamma_4^+) = 8 Q_c d^{abc} g^2 (13)^4 \left( \frac{\mu(t, s, u)}{\langle 13 \rangle \langle 24 \rangle \langle 14 \rangle \langle 23 \rangle} + \frac{\mu(t, u, s)}{\langle 13 \rangle \langle 24 \rangle \langle 12 \rangle \langle 34 \rangle} \right),
\tag{3.12}
$$
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 \[(3.8)\] is:

$$
|\mathcal{M}(g_1^-, g_2^-, g_3^+, \gamma_4^+)|^2 = 64 Q_c^2 d^{abc} d^{abc} g^4 \left| \frac{s \mu(s, t, u)}{u} + \frac{s \mu(s, u, t)}{t} \right|^2.
\tag{3.13}
$$

Taking into account all $4(N^2 - 1)^2$ possible initial polarization/color configurations and the formula \[80\]

$$
\sum_{a,b,c} d^{abc} d^{abc} = \frac{(N^2 - 1)(N^2 - 4)}{16N},
\tag{3.14}
$$
we obtain the average squared amplitude \[81\]

\[ |\mathcal{M}(gg \rightarrow g\gamma)|^2 = g^4 Q_c^2 C(N) \left\{ \left| \frac{s\mu(s, t, u)}{u} + \frac{s\mu(s, u, t)}{t} \right|^2 + (s \leftrightarrow t) + (s \leftrightarrow u) \right\}, \tag{3.15} \]

where

\[ C(N) = \frac{2(N^2 - 4)}{N(N^2 - 1)}. \tag{3.16} \]

Next, we consider the amplitude involving two SU(N) gluons \(g_1, g_2\) and two U(1) gauge bosons \(\gamma_3, \gamma_4\) associated to the same U(N) quiver:

\[ T^{a_1} = T^a, \quad T^{a_2} = T^b, \quad T^{a_3} = Q_c I, \quad T^{a_4} = Q_c I. \tag{3.17} \]

This amplitude can be obtained from \(gg \rightarrow g\gamma\) by replacing \(d^{abc}\) with \(\frac{1}{2}Q_c \delta^{ab}\). Hence at the level of squared amplitudes, summed over final polarizations and colors and averaged over initial polarizations and colors \[82\]

\[ |\mathcal{M}(gg \rightarrow \gamma\gamma)|^2 = \frac{4N Q_c^2}{N^2 - 4} |\mathcal{M}(gg \rightarrow g\gamma)|^2. \tag{3.18} \]

The two most interesting energy regimes of \(gg \rightarrow g\gamma\) 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. (3.15) becomes

\[ |\mathcal{M}(gg \rightarrow g\gamma)|^2 \approx g^4 Q_c^2 C(N) \frac{\pi^4}{4} (s^4 + t^4 + u^4) \quad (s, t, u \ll 1). \tag{3.19} \]

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)

\[ |\mathcal{M}(gg \rightarrow g\gamma)|^2 \approx 4g^4 Q_c^2 C(N) \frac{M_s^8 + t^4 + u^4}{M_s^4((s - M_s^2)^2 + (\Gamma M_s)^2)} \quad (s \approx M_s^2), \tag{3.20} \]

with the singularity (smeared with a width \(\Gamma\)) reflecting the presence of a massive string mode propagating in the \(s\) channel. (Further details are given in Appendix C.) It should be noted that because the \(g\gamma\) final state projects onto pure color octet, only the \(SU(3)\) adjoint string excitations \((G^*)\) contribute to Eq. (3.20) \[83\]. On the other hand, for \(gg \rightarrow \gamma\gamma\) only the color singlet excitation \((C_0^*)\) is present in the intermediate state.

An important modification needs to be introduced into Eq. (3.20), because it contains additively contributions from both angular momenta \(J = 0\) and \(J = 2\), corresponding to incoming helicities \((\pm \pm)\) and \((\pm \mp)\), respectively. In general these contributions would have different widths, for the \(G^*\) excitation: \(\Gamma^{J=0} = 75 (M_s/\text{TeV})\) GeV and
$\Gamma^{J=2} = 45 \text{(M}_s/\text{TeV}) \text{ GeV}$. These widths are premised on the assumption that corrections of order $(M_Z'/M_s)^2$ are negligible, both in obtaining matrix elements and in calculating phase space.

In what follows we will take $N = 3$ and set $g$ equal to the QCD coupling constant, $\alpha_s = (g^2/4\pi) \sim 0.1$. Before proceeding with numerical calculation, we need to make precise the value of $Q_c$. If we were considering the process $gg \rightarrow C^0 g$, where $C^0$ is the $U(1)$ gauge field tied to the $U(3)$ brane, then $Q_c = \sqrt{1/6}$ due to the normalization condition (3.6). However, for $gg \rightarrow \gamma g$ there are two additional projections: from $C^0$ to the hypercharge boson $Y$, giving a mixing factor $\kappa$; and from $Y$ onto a photon, providing an additional factor $\cos \theta_W$ ($\theta_W =$ Weinberg angle). The $C^0 - Y$ mixing coefficient is model dependent: in the minimal model $[65]$ 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 RG running of the couplings up to 2.5 TeV. It should be noted that in models $[66, 67]$ 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 $\kappa$ which can differ considerably from 0.12. In what follows, we take as a fiducial value $\kappa^2 = 0.02$. Thus, if (3.20) is to describe $gg \rightarrow \gamma g$, we modify our definition of $Q_c$ given in Eq. (3.5) to accommodate the additional mixings, and obtain

$$Q_c^2 = \frac{1}{6} \kappa^2 \cos^2 \theta_W \simeq 2.55 \times 10^{-2} \left(\kappa^2/0.02\right).$$

In the remainder of this chapter, we explore potential searches for Regge excitations of fundamental strings at the LHC.

### 3.3 High-$k_\perp$ Isolated Photons

In order to assess the possibility of discovery of signal above QCD background, we adopt the kind of signal introduced in $[84]$ to study detection of TeV-scale black holes at the LHC, namely a high-$k_\perp$ isolated $\gamma$ or $Z$. Thus, armed with parton distribution functions (CTEQ6D) $[85, 86]$, in what follows we calculate integrated cross sections $\sigma(pp \rightarrow \gamma + \text{jet})|_{k_\perp(\gamma)>k_{\perp\min}}$ for both the background QCD processes and for $gg \rightarrow \gamma g$, for an array of values for the string scale $M_s$.

#### 3.3.1 QCD background

The SM background for processes with a single photon in the final state originates in the parton tree level processes $gq \rightarrow \gamma q$, $g\bar{q} \rightarrow \gamma \bar{q}$ and $q\bar{q} \rightarrow \gamma g$,

$$2E' \frac{d\sigma}{d^3k'} \bigg|_{pp \rightarrow \gamma X} = \sum_{ijk} \int dx_a dx_b f_i(x_a, Q) f_j(x_b, Q) 2E' \frac{d\hat{\sigma}}{d^3k'} \bigg|_{ij \rightarrow \gamma k},$$

(3.22)
where \( x_a \) and \( x_b \) are the longitudinal fractions of momenta of the parent hadrons carried by the partons which collide, \( k' \) (\( E' \)) is the photon momentum (energy), \( d\hat{\sigma}/d^3k'|_{ij \to \gamma k} \) is the cross section for scattering of partons of type \( i \) and \( j \) according to elementary QCD diagrams, \( f_i(x_a, Q) \) and \( f_j(x_b, Q) \) are parton distribution functions, \( Q \) is the momentum transfer, and the sum is over the parton species: \( g, q = u, d, s, c, b \). In what follows, we focus on \( gq \to \gamma q \), which results in the dominant contribution to the total cross section. Corrections from the other two processes can be computed in a similar fashion. The hard parton-level cross section reads,

\[
2E' \frac{d\hat{\sigma}}{d^3k'}(g(k)q(p) \to \gamma(k')q(p)) = \frac{1}{(2\pi)^2} \frac{1}{2s} \delta[(k + p - k')^2] \frac{1}{4} \sum |M|^2
\]  

(3.23)

where the variables \( k, p, k' \) and \( p' \) in the parentheses are the momenta of the partons. Here, the amplitude for \( gq \to \gamma q \) is given by

\[
\frac{1}{4} \sum |M|^2 = \frac{1}{3} g^2 e^2 e_q^2 \left( \frac{\hat{s}}{\hat{s} + \hat{t}} + \frac{\hat{s} + \hat{u}}{\hat{t}} \right),
\]  

(3.24)

where \( \hat{s} = (k + p)^2 \), \( \hat{t} = (k - k')^2 \) and \( \hat{u} = (k - p')^2 \) are the Mandelstam variables in the parton level, \( g \) and \( e \) are the QCD and electromagnetic coupling constants, and \( e_q \) is the fractional electric charge of species \( q \). For completeness we note that for \( q\bar{q} \to g\gamma \),

\[
\frac{1}{4} \sum |M|^2 = \frac{8}{9} g^2 e^2 e_q^2 \left( \frac{-\hat{t}}{\hat{s} + \hat{t}} - \frac{\hat{s} + \hat{u}}{\hat{t}} \right).
\]  

(3.25)

In this process, we assume that the proton momenta take the explicit forms in the \( pp \) center of mass frame

\[
P_1 = (\sqrt{s}/2, 0, 0, \sqrt{s}/2), \quad P_2 = (\sqrt{s}/2, 0, 0, -\sqrt{s}/2)
\]  

(3.26)

and the final photon has momentum

\[
k_0' = k_\perp \cosh y, \quad k_\parallel = k_\perp \sinh y,
\]  

(3.27)

where \( k_\perp \) is the transverse momentum of the photon and \( y \) is called the longitudinal rapidity. The relation between the momenta of protons \( P_1, P_2 \) and those of the incoming partons \( k, p \) can be written with the longitudinal fractions \( x_a, x_b \):

\[
k = x_a P_1, \quad p = x_b P_2.
\]  

(3.28)

Using Eqs. (3.26), (3.27) and (3.28), we can re-write the argument of the delta function as

\[
(k + p - k')^2 = 2x_b P_2 \cdot (x_a P_1 - k') + \hat{t} = x_b x_a s - 2x_b P_2 \cdot k' + \hat{t} = x_a x_b s - \sqrt{s} x_b k_\perp e^y - \sqrt{s} x_a k_\perp e^{-y}.
\]  

