Infrared modifications of gravity have been proposed over the years as potentially useful phenomenological models. Massive (multi)-gravity is an interesting class of such theories. The true nature and the ultimate consistency of many of these models remains, however, unclear as one is frequently faced with important problems, e.g. classical instabilities and strong coupling problems which are hard to resolve without a known UV completion of the theory. In this note, we review a recent attempt to study some of these well known problems with the use of the AdS/CFT correspondence. In this context, products of large-$N$ conformal field theories coupled by multi-trace interactions in diverse dimensions are used to define quantum multi-gravity (multi-string/M theory) on a union of (asymptotically) AdS spaces. One-loop effects generate a small $O(1/N)$ mass for the gravitons and provide non-trivial examples of massive multi-graviton theories.

1 Introduction

One of the most vexing problems in theoretical physics is the cosmological constant problem – currently we observe that the universe accelerates at a rate compatible with a vacuum energy density approximately equal to $(10^{-27} \, M_P)^4$. This is a fundamental problem because it appears to challenge our ideas about the laws of nature simultaneously at long and short distance scales.

A popular approach to this problem is based on the idea that gravity should be modified in the infrared (IR) in an attempt to effectively screen the gravitational interaction. Since the Einstein term in gravity is a two-derivative term, any local IR modification of gravity proceeds via a potential for the graviton. A quadratic mass term is the leading non-trivial part of the potential, but higher order terms can be present. Giving a mass to the graviton is a possibility that was explored many years ago by Fierz and Pauli [1]. Later it was realized, following the important work of S. Deser and collaborators, that gravitational theories with a mass term are behaving quite unlike other theories and are typically accompanied by several problems, most notably ghost/tachyon instabilities and strong coupling issues.

At the same time, discussions of massive gravity in the context of the cosmological constant problem have been inconclusive. Using the cosmological equations for massive gravity [2] one finds that there is an effective cosmological acceleration at late times, which is equivalent to a constant $positive$ vacuum energy $\Lambda^4 \sim m_g^2 M_P^2$, [3]. If the mass of the graviton is of order the inverse Hubble scale today, $m_g \sim H_0^{-1}$, then $\Lambda \sim 10^{-3} \text{eV}$ in agreement with today's observations. This result is encouraging, but it turns out that higher terms in the graviton potential give more and more dominant contributions to the late time evolution of the universe [3] indicating that we need a much better understanding of the theory.

Our purpose in this short review of Ref. [4] is to discuss the fundamental issues of massive graviton theories in a context where an alternative description of the theory provides us with a much needed additional control. Before going into the details of this context let us briefly summarize the main features of the fundamental issues that we want to understand better.

It has been shown in [5–8], and more recently in [9–11], that massive graviton theories are generically unstable. The unstable mode (usually called a Boulware-Deser mode) is a ghost (i.e. a mode with a wrong sign kinetic term), and sometimes also a tachyon (a mode with a negative mass squared). The source of the...
instability lies in the non-linear structure of the kinetic terms. Although it is possible sometimes to fine-tune the instabilities away in a fixed background, small departures from the original background are enough to reinstate them. Even in cases where one considers non-Lorentz invariant potentials as in [12], and other cases reviewed in [13], the theory is stable only in a subclass of backgrounds. This behavior persists also in more exotic cases where the massive graviton is a resonance (e.g. in brane-induced gravity [14]).

Another peculiar feature of massive gravity is the appearance of a strong coupling scale at energies that are hierarchically smaller than the Planck scale. To be precise, if the graviton mass \( m_g \ll M_P \), then the theory becomes strongly coupled at energies comparable or larger than \( \Lambda_V = (m_g^3 M_P)^{1/5} \) [15], a scale that is hierarchically smaller than the Planck scale. The interactions that become strongest are those of the scalar mode of the graviton. By judiciously modifying the graviton potential [16] the scale where gravity becomes strong can be improved to \( \Lambda_{AGS} = (m_g^2 M_P)^{1/3} \).