(3.29)
so that
\[
\begin{align*}
\delta[(k+p-k')^2] &= \delta(x_a \cdot x_b \cdot s - \sqrt{s} x_b k_{\perp} e^y - \sqrt{s} x_a k_{\perp} e^{-y}) \\
&= \frac{1}{s [x_a - x_{\perp} e^y]} \delta \left( x_b - \frac{x_a x_{\perp} e^y}{x_a - x_{\perp} e^y} \right), \quad (3.30)
\end{align*}
\]
where \( x_{\perp} = k_{\perp} / \sqrt{s} \). The lower bound \( x_b > 0 \) implies \( x_a > x_{\perp} e^y \). The upper bound \( x_b < 1 \) leads to a stronger constraint
\[
x_a > \frac{x_{\perp} e^y}{1 - x_{\perp} e^{-y}}, \quad (3.31)
\]
which requires \( x_{\perp} e^y < 1 - x_{\perp} e^{-y} \), yielding \( x_{\perp} < (2 \cosh y)^{-1} \). Of course there is another completely symmetric term, in which \( g \) comes from \( P_2 \) and \( q \) comes from \( P_1 \). Putting all this together, the total contribution from \( gg \rightarrow qg \) reads
\[
\begin{align*}
\sigma_{gg \rightarrow qg}^{pp \rightarrow \gamma X} &= 2 \sum_q \int \frac{d^3k'}{2E'} \int dx_a \int dx_b f_g(x_a, Q) f_q(x_b, Q) \frac{1}{(2\pi)^2} \frac{1}{s [x_a - x_{\perp} e^y]} \\
&\times \frac{1}{2s} \delta \left( x_b - \frac{x_a x_{\perp} e^{-y}}{x_a - x_{\perp} e^y} \right) \frac{e^2 g^2 e_q^2}{3} \left( \hat{s} + \hat{t} + \hat{s} \right), \quad (3.32)
\end{align*}
\]
With the change of variables \( z = e^y \) and the relation
\[
\frac{d^3k'}{2E'} = \pi k_{\perp} dk_{\perp} dy = \frac{\pi k_{\perp} dk_{\perp} dz}{z}, \quad \frac{\hat{t}}{\hat{s}} = -\frac{\sqrt{s} k_{\perp} e^{-y}}{x_b s} = -\frac{x_{\perp}}{x_b z}, \quad (3.33)
\]
Eq. (3.32) can be re-written as
\[
\begin{align*}
\sigma_{gg \rightarrow qg}^{pp \rightarrow \gamma X} &= \frac{e^2 g^2}{12\pi s} \int_{x_{\perp,\min}}^{1/2} dx_{\perp} \int_{z_{\min}}^{z_{\max}} dz \int_{x_{a,\min}}^{1} dx_a f_g(x_a, Q) \left[ \sum_q e^2 f_q \left( \frac{x_a x_{\perp} z^{-1}}{x_a - x_{\perp} z}, Q \right) \right] \\
&\times \frac{1}{x_a^2} \left( \frac{x_{\perp} z}{x_a} + \frac{x_a}{x_{\perp} z} \right), \quad (3.34)
\end{align*}
\]
where the integration limits,
\[
z_{\max,\min} = \frac{1}{2} \left[ \frac{1}{x_{\perp}} \pm \frac{1}{x_{\perp}^{-1}} - 4 \right] \quad \text{and} \quad x_{a,\min} = \frac{x_{\perp} z}{1 - x_{\perp} z^{-1}}, \quad (3.35)
\]
are obtained from Eq. (3.31). In Fig. 3 we show the QCD background cross section \( vs \ k_{\perp,\min} \), as obtained through numerical integration of Eq. (3.34). To accommodate the minimal acceptance cuts on final state photons from the CMS [87] and ATLAS [88] proposals, an additional kinematic cut, \( |y| < 2.4 \), has been included in the calculation.
Figure 3: Different contributions to the QCD cross section for $pp \rightarrow \gamma + \text{jet}$ as a function of $k_{\perp,\text{min}}$. It is clearly seen that the $gq \rightarrow \gamma q$ process provides the dominant contribution.
3.3.2 The string signal

For the considerations in this Dissertation, the resonant cross section can be safely approximated by single poles in the Narrow-Width Approximation,

$$\Gamma \sqrt{s_0/\pi} \frac{\pi}{(\hat{s} - s_0)^2 + (\Gamma/\sqrt{s_0})^2} \frac{\pi}{\Gamma/\sqrt{s_0}} \delta(\hat{s} - s_0),$$  

where $s_0 = M_s^2$. The scattering proceeds through $J = 0$ and $J = 2$ angular momentum states, with the $M_s^8$ term in Eq. (3.20) originating from $J = 0$, and the $t^4 + u^4$ piece reflecting $J = 2$ activity. The widths of these two resonances are different, with $\Gamma^{J=0} = (3/4) \alpha_s M_s$, and $\Gamma^{J=2} = (9/20) \alpha_s M_s$. The average string amplitude square in Eq. (3.20) then becomes

$$|\mathcal{M}(gg \rightarrow g\gamma)|^2 \approx 4g^4 Q_c^2 C(N) \frac{\pi}{s_0^{5/2}} \left[ \frac{s_0^4}{\Gamma^{J=0}} + \frac{\hat{t}^4 + (\hat{t} + s_0)^4}{\Gamma^{J=2}} \right] \delta(\hat{s} - s_0)$$

$$= 4g^4 Q_c^2 C(N) \frac{\pi}{\alpha_s s_0^{3/2}} \left\{ \frac{4}{9} s_0^4 + \frac{20}{3} \left[ \hat{t}^4 + (\hat{t} + s_0)^4 \right] \right\} \delta(\hat{s} - s_0).$$  

Thus, the total cross section for single photon production in gluon fusion is given by

$$\sigma_{pp \rightarrow \gamma X}^{gg \rightarrow \gamma g} = \int \frac{d^3k'}{2E'} \int dx_a \int dx_b f_g(x_a, Q) f_g(x_b, Q) \frac{1}{(2\pi)^2} \frac{1}{2 \hat{s}} \delta(x_a x_b - x_b x_\perp z - x_a x_\perp z^{-1})$$

$$\times 4g^4 Q_c^2 C(N) \frac{\pi}{\alpha_s s_0^{3/2}} \left\{ \frac{4}{9} s_0^4 + \frac{20}{9} \left[ \hat{t}^4 + (\hat{t} + s_0)^4 \right] \right\} \delta(\hat{s} - s_0).$$  

We set $Q = M_s$, which is appropriate for the dual picture of string theory. We are aware that for $Q \sim M_s$, the parton distribution functions will receive significant corrections from the rapid increase of degrees of freedom. Fortunately, as noted elsewhere [89], at parton center-of-mass energies corresponding to low-lying string excitations the resonant cross section is largely insensitive to the details of the choice of $Q$. Since the second delta function can be written as

$$\delta(\hat{s} - s_0) = \frac{1}{x_a s} \delta(x_b - \frac{s_0}{x_a s}),$$

integration over $x_b$ leads to

$$\sigma_{pp \rightarrow \gamma X}^{gg \rightarrow \gamma g} = g^4 Q_c^2 C(N) \frac{1}{2 \alpha_s s_0^{3/2}} \int \frac{x \, dx \, dz}{z} \int dx_a f_g(x_a, Q) f_g(\tau_0/x_a, Q) \frac{1}{x_a}$$

$$\times \delta \left( \tau_0 - \frac{\tau_0 x_\perp z}{x_a} - \frac{x_a x_\perp}{z} \right) \left\{ \frac{4}{9} \tau_0 + \frac{20}{9} \left[ (x_a x_\perp z^{-1})^4 + (-x_a x_\perp z^{-1} + \tau_0)^4 \right] \right\},$$  

where $\tau_0 = s_0/s$. In order to proceed the integral over the argument $z$, we write

$$f(x) \equiv \tau_0 - \frac{\tau_0 x_\perp z}{x_a} - \frac{x_a x_\perp}{z}.$$  

(3.41)
Then, the delta function becomes of the form:
\[
\delta(f(z)) = \frac{1}{|f'(z_+)|} \delta(z - z_+) + \frac{1}{|f'(z_-)|} \delta(z - z_-), \tag{3.42}
\]
where \(z_\pm\) are the solutions to \(f(z) = 0\),
\[
z_\pm = \frac{x_a}{2x_\perp} \left(1 \pm \sqrt{1 - \frac{4x_\perp^2}{\tau_0}}\right). \tag{3.43}
\]
Using the identities
\[
\frac{1}{z_\pm |f'(z_\pm)|} = \left|\frac{\tau_0 x_\perp z_\pm}{x_a} - \frac{x_a x_\perp}{z_\pm}\right|^{-1} \tag{3.44}
\]
and
\[
\frac{16}{9} \tau_0^2 (5 x_\perp^4 - 10 x_\perp^2 \tau_0 + 4 \tau_0^2) = \left\{\frac{4}{9} \tau_0^4 + \frac{20}{9} [(x_a x_\perp z_\perp^{-1})^4 + (-x_a x_\perp z_\perp^{-1} + \tau_0)^4]\right\} + \left\{\frac{4}{9} \tau_0^4 + \frac{20}{9} [(x_a x_\perp z_\perp^{-1})^4 + (-x_a x_\perp z_\perp^{-1} + \tau_0)^4]\right\}, \tag{3.45}
\]
the integral over the \(z\) variable yields
\[
\sigma_{gg \rightarrow \gamma g}^{9g \rightarrow 9g} = \frac{8}{9} \frac{g^4 Q_c^2 C(N)}{C_s} \frac{\sqrt{\tau_0}}{\alpha_s \tau_0^2 s} \int_{x_{\perp,\min}}^{\sqrt{\tau_0}/2} \frac{dx_\perp}{\sqrt{1 - 4x_\perp^2/\tau_0}} \left(5 x_\perp^4 - 10 x_\perp^2 \tau_0 + 4 \tau_0^2\right) \times \int_{\tau_0}^{1} \frac{dx_a}{x_a} f_g(x_a, Q) f_g(\tau_0/x_a, Q), \tag{3.46}
\]
where the integration range has been derived from the conditions \(0 < x_b = \tau_0/x_a < 1\) and \(4x_\perp^2 < \tau_0\), which imply \(\tau_0 < x_a < 1\) and \(x_{\perp,\min} < x_\perp < \sqrt{\tau_0}/2\). Finally, integration over \(x_\perp\) leads to
\[
\sigma_{gg \rightarrow \gamma g}^{9g \rightarrow 9g} = \frac{1}{9} \frac{g^4 Q_c^2 C(N)}{C_s} \frac{\sqrt{\tau_0}}{\alpha_s \tau_0^2 s} \sqrt{1 - \frac{4x_{\perp,\min}^2}{\tau_0}} \left(5 \tau_0^2 - 6 \tau_0 x_{\perp,\min}^2 + 2 x_{\perp,\min}^4\right) \times \int_{\tau_0}^{1} \frac{dx_a}{x_a} f_g(x_a, Q) f_g(\tau_0/x_a, Q). \tag{3.47}
\]
In Fig. 4 we show the resonant cross section for \(M_s = 1\) TeV. It is evident that the background is significantly reduced for large \(k_{\perp,\min}\). At very large values of \(k_{\perp,\min}\), however, event rates become problematic. Note that all stringy corrections to the pure bosonic cross section given by Eq. (3.47) have similar factorizations. An illustration of the relative partonic luminosities of the different processes is shown in Fig. 5.
Figure 4: Behavior of the QCD cross section for $pp \to \gamma + \text{jet}$ (dot-dashed line) as a function of $k_{\perp,\text{min}}$. The string cross section overlying the QCD background is also shown as a solid line, for $M_s = 1 \text{ TeV}$. 
Figure 5: Relative contributions of initial state partons ($ij = gg, gq, g\bar{q}$, and $q\bar{q}$) to $\int_{x_0}^{1} \frac{f_i(x_a, Q)}{f_j(\tau_0/x_a, Q)} dx_a/x_a$, with varying string scale.
3.3.3 LHC discovery reach

In this section we explore the LHC discovery potential by computing the signal-to-noise ratio (signal/$\sqrt{\text{SM background}}$). In Fig. 6 we show the string cross section and number of events (before cuts) in a 100 fb$^{-1}$ run at the LHC, for $p_{T,\text{min}} = 300$ GeV, as a function of the string scale $M_s$. For a 300 GeV cut in the transverse momentum, the QCD cross section (shown in Fig. 4) is about $8 \times 10^3$ fb, yielding (for 100 fb$^{-1}$) $\sqrt{\text{SM background}} \approx 895$. A point worth noting at this juncture: 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 [90]. We will leave the exact nature of these cuts for the experimental groups, and present results for a generous range of direct photon reconstruction efficiency. To do so, we define the parameter

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

(3.48)

Therefore, the noise is increased by a factor of $\sqrt{\beta}$, over the direct photon QCD contribution. Our significant results are encapsulated in Fig. 7 where we show the discovery reaches of the LHC for several integrated luminosities and $\kappa^2 = 0.02$. A detailed study of the CMS potential for isolation of prompt-$\gamma$’s has been recently carried out [91], using GEANT4 simulations of $\gamma +$ jet events generated with Pythia. This analysis (which also includes $\gamma$’s produced in the decays of $\eta$, $K^0_s$, $\omega^0$, and bremsstrahlung photons emerging from high-$p_{\perp}$ jets) suggests $\beta \simeq 2$. Of course, considerations of detector efficiency further reduce the $S/N$ ratio by an additional factor $\epsilon$, where $1 < \epsilon \ll \sqrt{\beta}$. We conclude that discovery at the LHC would be possible for $M_s$ as large as 2.3 TeV.

We now consider gluon fusion into two photons. As can be seen in Eq. (3.18), the string amplitude for diphoton production is suppressed by a factor of $Q^4_c$. On the other hand, the QCD leading order contribution to diphoton production, given by the Born level process $q \bar{q} \rightarrow \gamma \gamma$, appears with an additional $\alpha_{em}$ so that it nearly compensates the extra factor $\kappa^2$. However, the restriction to $C^{0s}$ of the intermediate state, introduces an important dependence on the unknown $C^0$ mass, because the pole is shifted away from $M_s$. Additionally, the widths of the $C^{0s}$ excitation, $\Gamma^{J=0} = 150 \alpha_s (M_s/\text{TeV})$ GeV and $\Gamma^{J=2} = 75 \alpha_s (M_s/\text{TeV})$ GeV [83], are nearly twice that of the $G^*$. These two considerations vitiate any useful sensitivity of the diphoton channel.

We now briefly explore the potential of the ALICE to search for low mass string excitations. With this motivation, we extend our analysis to include heavy ions collisions.

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\footnote{Pb-Pb $\rightarrow \gamma +$ jet events can be identified by selecting a prompt photon and searching for the leading jet.}
Figure 6: Cross section for gluon fusion into $\gamma + \text{jet}|_{p_T(\gamma) > 300 \text{ GeV}}$ and expected number of events, for $100 \text{ fb}^{-1}$ and varying string scale [82].
Figure 7: Contours of 5σ discovery in the \((M_s, \text{detector efficiency})\) plane for different integrated luminosities and \(\kappa^2 = 0.02\).
the spirit of Ref. [93] we consider the unshadowed parton distribution functions, i.e.,

\[ R_{i/A}(x) = \frac{f_{i/A}(x, Q)}{Af_i(x, Q)} \simeq 1, \] (3.49)

where \( f_{i/A} \) and \( f_i \) are the parton distribution functions inside a free nucleus of mass \( A \) and free nucleon, respectively. For \( M_s > 1 \) TeV, this approximation holds because the LHC Pb-Pb collisions probe the minimum value of parton momentum at \( x_{\text{min}} \approx M_s^2/s \sim 0.033 \), where there are no shadowing effects. A comparison of the string cross section for gluon fusion into \( \gamma + \text{jet}_{k_{\perp}^\gamma > 300 \text{ GeV}} \) for \( pp \) and Pb-Pb collisions is shown in Fig. 5. However, the larger aggregate of partons also increase the SM background; namely, for \( k_{\perp,\text{min}} > 300 \) GeV, \( \sigma_{\text{Pb-Pb} \rightarrow \gamma X} \approx 2.8 \times 10^7 \) fb. This greatly decreases the sensitivity to D-brane models, which would require a Pb-Pb integrated luminosity of a few hundred pb\(^{-1}\). This is substantially larger than the present day estimate [94].

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particle in the opposite direction inside the ALICE central tracking system [92]. As photons emerge almost unaltered from dense medium, they provide a measurement of the original energy of the parton emitted in the opposite direction.
Chapter 4

Strings vs. Black Holes at the LHC

If nature has gracefully picked a sufficiently low-scale gravity, microscopic black holes can be produced at particle accelerators \cite{95}. In particular, the cross section for black hole production at the LHC is expected to be $\sim 100 \text{ pb}$ for a fundamental Planck scale $M_{10} \sim 1 \text{ TeV}$, which could turn the LHC into a black hole factory with a production rate of $\sim 1 \text{ Hz}$ \cite{84, 96}. The LHC Hawking temperature would be few hundred GeV, and so black holes would quickly evaporate into about a half dozen particles with large transverse momentum. In this chapter we discuss potential methods to discriminate the high-$k_\perp$ string decay products from the light descendants of black holes. Interestingly, one of these methods allows an increase of the LHC sensitivity for Regge recurrences of fundamental strings up to about 4 TeV.

4.1 Bump-Hunting

The discovery trigger described in the previous chapter, the observation of isolated photons at large transverse momentum, serves very well as a signature of new physics. However, as mentioned above, this criterion served also as a marker for Hawking radiation following production of TeV-scale black holes at the LHC. Given the particular nature of the string process we are considering, the production of a TeV-scale resonance and its subsequent 2-body decay, signatures in addition to large $k_\perp$ photons are available. Most apparently, one would hope that the resonance would be visible in data binned according to the invariant mass $M$ of the photon + jet, setting cuts on photon and jet rapidities, $|y_1|, |y_2| < y_{\text{max}} = 2.4$, respectively. With the definitions $Y \equiv \frac{1}{2}(y_1 + y_2)$ and $y \equiv \frac{1}{2}(y_1 - y_2)$, the cross section per
interval of $M$ for $pp \rightarrow \gamma + \text{jet} + X$ is given by

$$\frac{d\sigma}{dM} = \frac{M^3}{s} \sum_{ijk} \left[ \int_{-Y}^{0} dY f_i(x_a, M) f_j(x_b, M) \frac{dy}{\cosh^2 y} \frac{d\sigma}{dt} \bigg|_{ij \rightarrow k,\gamma} \right] \right] \right]$$

(4.1)

where $i, j, k$ are different partons, and the longitudinal fractions have the forms

$$x_a = M e^Y / \sqrt{s}, \quad x_b = M e^{-Y} / \sqrt{s}$$

(4.2)

(see Appendix D for details). The kinematics of the scattering provides the relation

$$k_\perp = \frac{M}{2 \cosh y},$$

(4.3)

which, when combined with the standard cut $k_\perp > k_{\perp,\text{min}}$, imposes a lower bound on $y$ to be implemented in the limits of integration. The $Y$ integration range in Eq. (4.1), $Y_{\text{max}} = \min\{\ln(\sqrt{s}/M), \ y_{\text{max}}\}$, comes from requiring $x_a, x_b < 1$ together with the rapidity cuts $|y_1|, |y_2| \leq 2.4$. Finally, the Mandelstam invariants occurring in the cross section are given by

$$\hat{s} = M^2,$$

$$\hat{t} = -M^2 e^{-y} / 2 \cosh y,$$

$$\hat{u} = -M^2 e^{+y} / 2 \cosh y.$$  

(4.4)

In Fig. 8 we show several representative plots of this cross section for different values of $M_s$. Standard bump-hunting methods, such as calculating cumulative cross sections

$$\sigma(M_0) = \int_{M_0}^\infty \frac{d\sigma}{dM} dM$$

(4.5)

and searching for regions with significant deviations from the QCD background, may allow for finding 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 $[M_s - 2\Gamma, M_s + 2\Gamma]$. This estimate of the signal would be roughly the same as that obtained through the inclusive cut $k_\perp > 300$ GeV. This follows from the relation (4.3): for $M$ in the range of $M_s > 2$ and for the significant contributing regions of $y$, the resulting $k_\perp$ cut in Eq. (4.3) does not differ significantly from the estimated 300 GeV. However, for $M_s > 2$ TeV, the background is significantly reduced, augmenting
Figure 8: $d\sigma/dM$ (units of fb/GeV) vs. $M$ (TeV) is plotted for the case of the SM QCD background (dashed) and (first resonance) string signal + background (solid) [82].
Figure 9: Signal-to-noise ratio for an integrated luminosity of 100 fb$^{-1}$ and $\beta = 2$. The solid line is for $\kappa^2 = 0.02$, the dot-dashed line is for $\kappa^2 = 0.05$, and the dashed line is for an optimistic case with $\kappa^2 = 0.1$ [95].
the LHC discovery reach. In Fig. 9 we show the signal-to-noise for different values of the mixing parameter $\kappa$, assuming $\beta = 2$. It is clearly seen that even for relatively small mixing, 100 fb$^{-1}$ of the LHC data could probe deviations from the SM physics associated with TeV-scale strings at a 5$\sigma$ significance, for $M_s \lesssim 4$ TeV. Should bumps be found, the D-brane model can be further differentiated from other TeV-scale resonant processes by the details of the angular distributions inherent in Eq. (3.20).

4.2 Z’s

Analytic [99] and numerical [100] studies have revealed that gravitational collapse takes place only at sufficiently high energies and small impact parameters, as conjectured years ago by Thorne [101]. A horizon forms when and only when a mass is compacted into a hoop whose circumference in every direction is less than $2\pi$ times its Schwarzschild radius up to a factor of order 1.

The LHC black holes would decay largely via the Hawking process [96], in which both the average number [102, 103] and the probability distribution of the number [104, 105, 106] of outgoing particles in each mode obey a thermal spectrum. In 10-dimensions, the emission rate per degree of particle freedom $i$ of particles of spin $s$ with initial total energy between $(\omega, \omega + d\omega)$ is found to be [107]

$$\frac{\dot{N}_i}{d\omega} = \frac{\sigma_s(\omega)\Omega_{d-3}\omega^{d-2}}{(d-2)(2\pi)^{d-1}} \left[ e^{\omega/T} - (-1)^{2s} \right]^{-1},$$ (4.6)

where $T = 7/(4\pi r)$ is the instantaneous Hawking temperature,

$$\Omega_{d-3} = \frac{2\pi^{(d-2)/2}}{\Gamma((d-2)/2)}$$ (4.7)

is the volume of a unit $(d-3)$-sphere,

$$r = \frac{1}{M_{10}} \left[ \frac{M}{M_{10}} 8 \pi^{3/2} \Gamma(9/2) \right]^{1/7}$$ (4.8)

is the instantaneous Schwarzschild radius of mass $M$ [108], and $\sigma_s(\omega)$ is the greybody absorption area due to the backscattering of part of the outgoing radiation of frequency $\omega$ into the black hole (a.k.a. the greybody factor) [109]. The SM fields live on a 3-brane ($d = 4$), while gravitons inhabit the entire spacetime ($d = 10$). The prevalent energies of the decay quanta are of $O(T \sim 1/r)$, resulting in $s$-wave dominance of the final state. Indeed, as the total angular momentum number of the emitted field increases, $\sigma_s(\omega)$ is rapidly suppressed [110, 111, 112, 113]. In the low energy limit, $\omega r \ll 1$, higher-order
terms are suppressed by a factor of $3(\omega r)^{-2}$ for fermions and by a factor of $25(\omega r)^{-2}$ for gauge bosons. For an average particle energy $\langle \omega \rangle$ of $\mathcal{O}(r^{-1})$, higher partial waves are also suppressed, although by a smaller factor. This strongly suggests that the black hole is sensitive only to the radial coordinate and does not make use of the extra angular modes available in the internal space \cite{114}. Actually, a recent detailed analysis \cite{115,116} has explicitly shown that the relative emission rate of the SM particles and the 10-dimensional bulk graviton is roughly 92:5. This implies that the power lost in the bulk is less than 15% of the total black hole mass, largely favoring the dominance of visible decay. Therefore, in what follows, we assume the Hawking evaporation process to be dominated by the SM brane modes and we neglect graviton emission during the Schwarzschild phase.