Thinking of the massive graviton actions as low-energy effective actions, the above scales should be viewed as UV cutoffs in the standard sense of quantum field theory. Beyond the energy of these scales one must know the UV completion of the theory in order to proceed. For massive graviton theories this UV completion is essentially unknown and this poses a major hurdle in making further progress. This peculiar IR behavior of massive gravity is reflecting an important feature: the failure of IR physics to decouple in a simple way from UV physics. The theory has two characteristic scales, \( m_g \) and \( M_P \), however, there is a dynamically generated intermediate scale where the effective low-energy description breaks down.

We should point out that neither \( \Lambda_V \) nor the improved \( \Lambda_{AGS} \) are cutoffs compatible with today’s experimental results on the gravitational force. Only a cutoff that would scale as \( \Lambda_\ast \sim \sqrt{m_g M_P} \) would be marginally compatible with experiment, but a Lorentz-invariant massive graviton theory with such a cutoff is currently unknown (nevertheless some examples in AdS backgrounds were presented in [4]).

Finally, it will be useful to introduce a related concept, that of multi-graviton theories. To formulate a massive gravity theory, or to write down a general potential for the graviton, a fiducial metric must be introduced. This metric is non-dynamical and is typically taken to be flat (constant curvature metrics are also considered). A natural generalization involves making the fiducial metric dynamical. In that case, we are dealing with a bi-gravity theory\(^1\)– a theory with two propagating gravitons, two in general distinct Planck scales and in the decoupling limit two diffeomorphism invariances. By adding a potential that couples the two gravitons together it is possible to break one of the diffeomorphism invariances and give mass to a linear combination of the two gravitons. Bi-gravity theories naturally generalize to multi-gravity theories involving more than two metrics. We will soon encounter such theories in the context string theory.

## 2 Conformal Field Theories and Multi-gravity in AdS spaces

In the presence of rather low UV cutoffs, it is sensible to seek out theories of modified gravity which are UV complete. The AdS/CFT correspondence provides a rather rare context where a theory of gravity (in asymptotically AdS spaces) has an implicit UV completion. The completion is provided by a lower dimensional quantum field theory without gravity living on the boundary of the AdS space. Examples where the gravity on the AdS side is massive are an interesting arena where the problems of massive gravity outlined above can be discussed from a totally new and different perspective: the perspective of the non-gravitational field theory on the boundary.

A natural set of AdS/CFT examples with massive gravity can be obtained following a rather simple idea. In standard AdS/CFT pairs the graviton in AdS maps to the stress-energy tensor of the conformal field theory (CFT). Accordingly, diffeomorphism invariance in AdS maps to energy-momentum conservation in the CFT. Deforming the CFT in a way that violates the energy-momentum conservation corresponds in AdS to a gravitational Higgs effect where diffeomorphism invariance gets broken and the graviton acquires a mass.

\(^1\) Bi-gravity was first introduced in [17, 18] in the context of the strong interactions. For a recent detailed classical study of such theories and their solutions see [19–21].
There are several ways to violate the stress-energy conservation on the boundary. One of them involves the introduction of $(d-1)$-dimensional defects [22]. Another involves product conformal field theories deformed by multi-trace deformations [23, 24]. In what follows we will discuss massive (multi)-gravity in the second setup.

In order to be concrete, we will consider the following instructive prototype (several generalizations can be found in [4]). Our setup consists of two $d$-dimensional large-$N$ CFTs (CFT$_1$ and CFT$_2$) which are coupled by a double-trace interaction $gO_1O_2$, where $O_{1,2}$ are single-trace (scalar) operators in CFT$_{1,2}$ respectively. The two CFTs are generic – we will only demand that they are defined on the same space and that each of them has a dual description in terms of a gravity (string/M) theory on a space of the form $AdS_{d+1} \times M_i$, where $M_i$ ($i = 1, 2$) is some compact space. Coupling the CFTs by a multiple trace deformation is the generic way to make the CFTs communicate without spoiling their individual gauge symmetries. In order to preserve the standard large-$N$ counting the double-trace coupling constant $g$ must scale with $N$ as $O(N^0)$.