Altogether, the average total emission rate for particle species $i$ is

$$\frac{d\langle N \rangle}{dt} = \frac{1}{2\pi} \left( \sum c_i g_i \Gamma_i \right) \zeta(3) \Gamma(3) r^2 T^3,$$

(4.9)

where $c_i$ is the number of internal degrees of freedom of particle species $i$, $g_i = 1 (3/4)$ for bosons (fermions),

$$\Gamma_i = \frac{1}{4\pi r^2} \int \frac{\sigma_s(\omega) \omega^2 d\omega}{e^{\omega/T} \pm 1} \left[ \int \frac{\omega^2 d\omega}{e^{\omega/T} \pm 1} \right]^{-1},$$

(4.10)

and $\Gamma_i = 0.60$ ($\Gamma_i = 0.66$) for bosons (fermions) \cite{117}. This implies that black holes decay with roughly equal probability to all degrees of freedom of the SM particles. Since there are six charged leptons, one $Z$ boson, and one photon, we expect $\sim 10\%$ of the particles to be hard primary leptons and 2% of the particles to be hard photons and $Z$’s, each carrying hundreds of GeV of energy.

We now discuss some interesting contrast of $\gamma$ and $Z$ production in $D$-brane models that can serve as an additional marker for discovery of string recurrences. Ignoring the $Z$-mass (i.e., keeping only transverse $Z$’s), and assuming that cross sections $\times$ branching into lepton pairs are large enough for complete reconstruction to $pp \rightarrow Z + \text{jet}$, the quiver contribution to the signal is suppressed relative to the photon signal by a factor of $\tan^2 \theta_W = 0.29$. The SM ratio ($Z$ background)/($\gamma$ background) is roughly 0.92 for processes involving $u$ (or $\bar{u}$) quarks, and 4.7 for processes involving $d$ (or $\bar{d}$) quark. Thus, even if $d$ quark processes are ignored, one obtains a signal-to-noise ratio $(S/N)_Z = 0.29/\sqrt{0.92} = 0.30 (S/N_\gamma)$. Keeping the $d$ quarks will only lead to more suppression of $(S/N)_Z$. This implies that if the high-$k_{\perp}$ photons, as predicted by the TeV string model, are discovered at 5$\sigma$, they will not be

\footnote{It is worth pausing to note that $\pi^0$ misidentification does not play a role in the $Z$ channel, and so this tends to decrease the QCD background. On the other hand, the string signal will suffer some suppression because of finite mass effects. These systematics (which have opposite effects on $(S/N)_Z$) were not considered in the preceding discussion.}
accompanied by any significant deviation of $pp \rightarrow Z + \text{jet}$ from the SM predictions. This differs radically from the evaporation of black holes produced at the LHC. In that case, production of high-$k_\perp Z$ and $\gamma$ are comparable. The suppression of high-$k_\perp Z$ production, whose origin lies in the particular structure of the quiver model, will hold true for all the low-lying levels of the string.
Chapter 5

Dijet Signals of Low Mass Strings at the LHC

The string amplitudes that involve four matter fields depend on the details of the D-brane geometry and how the D-branes are embedded into the compact Calabi-Yau space. This is because modes of the internal geometry can be exchanged during the four fermion scattering processes. However, as shown in Ref. [9], the poles of the amplitudes for two gauge bosons and two matter fermions are due to the exchanges of massless gauge bosons and universal string Regge excitations only, and so computations of the respective average square amplitudes can be performed in a model independent and universal way. In this chapter, we extend our search for string signals at the LHC, by including scattering processes which involve four gauge bosons as well as two gauge bosons and two fermions.

5.1 \( pp \rightarrow \text{dijet} \)

As noted in Ref. [118], string signals are likely to show up in the dijet channel. 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 \( 2 \rightarrow 2 \) scattering amplitudes, computed at the leading order in string perturbation theory, are collected in Ref. [9]. The average square amplitudes are given by the following:

\[
|M(gg \rightarrow gg)|^2 = g^4 \left( \frac{1}{s^2} + \frac{1}{t^2} + \frac{1}{u^2} \right) \left[ \frac{2N^2}{N^2-1} (s^2 V_s^2 + t^2 V_t^2 + u^2 V_u^2) + \frac{4(3-N^2)}{N^2(N^2-1)} (s V_s + t V_t + u V_u)^2 \right],
\]  

(5.1)
\[ |\mathcal{M}(gg \to q\bar{q})|^2 = g^4 N_f \frac{t^2 + u^2}{s^2} \left[ \frac{1}{2N u t} (t V_t + u V_u)^2 - \frac{N}{N^2 - 1} V_t V_u \right], \]  
(5.2)

\[ |\mathcal{M}(q\bar{q} \to gg)|^2 = g^4 \frac{t^2 + u^2}{s^2} \left[ \frac{(N^2 - 1)^2}{2N^3} \frac{1}{u t} (t V_t + u V_u)^2 - \frac{N^2 - 1}{N} V_t V_u \right], \]  
(5.3)

and

\[ |\mathcal{M}(gg \to gg)|^2 = g^4 \frac{s^2 + u^2}{t^2} \left[ V_s V_u - \frac{N^2 - 1}{2N^2} \frac{1}{s u} (s V_s + u V_u)^2 \right], \]  
(5.4)

where the string “formfactor” functions of the Mandelstam variables are defined as

\[ V_t = V(s, t, u), \quad V_u = V(t, u, s), \quad V_s = V(u, s, t), \]  
(5.5)

with

\[ V(s, t, u) = \frac{\Gamma(1 - s/M_s^2) \Gamma(1 - u/M_s^2)}{\Gamma(1 + t/M_s^2)}. \]  
(5.6)

Before proceeding, we pause to review 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 \( C^0 \) and its excitations to accompany the gluon and its excitations. The first excitation of the \( C^0 \) will be denoted by \( C^{0*} \).

In the following we isolate the contribution from the first resonant state in Eqs. (5.1) - (5.4). For partonic center of mass energies \( \sqrt{s} < M_s \), contributions from the Veneziano functions are strongly suppressed, as \( \sim (\sqrt{s}/M_s)^8 \), over standard model processes; see Eq. (5.19). (Corrections to SM processes at \( \sqrt{s} \ll M_s \) are of order \( (\sqrt{s}/M_s)^4 \).) In order to factorize amplitudes on the poles due to the lowest massive string states, it is sufficient to consider \( s = M_s^2 \). In this limit, \( V_s \) is regular while

\[ V_t = \frac{u}{s - M_s^2}, \quad V_u = \frac{t}{s - M_s^2}. \]  
(5.7)

Thus the s-channel pole term of the average square amplitude (5.1) can be rewritten as

\[ |\mathcal{M}(gg \to gg)|^2 = 2 \frac{g^4 M_s^4}{M_s^8} \left( \frac{N^2 - 4 + (12/N^2)}{N^2 - 1} \right) \frac{M_s^8 + t^4 + u^4}{(s - M_s^2)^2}. \]  
(5.8)

\(^1\)Note that the contributions of single poles to the cross section are antisymmetric about the position of the resonance, and vanish in any integration over the resonance. Let the amplitude be \( a + b/D \) in the vicinity of the pole, where \( a \) and \( b \) are real, \( D = x + i \epsilon, x = s - M_s^2, \) and \( \epsilon = \Gamma M_s. \) Then, since \( \text{Re}(1/D) = x/|D|^2, \) the cross section becomes \( \sigma \propto a^2 + b^2/|D|^2 + 2 a b x/|D|^2 \approx a^2 + b^2 \pi \delta(x)/\epsilon + 2 a b \pi x \delta(x)/\epsilon. \) Integrating over the width of the resonance, one obtains \( a^2 \epsilon + b^2 \pi /\epsilon \approx b \pi, \) because \( b \propto \epsilon, a \propto g^2 \) and \( \epsilon \propto g^2. \)
The singularity at \( s = M_s^2 \) needs softening to a Breit-Wigner form, reflecting the finite decay widths of resonances propagating in the s channel. Due to averaging over initial polarizations, Eq. (5.8) contains additive contributions from both spin \( J = 0 \) and spin \( J = 2 \) \( U(3) \) bosonic Regge recurrences \((G^* \text{ and } C^{0*})\) in the notation of Ref. [83], created by the incident gluons in the helicity configurations \((\pm \pm)\) and \((\pm \mp)\), respectively. The \( M_s^8 \) term in Eq. (5.8) originates from \( J = 0 \), and the \( t^4 + u^4 \) piece reflects \( J = 2 \) activity. Since the resonance widths depend on the spin and on the identity of the intermediate state \((G^*, C^{0*})\) the pole term (5.8) should be smeared as

\[
|M(gg \to gg)|^2 = 2 \frac{g^4}{M_s^4} \left( \frac{N^2 - 4 + (12/N^2)}{N^2 - 1} \right)
\]

\[
\times \left\{ W_{G^*}^{gg \to gg} \left[ \frac{M_s^8}{(s - M_s^2)^2 + (\Gamma_{G^*}^{J=0} M_s)^2} + \frac{t^4 + u^4}{(s - M_s^2)^2 + (\Gamma_{G^*}^{J=2} M_s)^2} \right] \right. 
\]

\[\left. + W_{C^{0*}}^{gg \to gg} \left[ \frac{M_s^8}{(s - M_s^2)^2 + (\Gamma_{C^{0*}}^{J=0} M_s)^2} + \frac{t^4 + u^4}{(s - M_s^2)^2 + (\Gamma_{C^{0*}}^{J=2} M_s)^2} \right] \right\},
\]

where \( \Gamma_{G^*}^{J=0} = 75 (M_s/\text{TeV}) \text{ GeV} \), \( \Gamma_{C^{0*}}^{J=0} = 150 (M_s/\text{TeV}) \text{ GeV} \), \( \Gamma_{G^*}^{J=2} = 45 (M_s/\text{TeV}) \text{ GeV} \), and \( \Gamma_{C^{0*}}^{J=2} = 75 (M_s/\text{TeV}) \text{ GeV} \) are the total decay widths for intermediate states \( G^* \) and \( C^{0*} \), with angular momentum \( J \). The associated weights of these two intermediate states are given in terms of the probabilities for the various entrance and exit channels

\[
W_{G^*}^{gg \to gg} = \frac{(\Gamma_{G^*}^{J=0} - GG)^2}{(\Gamma_{G^*}^{J=0} - GG)^2 + (\Gamma_{C^{0*} - GG}^{J=0})^2} = 0.09 \quad (5.10)
\]

and

\[
W_{C^{0*}}^{gg \to gg} = \frac{(\Gamma_{C^{0*} - GG}^{J=0})^2}{(\Gamma_{G^*}^{J=0} - GG)^2 + (\Gamma_{C^{0*} - GG}^{J=0})^2} = 0.91 \quad (5.11)
\]

A similar calculation transforms Eq. (5.2) near the pole into

\[
|M(gg \to q\bar{q})|^2 = \frac{g^4}{M_s^4} N_f \left( \frac{N^2 - 2}{N(N^2 - 1)} \right) \left[ W_{G^*}^{gg \to q\bar{q}} \frac{ut(u^2 + t^2)}{(s - M_s^2)^2 + (\Gamma_{G^*}^{J=2} M_s)^2} \right.
\]

\[\left. + W_{C^{0*}}^{gg \to q\bar{q}} \frac{ut(u^2 + t^2)}{(s - M_s^2)^2 + (\Gamma_{C^{0*}}^{J=2} M_s)^2} \right], \quad (5.12)
\]

where

\[
W_{G^*}^{gg \to q\bar{q}} = W_{G^*}^{qq \to gg} = \frac{\Gamma_{G^*}^{J=0} - GG \Gamma_{G^*}^{J=0} - q\bar{q}}{\Gamma_{G^*}^{J=0} - GG \Gamma_{G^*}^{J=0} - q\bar{q} + \Gamma_{C^{0*} - GG}^{J=0} \Gamma_{C^{0*} - q\bar{q}}} = 0.24 \quad (5.13)
\]

and

\[
W_{C^{0*}}^{gg \to q\bar{q}} = W_{C^{0*}}^{qq \to gg} = \frac{\Gamma_{C^{0*} - GG}^{J=0} \Gamma_{C^{0*} - q\bar{q}}}{\Gamma_{G^*}^{J=0} - GG \Gamma_{G^*}^{J=0} - q\bar{q} + \Gamma_{C^{0*} - GG}^{J=0} \Gamma_{C^{0*} - q\bar{q}}} = 0.76 \quad (5.14)
\]
Near the $s$ pole Eq. (5.3) becomes

$$|\mathcal{M}(q\bar{q} \rightarrow gg)|^2 = \frac{g^4}{M_s^4} \left( \frac{(N^2 - 2)(N^2 - 1)}{N^3} \right) \left[ W_{qq\rightarrow gg}^{s} \frac{ut(u^2 + t^2)}{(s - M_s^2)^2 + (\Gamma^{J=2/M_s^2})^2} \right] \left[ W_{qq\rightarrow gg}^{c} \frac{ut(u^2 + t^2)}{(s - M_s^2)^2 + (\Gamma^{J=2/M_s^2})^2} \right],$$

(5.15)

whereas Eq. (5.4) can be rewritten as

$$|\mathcal{M}(qg \rightarrow qg)|^2 = \frac{g^4}{M_s^4} \left( \frac{N^2 - 1}{2N^2} \right) \left[ M_s^4 u \frac{1}{(s - M_s^2)^2 + (\Gamma^{J=1/M_s^2})^2} \right] \left[ M_s^4 u \frac{1}{(s - M_s^2)^2 + (\Gamma^{J=3/M_s^2})^2} \right].$$

(5.16)

The total decay widths for the $q^*$ excitation are: $\Gamma^{J=1/2}_{q^*} = \Gamma^{J=3/2}_{q^*} = 37 (M_s/\text{TeV}) \text{ GeV}$ [83]. Superscripts $J = 2$ are understood to be inserted on all the $\Gamma$’s in Eqs. (5.10), (5.11), (5.13), (5.14). Equation (5.9) reflects the fact that weights for $J = 0$ and $J = 2$ are the same [83]. In what follows we set $N = 3$ and $N_f = 6$.

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, $|y_1|, |y_2| \leq 1$ [119] and transverse momenta $p_T^{1,2} > 50 \text{ GeV}$. In Fig. 10 we show a representative plot of the invariant mass spectrum, for $M_s = 2 \text{ TeV}$, 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 addition of the $t$-channel exchange process $qq \rightarrow qq$. Our calculation, making use of the CTEQ6 parton distribution functions [85, 86] agrees with that presented in [119].

We now estimate (at parton level) the LHC discovery reach. To do so, we calculate a signal-to-noise ratio, with the signal rate estimated in the invariant mass window $[M_s - 2\Gamma, M_s + 2\Gamma]$. As usual, the noise is defined as the square root of the number of background events in the same dijet events in the same integrated luminosity.

The top two and bottom curves in Fig. 11 show the behavior of the signal-to-noise ($S/N$) ratio as a function of the string scale 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* [3]. 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 ($S/N > 13$). The bottom curve, corresponding data collected in a very early

---

脚注：

2 This intersects with the range of string scales consistent with correct weak mixing angle found in the minimal quiver standard model [66].
Figure 10: $d\sigma/dM$ (units of fb/GeV) vs. $M$ (TeV) is plotted for the case of the 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 ($gg \rightarrow gg$, $gg \rightarrow q\bar{q}$, $gg \rightarrow qg$, and $q\bar{q} \rightarrow gg$) [20].
run of 100 $pb^{-1}$, shows that a resonant mass as large as 4.0 TeV can be observed with 10$\sigma$ significance! Once more, we stress that these results contain no unknown parameters. They depend only on the D-brane construct for the standard model, and are independent of compactification details.

The amplitudes for the four-fermion processes like quark-antiquark scattering are more complicated because the respective formfactors 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, see Fig. 5. The collisions of valence quarks occur at higher luminosity; however, there are no Regge recurrences appearing in the $s$-channel of quark-quark scattering [9].

5.2 $pp \rightarrow \gamma + \text{jet}$

In this section we estimate corrections from scattering of two gauge bosons and two matter fermions to the $pp \rightarrow \gamma + \text{jet}$ channel. From our dijet analysis in the previous section and Fig. 5, it is easily seen that the average square amplitude dominating $pp \rightarrow \gamma + \text{jet}$ reads [9]

$$|M(qg \rightarrow q\gamma)|^2 = -g^2 Q_c^2 \frac{1}{N} \frac{s^2 + u^2}{s u t^2} (s V_s + u V_u)^2,$$

where, as we defined in Chapter 3, $Q_c$ is the product of the $U(1)$ charge of the fundamental representation ($\sqrt{1/6}$) followed by successive projections onto the hypercharge and then onto the photon ($\cos \theta_W$). For comparison with our dijet analysis, we also show in Fig. 11 a fourth curve, for the process $pp \rightarrow \gamma + \text{jet}$, taking into account all the possible contributions. The approximate equality of the background due to misidentified $\pi^0$’s and the QCD background, across a range of large $k_\perp$ as implemented in Chapter 4, is maintained as an approximate equality over a range of invariant $\gamma$-jet invariant masses with the rapidity cuts imposed. When considering contributions from scattering processes with two gauge bosons and two fermions, the LHC discovery reach in the $pp \rightarrow \gamma + \text{jet}$ channel is extended up to $M_\gamma \sim 5.0$ TeV.
Figure 11: $pp \to \text{dijet}$ signal-to-noise ratio for three integrated luminosities. For comparison, we also show the signal-to-noise of $pp \to \gamma + \text{jet}$, for the minimal quiver standard model [120].
Chapter 6

Conclusion

In the first part of this Dissertation, we studied the six dimensional Salam-Sezgin model, where a solution of the form $\text{Minkowski}_4 \times S^2$ is known to exist, with a $U(1)$ monopole serving as background in the two-sphere. This model circumvents the hypotheses of the no-go theorem of Maldacena and Nuñez, and then when lifted to string theory, can show a dS phase. In our analysis we have allowed for time dependence of the six-dimensional moduli fields and metric (with a Robertson-Walker form). Time dependence in these fields vitiates invariance under the supersymmetry transformations. With these constructs, we have obtained the following results:

- In terms of linear combinations of the $S^2$ moduli field and the six dimensional dilaton, the effective potential consists of $(a)$ a pure exponential function of a quintessence field (this piece vanishes in the supersymmetric limit of the static theory) and $(b)$ a part which is a source of cold dark matter, with a mass proportional to an exponential function of the quintessence field. This presence of a VAMP CDM candidate is inherent in the model.

- If the monopole strength is precisely at the value prescribed by supersymmetry, the model is in gross disagreement with present cosmological data - there is no accelerative phase, and the contribution of energy from the quintessence field is purely kinetic. However, a miniscule deviation of $O(10^{-120})$ from this value permits a qualitative match with data. Contribution from the VAMP component to the matter energy density can be as large as about 7% without having negative impact on the fit. The emergence of a VAMP CDM candidate as a necessary companion of dark energy has been a surprising aspect of the present findings, and perhaps encouraging for future exploration of candidates which can assume a more prominent role in the CDM sector.
• In our model, the exponential potential has behavior \( V_Y \sim e^{\lambda Y / M_{Pl}} \), with \( Y \) the quintessence field and \( \lambda = \sqrt{2} \). The asymptotic behavior of the scale factor for exponential potentials is \( e^{h(t)} \approx t^{2/\lambda^2} \), so that for our case \( h \approx \ln t \), leading to a conformally flat Robertson-Walker metric for large times. The evolution from constant velocity expansion to a brief accelerated phase in the neighborhood of our era makes the model phenomenologically viable. In the case that the supersymmetry condition \((b^2 \xi = 1)\) is imposed, and there is neither radiant energy nor dark matter except for the \( X \) contribution, we find for large times that the growth of the scale factor is given by \( e^{h(t)} \approx \sqrt{t} \), so that even in this case the asymptotic metric is Robertson-Walker rather than Minkowski. Moreover, and rather intriguingly, the scale factor is what one would find with radiation alone.

In sum, in spite of the shortcomings of the model (not a perfect fit, requirement of a tiny deviation from supersymmetric prescription for the monopole embedding), it has provided stimulating new, and unifying, look at the dark energy and dark matter puzzles.

On a separate track, the LHC program will include the identification of events with single prompt high-\( k_T \) photons as probes of new physics. In the second part of this Thesis, we have shown that this channel is uniquely suited to search for experimental evidence of TeV-scale open string theory. At the parton level, we analyzed single photon production in gluon fusion, \( gg \rightarrow \gamma g \), 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 (which is actually determined in the minimal theory) – and it is otherwise model-(compactification-) independent. We calculated cross sections for Regge recurrences of fundamental strings, as well as the QCD background. (A vital part of the background discussion concerned the minimization of misidentified \( \pi^0 \)'s emerging from high-\( p_T \) jets.) We showed that even for relatively small mixing, 100 fb\(^{-1}\) of the LHC data could probe deviations from the SM physics associated with TeV-scale strings at a 5\( \sigma \) significance, for \( M_{\text{string}} \) as large as 4 TeV.

Another channel that can provide a clean signal of new physics at the LHC is \( pp \rightarrow \) dijet. 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 made use of these amplitudes evaluated near the first resonant pole to determine the discovery potential of the LHC for the first Regge excitations of the quark and gluon. We found that,
remarkably, the reach of the 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. For string scales as high as 5.0 TeV, observations of resonant structures in $pp \to \gamma + \text{jet}$ (considering parton subprocesses $gg \to \gamma g$, $gq \to \gamma q$, $g\bar{q} \to \gamma \bar{q}$, and $q\bar{q} \to \gamma g$), can provide interesting corroboration of string physics at the TeV-scale.

All in all, a new era, with experimental measurements of string physics, may be close at hand.
Appendix A

Proof of the no-go theorem

For completeness, in this Appendix, we provide the proof of the no-go theorem of Maldacena and Nunez [36].