Without the double-trace interaction the field theory on the boundary is the product of two decoupled CFTs. This theory maps to a product of two string/M theories on a product manifold of the form $\prod_i AdS_{d+1} \times M_i$ [23, 24]. In cases, where the (super)gravity approximation is valid this theory reduces to a bi-(super)gravity theory. This is a trivial bi-gravity theory in the sense that there are two gravitons that do not communicate with each other.

The presence of the double-trace coupling $g$ on the boundary makes the dual bi-gravity theory a far more interesting system. The two CFTs are now coupled non-trivially. The total stress-energy tensor

$$T^\text{tot}_{\mu\nu} = T^{(1)}_{\mu\nu} + T^{(2)}_{\mu\nu} - g\eta_{\mu\nu}O_1O_2$$

remains conserved. An orthogonal combination of the stress-energy tensors, however, is no longer conserved. Roughly speaking (for a more precise expression see [24])

$$\partial^\mu(T^{(1)}_{\mu\nu} - T^{(2)}_{\mu\nu}) \sim g[(\partial_\nu O_1)O_2 - O_1(\partial_\nu O_2)].$$

In the bulk product of AdS spaces a corresponding breaking of a subgroup of the diffeomorphism symmetries occurs. Simultaneously, a linear combination of the gravitons becomes massive. Field theory predicts [23, 24] a mass determined by the scaling dimension of the non-conserved stress-energy tensor

$$m_g^2 = d(\Delta_T - d) \sim \frac{g^2}{N^2}.$$  

In gravity this mass comes about in the following way.

The effect of multiple trace deformations in the AdS/CFT correspondence was analyzed at tree-level (i.e. to leading order in the $1/N$ expansion) some time ago in Refs. [25–29]. In situations where we can trust the supergravity approximation, the main effect of field theory multi-trace deformations in the bulk is a change of the standard Dirichlet or Neumann boundary conditions to mixed boundary conditions for the (scalar) fields $\phi_i$ that are dual to the operators $O_i$ participating in the deformation.

References [23, 24] (see [30, 31] for related earlier work) observed that the mixed boundary conditions facilitate a scalar loop correction to the graviton propagator in the bulk of the form depicted in Fig.1. This correction induces a non-zero mass term for a linear combination of the gravitons. The mass is of order $g/N$ as anticipated from the field theory on the boundary. A detailed computation of this mass can be found in [24] (a discussion of the graviton mass matrix in a more general context of a network of interacting CFTs can be found in appendix D of [4]).

Further properties of this system can be read off the dual gauge theory. The double-trace deformation can be relevant, marginal or irrelevant depending on the scaling dimensions $\Delta_1, \Delta_2$ of the operators $O_1, O_2$. In general, the double-trace deformation breaks the original conformal invariance and the couplings of the gauge theory run in the standard renormalization group (RG) sense. We would like to know the possible endpoints of these flows.
\textbf{Fig. 1} A diagrammatic representation of the one-loop amplitude that gives a mass to a linear combination of two gravitons, each living on a different AdS space. The mixed boundary conditions for the dual scalars $\phi_1, \phi_2$ give a non-vanishing 1-2 graviton propagator.

To uncover the full structure of the RG running in the space of multi-trace couplings consider the general deformation

$$\delta S = \int d^d x \left[ N \sum_{i=1,2} g_i O_i + \sum_{i,j=1,2} g_{ij} O_i O_j + \frac{1}{N} \sum_{i,j,k=1,2} g_{ijk} O_i O_j O_k + \cdots \right], \quad (4)$$

where the dots indicate higher multi-trace interactions. It will be convenient to denote collectively the single-trace, double-trace, etc couplings as

$$g^{(1)} = \{g_1, g_2\}, \quad g^{(2)} = \{g_{11}, g_{12}, g_{22} \}, \cdots \quad (5)$$

We will assume that the scaling dimensions $\Delta_1, \Delta_2 < \frac{d}{2}$ so that the double-trace couplings are relevant.