**Theorem of Maldacena and Nunez:**
Consider a $D$ dimensional gravity theory which is compactified on $d$ dimensions. If the $D$ dimensional gravity theory satisfies the following conditions:

1. the gravity action does not contain higher curvature corrections;
2. the potential is non-positive, $V \leq 0$;
3. the theory contains massless fields with positive kinetic terms;
4. the internal manifold is compact without boundary and its volume is finite;

then there are no non-singular spontaneous compactifications to Minkowwski or de-Sitter of the form

$$ds^2 = \Omega^2(y) (dx_d^2 + \hat{g}_{mn} dy^m dy^n), \quad (A.1)$$


Proof:
From assumption 1, we can write Einstein’s equations in $D$ dimensions

$$R_{MN} = T_{MN} - \frac{1}{D-2} g_{MN} T_{L}^{L}. \quad (A.2)$$
Using the metric form, Einstein’s equations can be re-written as

$$ R_{\mu\nu} = R_{\mu\nu}(\eta) - \eta_{\mu\nu}\left(\nabla^2 \log \Omega + (D - 2)(\nabla \log \Omega)^2\right) $$

$$ = T_{\mu\nu} - \frac{1}{D - 2} \Omega^2 \eta_{\mu\nu} T^L_L, \quad (A.3) $$

where $\nabla$ is the covariant derivative operator of the metric $\hat{g}$, and its indices are contracted with $\hat{g}$. Taking the trace over $\eta$ on both sides we find

$$ \nabla^2 \log \Omega + (D - 2)(\nabla \log \Omega)^2 = \frac{1}{(D - 2)\Omega^{D-2}} \nabla^2 \Omega^{D-2} $$

$$ = R(\eta) + \Omega^2 \left(-T^\mu_\mu + \frac{d}{D - 2} T^L_L\right). \quad (A.4) $$

Now, we define

$$ \tilde{T} \equiv -T^\mu_\mu + \frac{d}{D - 2} T^L_L. \quad (A.5) $$

First, we shall show that $\tilde{T}$ is non-negative. For a given potential of matter fields $V$, we have $T_{MN} \sim -V g_{MN}$, and

$$ \tilde{T} \sim V d - \frac{d}{D - 2} DV = -\frac{2d}{D - 2} V \geq 0, \quad (A.6) $$

where the last equality follows from assumption 2. The energy momentum tensor of $n$-form fields takes the form

$$ T_{MN} = F_{M L_1 \cdots L_{n-1}} F^L_N L_1 \cdots L_{n-1} - \frac{1}{2n} g_{MN} F^2, \quad (A.7) $$

which in turn gives

$$ T^\mu_\mu = F_{\mu L_1 \cdots L_{n-1}} F^{\mu L_1 \cdots L_{n-1}} - \frac{d}{2n} F^2. \quad (A.8) $$

Hence,

$$ \tilde{T} = -F_{\mu L_1 \cdots L_{n-1}} F^{\mu L_1 \cdots L_{n-1}} + \frac{d}{D - 2} \left(1 - \frac{1}{n}\right) F^2. \quad (A.9) $$

The space time indices of non-vanishing components of $F$ could be completely along the internal dimensions or, if $n \geq d$, they could have $d$ out of $n$ indices along the $d$ dimensions and the rest along the internal dimensions. Otherwise the isometry of $R^d$ or $dS^d$ is broken. They separately contribute to $\tilde{T}$. In the former case, $F^2 \geq 0$. It follows from Eq. (A.8) that $\tilde{T} \geq 0$ for $n > 1$ and $\tilde{T} = 0$ for $n = 1$. In the latter case, $F^2 < 0$ and

$$ F_{\mu L_1 \cdots L_{n-1}} F^{\mu L_1 \cdots L_{n-1}} = \frac{d}{n} F^2. \quad (A.10) $$

It again follows

$$ \tilde{T} = \left[-\frac{d}{n} + \frac{d}{D - 2} \left(1 - \frac{1}{n}\right)\right] F^2 = -\frac{d(D - 2 - n + 1)}{n(D - 2)} F^2 \geq 0. \quad (A.11) $$
Consequently we have in general
\[ \hat{T} \geq 0. \quad (A.12) \]

Combined with Eq. (A.1) and the assumption that our $d$-dimensional space is Minkowski or de Sitter with non-negative scalar curvature, this means that
\[ \Omega^{D-2} \hat{\nabla}^2 \Omega^{D-2} \geq 0. \quad (A.13) \]

The equality holds only if the right hand side of Eq. (A.1) is zero so that the $d$ dimensional space is Minsowski space. Since the internal space is compact, $\Omega$ is bounded below and above. Hence, integrating this over the internal space by parts, we obtain
\[ \int d^{D-d}y \sqrt{-\hat{g}} (\hat{\nabla} \Omega^{D-2})^2 \leq 0. \quad (A.14) \]

The left hand side is positive-definite so that this is valid only if $\Omega$ is constant and the equality holds. This implies that the right hand side of Eq. (A.1) vanishes, hence dS space is not allowed and the only $n$ forms that can be turned on are the $n = 1$, $D - 1$ forms. On the other hand, since the spontaneous compactification does not allow $\Omega = \text{constant}$, Minkowski space is also forbidden. In addition, it follows that the effective Newton constant is finite since the $d$ dimensional Newton constant is given by
\[ \frac{1}{G_N^d} \sim \int d^d y \sqrt{-g} \Omega^{(d-2)}. \quad (A.15) \]
Appendix B

The String Connection

In this Appendix we briefly comment on how the six dimensional solution derived in Chapter 2 reads in string theory. To this end, we use the uplifting formulae developed by Cvetic, Gibbons and Pope [42]; we will denote with the subscript “cgp” the quantities of that paper and with “us,” quantities in this Dissertation. Let us more specifically look at Eq. (34) in Ref. [42], where the authors described the six dimensional Lagrangian they uplifted to Type I string theory. By simple inspection, we can see that the relation between their variables and fields with the ones we used in Eq. (2.2) is

\[ \text{\phi}_{\text{cgp}} = -2\text{\phi}_{\text{us}}, \quad F_2|_{\text{cgp}} = \sqrt{G_6}F_2|_{\text{us}}, \quad H_3|_{\text{cgp}} = \sqrt{G_6/3}G_3|_{\text{us}}, \quad \text{and} \quad \tilde{g}^2|_{\text{cgp}} = \xi/(8G_6)|_{\text{us}}. \]

Our six dimensional background is determined by

- the (string frame) metric
  \[ ds_6^2 = e^{2f}[dt^2 + e^{2h}dx_3^2 + r_c^2d\sigma_2^2], \]
- the gauge field
  \[ F_{\vartheta \phi} = -b \sin \vartheta, \]
- and the \( t \)-dependent functions
  \[ h(t) = \sqrt{G_4}(X - Y)/4, \quad f(t) = \sqrt{G_4}(X + Y)/2. \]

Identifying these expressions with those in Eqs. (47), (48) and (49) of Ref. [42], one obtains a full Type I or Type IIB configuration, consisting of a 3-form (denoted by \( F_3 \)),

\[
F_3 = \frac{8G_6 \sinh \hat{\rho} \cosh \hat{\rho}}{\xi \cosh^2 2\hat{\rho}} d\hat{\rho} \wedge \left( d\alpha - \sqrt{\frac{\xi}{8G_6}} b \cos \vartheta d\varphi \right) \wedge \left( d\beta + \sqrt{\frac{\xi}{8G_6}} b \cos \vartheta d\varphi \right) \\
- \frac{\sqrt{2}G_6b}{\sqrt{\xi} \cosh 2\hat{\rho}} \sin \vartheta d\vartheta \wedge d\varphi \wedge \left[ \cosh^2 \hat{\rho} \left( d\alpha - \sqrt{\frac{\xi}{8G_6}} b \cos \vartheta d\varphi \right) \right. \\
- \left. \sinh^2 \hat{\rho} \left( d\beta + \sqrt{\frac{\xi}{8G_6}} b \cos \vartheta d\varphi \right) \right], \quad (B.1)
\]

a dilaton (denoted by \( \hat{\phi} \))

\[ e^{2\hat{\phi}} = \frac{e^{2\phi}}{\cosh(2\hat{\rho})}, \quad (B.2)\]
and a ten dimensional metric that in the string frame reads

$$ds_{\text{str}}^2 = e^\phi ds_6^2 + dz^2 + \frac{4G_6}{\xi} \left[ d\hat{\rho}^2 + \frac{\cosh^2 \hat{\rho}}{\cosh 2\rho} \left( d\alpha - \sqrt{\frac{\xi}{8G_6}} b \cos \vartheta d\varphi \right)^2 + \frac{\sinh^2 \hat{\rho}}{\cosh 2\rho} \left( d\beta + \sqrt{\frac{\xi}{8G_6}} b \cos \vartheta d\varphi \right)^2 \right],$$

where $\hat{\rho}, z, \alpha,$ and $\beta$ denote the four extra coordinates. It is important to stress that though the uplifted procedure described above implies a non-compact internal manifold, the metric in Eq. (B.3) can be interpreted within the context of [40] (i.e., $0 \leq \hat{\rho} \leq L,$ with $L \gg 1$ an infrared cutoff where the spacetime smoothly closes up) to obtain a finite volume for the internal space and consequently a non-zero but tiny value for $G_6.$
Appendix C

Pole residues of the Veneziano form factor

Consider the product of Gamma functions

\[ \Gamma(n) \Gamma(1-n) = \frac{\pi}{\sin(n\pi)} . \]  \hspace{1cm} (C.1)

In the limit \( 1 - n = \epsilon \ll 1, \sin(n\pi) = \sin(\pi - \pi\epsilon) = \sin(\pi) - \pi\epsilon \cos(\pi) = \pi\epsilon, \) and so

\[ \Gamma(1 - \epsilon) \Gamma(\epsilon) = \frac{\pi}{\pi\epsilon} = \frac{1}{\epsilon} , \]  \hspace{1cm} (C.2)

which in turn leads to

\[ \lim_{n \to 1} \Gamma(1 - n) = \frac{1}{1 - n} . \]  \hspace{1cm} (C.3)

Therefore, in the limit of \( s \to 1, \)

\[ \mu(s, t, u) = \frac{\Gamma(1 - u)}{\Gamma(1 + t)} \frac{1}{(1 - s)} = \frac{\Gamma(2 + t)}{\Gamma(1 + t)} \frac{1}{1 - s} = \frac{1 + t}{1 + s} \]  \hspace{1cm} (C.4)

and

\[ \mu(s, u, t) = \frac{\Gamma(1 - t)}{\Gamma(1 + u)} \frac{1}{(1 - s)} = \frac{\Gamma(2 + u)}{\Gamma(1 + u)} \frac{1}{1 - s} = \frac{1 + u}{1 + s} . \]  \hspace{1cm} (C.5)

We can now expand the string squared amplitude,

\[ |\mathcal{M}(gg \to \gamma g)|^2 = \left| \frac{s}{u} \mu(s, t, u) + \frac{s}{t} \mu(s, u, t) \right|^2 + \left| \frac{t}{u} \mu(s, t, u) + \frac{t}{s} \mu(t, u, s) \right|^2 + \left| \frac{u}{s} \mu(u, t, s) + \frac{u}{t} \mu(u, s, t) \right|^2 , \]  \hspace{1cm} (C.6)
near the pole yielding

$$|\mathcal{M}(gg \rightarrow \gamma g)|^2 \propto \left| \frac{s_0}{u} \frac{\Gamma(1-u)}{\Gamma(1+t)} + \frac{s_0}{t} \frac{\Gamma(1-t)}{\Gamma(1+u)} \right|^2 \frac{1}{(s-s_0)^2}$$

$$+ \left| -\frac{t}{u} \frac{\Gamma(1-u)}{\Gamma(1+t)} + \frac{t}{s_0} \left[ \frac{\Gamma(1-t)}{\Gamma(1+u)} - \frac{\Gamma(1-u)}{\Gamma(1+t)} \right] \right|^2 \frac{1}{(s-s_0)^2}$$

$$+ \left| \frac{u}{s_0} \left( \frac{\Gamma(1-u)}{\Gamma(1+t)} - \frac{\Gamma(1-t)}{\Gamma(1+u)} \right) - \frac{u}{t} \frac{\Gamma(1-t)}{\Gamma(1+u)} \right|^2 \frac{1}{(s-s_0)^2} , \quad (C.7)$$

where we have restored the string scale, $s_0 = M_s^2$. Equivalently,

$$|\mathcal{M}(gg \rightarrow \gamma g)|^2 \propto \left\{ \left| \frac{s_0}{u} \frac{A}{u} + \frac{s_0}{t} \frac{B}{t} \right|^2 + \left| -\frac{t}{u} A + \frac{t}{s_0} (B-A) \right|^2 \right\} \frac{1}{(s-s_0)^2} , \quad (C.8)$$

where

$$A = \frac{\Gamma(1-u)}{\Gamma(1+t)} = 1 + t = -u \quad (C.9)$$

and

$$B = \frac{\Gamma(1-t)}{\Gamma(1+u)} = 1 + u = -t \quad (C.10)$$

are obtained from Eqs. (C.4) and (C.5). Then, Eq. (C.8) becomes

$$|\mathcal{M}(gg \rightarrow \gamma g)|^2 \propto 4s_0^2 + \left| t + \frac{t}{s_0} (-t + u) \right|^2 + \left| u + \frac{u}{s_0} (-u + t) \right|^2$$

$$\propto \left[ 4s_0^2 + \frac{4t^4 + 4u^4}{s_0^2} \right] \frac{1}{(s-s_0)^2} , \quad (C.11)$$

where we have used the Mandelstam relation: $u = -s_0 - t$. Finally, the singularity is smeared with a width $\Gamma$ to obtain Eq. (3.20).
Appendix D

Invariant mass spectrum

In this Appendix D, we shall derive the invariant mass formula Eq. (4.1). For this purpose, we write the total cross-section for the process $pp \to \gamma + \text{jet}$:

$$
\sigma_{| pp \to \gamma + \text{jet}} = \int_0^1 dx_a \int_0^1 dx_b \sum_{ijk} f_i(x_a) f_j(x_b) \, \sigma_{| ij \to k, \gamma} \tag{D.1}
$$

Relation (3.28) lets us convert the integral in Eq. (D.1) into an integral over the parameters $M^2, Y$. The Jacobian of the change of variables is

$$
\frac{\partial(M^2, Y)}{\partial(x_a, x_b)} = \begin{vmatrix}
    x_b s & x_b s \\
    \frac{1}{2} x_a & -\frac{1}{2} x_b
\end{vmatrix} = s. \tag{D.2}
$$

Hence, we obtain

$$
\frac{d\sigma}{dM^2}_{| pp \to \gamma + \text{jet}} = \frac{1}{s} \int dY \int \frac{dt}{t} f_i(x_a) f_j(x_b) \, \frac{d\sigma}{dt}_{| ij \to k, \gamma}. \tag{D.3}
$$

Eq. (4.4) gives us

$$
\frac{dt}{dy} = \frac{M^2}{2 \cosh^2 y}, \tag{D.4}
$$

so that the invariant mass spectrum can be written as

$$
\frac{d\sigma}{dM}_{| pp \to \gamma + \text{jet}} = \frac{M^3}{s} \int dY \int \frac{dy}{\cosh^2 y} f_i(x_a) f_j(x_b) \, \frac{d\sigma}{dt}_{| ij \to k, \gamma}. \tag{D.5}
$$

Since we set cuts in the photon and jet rapidities

$$
|y_1| = |y + Y| < y_{\text{max}} = 2.4, \quad |y_2| = |y - Y| < y_{\text{max}} = 2.4, \tag{D.6}
$$

Eq. (D.5) can be expressed with integral limits as Eq. (4.1).
Bibliography

[1] W. M. Yao et al. [Particle Data Group], J. Phys. G 33, 1 (2006).

[2] M. B. Green, J. H. Schwarz and E. Witten, *Superstring Theory Vol. 1: Introduction*, (Uk: University Press, Cambridge, 1987).

[3] M. B. Green, J. H. Schwarz and E. Witten, *Superstring Theory. Vol. 2: Loop Amplitudes, Anomalies And Phenomenology*, (Uk: University Press, Cambridge 1987).

[4] T. Kaluza, Sitzungsber. Preuss. Akad. Wiss. Berlin (Math. Phys.) 1921, 966 (1921).

[5] O. Klein, Z. Phys. 37, 895 (1926) [Surveys High Energ. Phys. 5, 241 (1986)].

[6] N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, Phys. Lett. B 429, 263 (1998) [arXiv:hep-ph/9803315].

[7] L. Randall and R. Sundrum, Phys. Rev. Lett. 83, 3370 (1999) [arXiv:hep-ph/9905221].

[8] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, Phys. Lett. B 436, 257 (1998) [arXiv:hep-ph/9804398].

[9] D. Lust, S. Stieberger and T. R. Taylor, arXiv:0807.3333 [hep-th].

[10] R. Blumenhagen, B. Kors, D. Lüst and S. Stieberger, Phys. Rept. 445, 1 (2007) [arXiv:hep-th/0610327].

[11] E. Komatsu et al. [WMAP Collaboration], arXiv:0803.0547 [astro-ph].

[12] K. A. Olive, G. Steigman and T. P. Walker, Phys. Rept. 333, 389 (2000) [arXiv:astro-ph/9905320].

[13] S. Cole et al. [The 2dFGRS Collaboration], Mon. Not. Roy. Astron. Soc. 362, 505 (2005) [arXiv:astro-ph/0501174].

[14] M. Tegmark et al., Phys. Rev. D 74, 123507 (2006) [arXiv:astro-ph/0608632].
[15] A. G. Riess et al. [Supernova Search Team Collaboration], Astron. J. 116, 1009 (1998) [arXiv:astro-ph/9805201].

[16] S. Perlmutter et al. [Supernova Cosmology Project Collaboration], Astrophys. J. 517, 565 (1999) [arXiv:astro-ph/9812133].

[17] N. A. Bahcall, J. P. Ostriker, S. Perlmutter and P. J. Steinhardt, Science 284, 1481 (1999) [arXiv:astro-ph/9906463].

[18] M. E. Peskin, J. Phys. Soc. Jap. 76, 111017 (2007) [arXiv:0707.1536 [hep-ph]].

[19] B. Ratra and P. J. E. Peebles, Phys. Rev. D 37, 3406 (1988).

[20] I. Zlatev, L. M. Wang and P. J. Steinhardt, Phys. Rev. Lett. 82, 896 (1999) [arXiv:astro-ph/9807002].

[21] P. G. Ferreira and M. Joyce, Phys. Rev. Lett. 79, 4740 (1997) [arXiv:astro-ph/9707286].

[22] P. G. Ferreira and M. Joyce, Phys. Rev. D 58, 023503 (1998) [arXiv:astro-ph/9711102].

[23] E. J. Copeland, A. R. Liddle and D. Wands, Phys. Rev. D 57, 4686 (1998) [arXiv:gr-qc/9711068].

[24] P. J. E. Peebles and A. Vilenkin, Phys. Rev. D 59, 063505 (1999) [arXiv:astro-ph/9810509].

[25] L. Parker and A. Raval, Phys. Rev. D 60, 063512 (1999) [Erratum-ibid. D 67, 029901 (2003)] [arXiv:gr-qc/9905031].

[26] L. Parker and A. Raval, Phys. Rev. D 60, 123502 (1999) [Erratum-ibid. D 67, 029902 (2003)] [arXiv:gr-qc/9908013].

[27] L. Parker and D. A. T. Vanzella, Phys. Rev. D 69, 104009 (2004) [arXiv:gr-qc/0312108].

[28] R. R. Caldwell, W. Komp, L. Parker and D. A. T. Vanzella, Phys. Rev. D 73, 023513 (2006) [arXiv:astro-ph/0507622].

[29] M. Gasperini, F. Piazza and G. Veneziano, Phys. Rev. D 65, 023508 (2002) [arXiv:gr-qc/0108016].

[30] A. Salam and E. Sezgin, Phys. Lett. B 147, 47 (1984).

[31] S. Weinberg, Phys. Rev. Lett. 59, 2607 (1987).
[32] R. Bousso and J. Polchinski, JHEP 0006, 006 (2000) [arXiv:hep-th/0004134].

[33] L. Susskind arXiv:hep-th/0302219.

[34] M. R. Douglas, JHEP 0305, 046 (2003) [arXiv:hep-th/0303194].

[35] M. R. Douglas and S. Kachru, arXiv:hep-th/0610102.

[36] J. M. Maldacena and C. Nuñez, Int. J. Mod. Phys. A 16, 822 (2001) [arXiv:hep-th/0007018].

[37] G. W. Gibbons, “Aspects of Supergravity Theories,” lectures given at GIFT Seminar on Theoretical Physics, San Feliu de Guixols, Spain, 1984. Print-85-0061 (CAMBRIDGE), published in GIFT Seminar 1984:0123.

[38] G. W. Gibbons and C. M. Hull, arXiv:hep-th/0111072.

[39] P. K. Townsend and M. N. R. Wohlfarth, Phys. Rev. Lett. 91, 061302 (2003) [arXiv:hep-th/0303097].

[40] S. B. Giddings, S. Kachru and J. Polchinski, Phys. Rev. D 66, 106006 (2002) [arXiv:hep-th/0105097].

[41] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, Phys. Rev. D 68, 046005 (2003) [arXiv:hep-th/0301240].

[42] M. Cvetic, G. W. Gibbons and C. N. Pope, Nucl. Phys. B 677, 164 (2004) [arXiv:hep-th/0308026].

[43] J. J. Halliwell, Nucl. Phys. B 286, 729 (1987).

[44] Y. Aghababaie, C. P. Burgess, S. L. Parameswaran and F. Quevedo, JHEP 0303, 032 (2003) [arXiv:hep-th/0212091].

[45] Y. Aghababaie, C. P. Burgess, S. L. Parameswaran and F. Quevedo, Nucl. Phys. B 680, 389 (2004) [arXiv:hep-th/0304256].

[46] G. W. Gibbons, R. Guven and C. N. Pope, Phys. Lett. B 595, 498 (2004) [arXiv:hep-th/0307238].

[47] Y. Aghababaie et al., JHEP 0309, 037 (2003) [arXiv:hep-th/0308064].

[48] R. Bean, S. H. Hansen and A. Melchiorri, Nucl. Phys. Proc. Suppl. 110, 167 (2002) [arXiv:astro-ph/0201127].
[49] M. Tegmark et al. [SDSS Collaboration], Phys. Rev. D 69, 103501 (2004) [arXiv:astro-ph/0310723].

[50] D. Comelli, M. Pietroni and A. Riotto, Phys. Lett. B 571, 115 (2003) [arXiv:hep-ph/0302080].