We would like to determine the RG flows and the fixed points associated to the one-loop $\beta$-functions

$$\beta_i = \frac{\partial g_i}{\partial t}, \quad \beta_{ij} = \frac{\partial g_{ij}}{\partial t}, \quad \beta_{ijk} = \frac{\partial g_{ijk}}{\partial t}, \cdots \quad (6)$$

where $t$ is the logarithm of the RG scale.

Performing an expansion in conformal perturbation theory up to second order in the couplings one finds the following type of equations

$$\dot{g}^{(1)} = (d - \Delta^{(1)}) g^{(1)} + C^{(1)(1)} g_1^2 + C^{(1)(2)} g_1 g_2 + \cdots \quad (7a)$$

$$\dot{g}^{(2)} = (d - \Delta^{(2)}) g^{(2)} + C^{(2)(1)} g_1^2 + C^{(2)(2)} g_1 g_2 + C^{(2)(1)} g_2^2 + C^{(2)(2)} g_1 g_2 + \cdots \quad (7b)$$

$$\dot{g}^{(3)} = (d - \Delta^{(3)}) g^{(3)} + C^{(3)(1)} g_1^3 + C^{(3)(2)} g_1 g_2^2 + C^{(3)(1)} g_1 g_2 + \cdots \quad (7c)$$
circle of the double-trace coupling space where individually deformed to a new fixed point with their leading contribution in the for the late RG time behavior of the system. Either it is driven towards the most attractive fixed point from them, there are always directions where the RG flow is repelling. There are two distinct possibilities as is apparent from the right plot in Fig. 2, these points are repellors of the RG flow. When we perturb away from them, there are always directions where the RG flow is repelling. There are two distinct possibilities for the late RG time behavior of the system. Either it is driven towards the most attractive fixed point A,

$$\Delta_1$$

We included terms up to four-trace couplings in the above expressions hoping that the general structure is evident. $C_{(m)(n)(\rho)}$ is shorthand for a set of 3-point function coefficients normalized in such a way that their leading contribution in the $1/N$ expansion is of order $N^0$. More details can be found in [4]. In what follows we would like to bring forward some of the key properties implied by these equations.

First, to leading order in $1/N$ we can consistently set all the single-trace couplings $g(1)$ to zero at all scales. In search of fixed points, we observe that $\beta(2)$ depend only $g(2)$ and $\beta(i) = 0, i \geq 3$, determine straightforwardly, in our order of approximation, the fixed point values of $g(2)$, $i \geq 3$, once the fixed point values of the double-trace couplings $g(2)$ is computed. Hence, it is enough to solve the double-trace fixed point equations. With appropriate normalization these equations, which are exact in the large-$N$ limit, read

\begin{align}
(d - 2\Delta_1)g_{11} - 8g_{11}^2 - 8g_{12}^2 = 0, \quad (d - 2\Delta_2)g_{22} - 8g_{22}^2 - 8g_{12}^2 = 0, \\
(d - \Delta_1 - \Delta_2)g_{12} - 8g_{12}(g_{11} + g_{22}) = 0.
\end{align}

(8a)

(8b)

This is a simple set of quadratic equations with the following solutions.

Non-interacting fixed points (i.e. points with $g_{12} = 0$) exist when

$$g_{11}(8g_{11} - (d - 2\Delta_1)) = 0, \quad g_{22}(8g_{22} - (d - 2\Delta_2)) = 0.$$  

(9)

Assuming $\Delta_1, \Delta_2 \neq d/2$ these are points where the CFTs 1 and 2 are either undeformed ($g_{ii} = 0$) or individually deformed to a new fixed point with $g_{ii} = \frac{1}{2}(d - 2\Delta_i)$. These theories are well known [25].