[51] U. Franca and R. Rosenfeld, Phys. Rev. D 69, 063517 (2004) [arXiv:astro-ph/0308149].

[52] R. M. Wald, “General Relativity,” (University of Chicago Press, Chicago, 1984).

[53] J. Vinet and J. M. Cline, Phys. Rev. D 71, 064011 (2005) [arXiv:hep-th/0501098].

[54] J. Preskill, M. B. Wise and F. Wilczek, Phys. Lett. B 120, 127 (1983).

[55] M. B. Hoffman, arXiv:astro-ph/0307350.

[56] L. Anchordoqui and H. Goldberg, Phys. Rev. D 68, 083513 (2003) [arXiv:hep-ph/0306084].

[57] U. J. Lopes Franca and R. Rosenfeld, JHEP 0210, 015 (2002) [arXiv:astro-ph/0206194].

[58] P. J. Steinhardt, L. M. Wang and I. Zlatev, Phys. Rev. D 59, 123504 (1999) [arXiv:astro-ph/9812313].

[59] L. Anchordoqui, H. Goldberg, S. Nawata and C. Nuñez, Phys. Rev. D 76, 126005 (2007) [arXiv:0704.0928 [hep-ph]].

[60] S. M. Carroll, Phys. Rev. Lett. 81, 3067 (1998) [arXiv:astro-ph/9806099].

[61] J. Polchinski, String Theory, Cambridge University Press (1998)

[62] J. D. Lykken, Phys. Rev. D 54, 3693 (1996) [arXiv:hep-th/9603133].

[63] J. Dai, R. G. Leigh and J. Polchinski, Mod. Phys. Lett. A 4, 2073 (1989).

[64] J. E. Paton and H. M. Chan, Nucl. Phys. B 10, 516 (1969).

[65] D. Berenstein and S. Pinansky, Phys. Rev. D 75, 095009 (2007) [arXiv:hep-th/0610104].

[66] I. Antoniadis, E. Kiritsis and T. N. Tomaras, Phys. Lett. B 486, 186 (2000) [arXiv:hep-ph/0004214].
[67] R. Blumenhagen, B. Kors, D. Lüst and T. Ott, Nucl. Phys. B 616, 3 (2001) [arXiv:hep-th/0107138].

[68] See e.g., D. M. Ghilencea, L. E. Ibanez, N. Irges and F. Quevedo, JHEP 0208, 016 (2002) [arXiv:hep-ph/0205083].

[69] P. Burikham, T. Figy and T. Han, Phys. Rev. D 71, 016005 (2005) [Erratum-ibid. D 71, 019905 (2005)] [arXiv:hep-ph/0411094].

[70] K. Cheung and Y. F. Liu, Phys. Rev. D 72, 015010 (2005) [arXiv:hep-ph/0505241].

[71] G. Domokos and S. Kovesi-Domokos, Phys. Rev. Lett. 82, 1366 (1999) [arXiv:hep-ph/9812260].

[72] S. Cullen, M. Perelstein and M. E. Peskin, Phys. Rev. D 62, 055012 (2000) [arXiv:hep-ph/0001166].

[73] S. J. Parke and T. R. Taylor, Phys. Rev. Lett. 56, 2459 (1986).

[74] S. Stieberger and T. R. Taylor, Phys. Rev. D 74, 126007 (2006) [arXiv:hep-th/0609175].

[75] S. Stieberger and T. R. Taylor, Phys. Rev. Lett. 97, 211601 (2006) [arXiv:hep-th/0607184].

[76] M. L. Mangano and S. J. Parke, Phys. Rept. 200, 301 (1991) [arXiv:hep-th/0509223].

[77] L. J. Dixon, arXiv:hep-ph/9601359.

[78] G. Veneziano, Nuovo Cim. A 57, 190 (1968).

[79] G. Veneziano, Phys. Rept. 9, 199 (1974).

[80] T. van Ritbergen, A. N. Schellekens and J. A. M. Vermaseren, Int. J. Mod. Phys. A 14, 41 (1999) [arXiv:hep-ph/9802376].

[81] L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Phys. Rev. Lett 100, 171603 (2008) [arXiv:0712.0386 [hep-ph]].

[82] L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Phys. Rev. D 78, 016005 (2008) [arXiv:0804.2013 [hep-ph]].

[83] L. A. Anchordoqui, H. Goldberg and T. R. Taylor, arXiv:0806.3420 [hep-ph].
[84] S. Dimopoulos and G. L. Landsberg, Phys. Rev. Lett. 87, 161602 (2001) [arXiv:hep-ph/0106295].

[85] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. Nadolsky and W. K. Tung, JHEP 0207, 012 (2002) [arXiv:hep-ph/0201195].

[86] D. Stump, J. Huston, J. Pumplin, W. K. Tung, H. L. Lai, S. Kuhlmann and J. F. Owens, JHEP 0310, 046 (2003) [arXiv:hep-ph/0303013].

[87] G. L. Bayatian et al. [CMS Collaboration], J. Phys. G 34 995 (2007).

[88] W. W. Armstrong et al. [ATLAS Collaboration], CERN/LHCC 94-43.

[89] L. A. Anchordoqui, J. L. Feng, H. Goldberg and A. D. Shapere, Phys. Rev. D 65 (2002) 124027 [arXiv:hep-ph/0112247].

[90] D. V. Bandurin and N. B. Skachkov, Eur. Phys. J. C 37, 185 (2004).

[91] P. Gupta, B. C. Choudhary, S. Chatterji, S. Bhattacharya and R. K. Shivpuri, arXiv:0705.2740 [hep-ex].

[92] G. Conesa, H. Delagrange, J. Diaz, Y. V. Kharlov and Y. Schutz, Nucl. Instrum. Meth. A 585, 28 (2008) [arXiv:0711.2431 [physics.data-an]].

[93] A. Chamblin and G. C. Nayak, Phys. Rev. D 66, 091901 (2002) [arXiv:hep-ph/0206060].

[94] A. Dainese, J. Phys. G 35, 044046 (2008) [arXiv:0710.3052 [nucl-ex]].

[95] T. Banks and W. Fischler, arXiv:hep-th/9906038.

[96] S. B. Giddings and S. D. Thomas, Phys. Rev. D 65, 056010 (2002) [arXiv:hep-ph/0106219].

[97] E. Eichten, I. Hinchliffe, K. D. Lane and C. Quigg, Rev. Mod. Phys. 56, 579 (1984) [Addendum-ibid. 58, 1065 (1986)].

[98] L. A. Anchordoqui, in Proceedings of the 16th International Conference on Supersymmetry and the Unification of Fundamental Interactions (SUSY08), Seoul, Korea, June (2008), arXiv:0806.3782 [hep-ph].

[99] D. M. Eardley and S. B. Giddings, Phys. Rev. D 66, 044011 (2002) [arXiv:gr-qc/0201034].
[100] H. Yoshino and Y. Nambu, Phys. Rev. D 67, 024009 (2003) [arXiv:gr-qc/0209003].

[101] K. S. Thorne, in Magic Without Magic: John Archibald Wheeler, edited by J. Klauder (Freeman, San Francisco, 1972) p.231.

[102] S. W. Hawking, Nature 248, 30 (1974).

[103] S. W. Hawking, Commun. Math. Phys. 43, 199 (1975) [Erratum-ibid. 46, 206 (1976)].

[104] L. Parker, Phys. Rev. D 12, 1519 (1975).

[105] R. M. Wald, Commun. Math. Phys. 45, 9 (1975).

[106] S. W. Hawking, Phys. Rev. D 14, 2460 (1976).

[107] T. Han, G. D. Kribs and B. McElrath, Phys. Rev. Lett. 90, 031601 (2003) [arXiv:hep-ph/0207003].

[108] R. C. Myers and M. J. Perry, Annals Phys. 172, 304 (1986).

[109] D. N. Page, Phys. Rev. D 13, 198 (1976).

[110] P. Kanti and J. March-Russell, Phys. Rev. D 66, 024023 (2002) [arXiv:hep-ph/0203223].

[111] D. Ida, K. y. Oda and S. C. Park, Phys. Rev. D 67, 064025 (2003) [Erratum-ibid. D 69, 049901 (2004)] [arXiv:hep-th/0212108].

[112] P. Kanti and J. March-Russell, Phys. Rev. D 67, 104019 (2003) [arXiv:hep-ph/0212199].

[113] C. M. Harris and P. Kanti, JHEP 0310, 014 (2003) [arXiv:hep-ph/0309054].

[114] R. Emparan, G. T. Horowitz and R. C. Myers, Phys. Rev. Lett. 85, 499 (2000) [arXiv:hep-th/0003118].

[115] V. Cardoso, M. Cavaglia and L. Gualtieri, Phys. Rev. Lett. 96, 071301 (2006) [Erratum-ibid. 96, 219902 (2006)] [arXiv:hep-th/0512002].

[116] V. Cardoso, M. Cavaglia and L. Gualtieri, JHEP 0602, 021 (2006) [arXiv:hep-th/0512116].

[117] L. A. Anchordoqui, J. L. Feng, H. Goldberg and A. D. Shapere, Phys. Lett. B 594, 363 (2004) [arXiv:hep-ph/0311365].
[118] P. Meade and L. Randall, JHEP 0805, 003 (2008) [arXiv:0708.3017 [hep-ph]].

[119] M. Cardaci et al., “CMS Search Plans and Sensitivity to New Physics using Dijets,” CMS AN-2007/039.

[120] L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Nawata, S. Stieberger, and T. R. Taylor, 
*Dijet signals of low mass strings at the LHC*, arXiv:0808.0497 [hep-ph].
Curriculum Vitae

Title of Dissertation

Manifestations of String Theory in Astrophysical Data and at the LHC

Full Name

Satoshi Nawata

Place and Date of Birth

Tokyo, Japan February 8, 1982

Colleges and Universities, Years attended and degrees

Tokyo Institute of Technology 2002-2004, BA

University of Wisconsin–Milwaukee 2005-2008, Ph.D.

Publications

• L. Anchordoqui, H. Goldberg, S. Nawata and C. Nuñez, Cosmology from String Theory, Phys. Rev. D 76, 126005 (2007) [arXiv:0704.0928 [hep-ph]].

• L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Jet signals for low mass strings at the Large Hadron Collider, Phys. Rev. Lett. 100, 171603 (2008) [arXiv:0712.0386 [hep-ph]].

• L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Direct photons as probes of low mass strings at the CERN LHC, Phys. Rev D 78, 016005 (2008) [arXiv:0804.2013 [hep-ph]].
• L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Nawata, S. Stieberger, and T. R. Taylor,
  *Dijet signals of low mass strings at the LHC*, arXiv:0808.0497 [hep-ph].

Awards and Fellowships

  2008 Institute for Advanced Study, School of Natural Sciences, Princeton, NJ
    Attendance Fellowship: “Strings and Phenomenology”
  2008 Kavli Institute for Theoretical Physics, University of California, Santa Barbara, CA
    Attendance Fellowship: “Gauge Theory and Langlands Duality”
  2008 UWM Chancellor’s Graduate Student Fellowship
  2007 Institute for Advanced Study, School of Mathematics, Princeton, NJ
    Attendance Fellowship: “Workshop on Gauge Theory and Representation Theory”
  2007 UWM Chancellor’s Graduate Student Fellowship
  2006 UWM Chancellor’s Graduate Student Fellowship
  2005 UWM Chancellor’s Graduate Student Fellowship

Major Department

Physics

______________________________
Professor Luis Anchordoqui           Date