Interacting fixed points (i.e. fixed points with $g_{12} \neq 0$) exist only when $\Delta_1 = \Delta_2$ or $\Delta_1 + \Delta_2 = d$ with $\Delta_1, \Delta_2 \neq d/2$ in both cases. These points lie on a circle in the $(g_{11}, g_{22}, g_{12})$ space. In a parametrization of the double-trace coupling space where $\Delta_1 + \Delta_2 = d$ at $g_{11} = g_{22} = g_{12} = 0$ this circle - denoted as circle C in the left figure in Fig. 2 - is described by the equations

$$g_{11} = -g_{22} \equiv f, \quad g_{12} \equiv g \quad a f^2 - af + 4f^2 = 0, \quad 2a \equiv d - 2\Delta_1 = 2\Delta_2 - d.$$  

(10)

As is apparent from the right plot in Fig. 2 these points are repellors of the RG flow. When we perturb away from them, there are always directions where the RG flow is repelling. There are two distinct possibilities for the late RG time behavior of the system. Either it is driven towards the most attractive fixed point A,
where $g_{12} = 0$ and the interaction between the CFTs is lost, or it is driven towards stronger and stronger coupling where the $1/N$ expansion eventually breaks down.

Incorporating the $1/N$ corrections does not modify this picture dramatically. In perturbation theory, these corrections will slightly shift the fixed point values determined at leading order. An important new element should be noted however. Already at the next-to-leading order there is a running of the single-trace couplings which is sourced by the evolution of the double-trace couplings.

3 Lessons for Massive Gravity

We can now ask what are the consequences of these properties in gauge theory for the dual massive multi-gravity theory on the AdS side of the correspondence. More importantly, what lessons can we extract for modified theories of gravity from this context?

In the setup we have just described, the dual gravitational theory is at tree-level (i.e. to leading order in the $1/N$ expansion) a string/M-theory on a union of AdS spaces. In the supergravity regime (whenever this is appropriate), the low-energy dynamics is described by a bi-gravity theory with mixed boundary conditions for a set of scalar fields. At this order, there are two massless non-interacting gravitons. Boundary RG flows do not manifest in the bulk as radial running of the background solution. This is consistent with the fact that in gauge theory there is no running of single-trace couplings at leading order in $1/N$.

The plot thickens beyond tree-level. First, one can argue on general grounds [32] that the multi-trace deformed gauge theory defines a new type of string theory, called non-local string theory (NLST), with non-standard features. This theory is non-local on the worldsheet – one has to extend the perturbative expansion to include contact interactions between disconnected worldsheets. A clear manifestation of the rules of this new perturbative expansion has been given for 2D string theories in [33]. NLSTs are also non-local in spacetime. Typically, the non-locality scale is set by the AdS scale [34].

Whatever the complicated structure of our multi-NLSTs is, we have seen that one-loop effects generate a non-trivial graviton potential. One can imagine, in situations where supergravity can be trusted, that the low-energy dynamics is described by a set of low-energy effective actions in the general spirit of bi-gravity theories of Damour and Kogan [19] (extra matter besides the gravitons has to be included in these actions). The tree-level AdS$_5 \times$ AdS$_5$ background is not a solution of the equations of motion of this action. In agreement with the structure of RG flows in gauge theory, the one-loop effects backreact and deform the tree-level background.

Gauge theory makes certain predictions for the structure of this new bi-gravity action. First, the equations of motion of this action must have an interacting AdS$_5 \times$ AdS$_5$ product solution where some of the scalar fields have a non-trivial profile and gravity is massive. These solutions map to the interacting products of CFTs lying on the fixed point circle $C$ in Fig. 2. Furthermore, there should be domain wall solutions interpolating between a ‘UV’ AdS$_5 \times$ AdS$_5$ geometry with a massive graviton and an ‘IR’ AdS$_5 \times$ AdS$_5$ geometry with two massless non-interacting gravitons. These solutions would correspond to the gauge theory RG flows from the fixed circle $C$ to the attractive point $A$. One could think of them as a dynamical way of switching off the graviton mass. In AdS space there is no vDVZ discontinuity [35,36] which could prevent the existence of such solutions.

It is interesting to ask if we can use the gauge theory to learn anything about the typical problems that plague the massive graviton theories – instabilities and strong coupling issues. A naive calculation of strong coupling scales in massive gravity around AdS (repeating the arguments of [16]) shows [4] that the graviton interactions become strong at a scale set by the AdS scale. This scale appears, however, to be a fake. There are several precedents of such fake scales – Kaluza-Klein gravity and string theory are such examples. In these cases, it is important to consider all the degrees of freedom in the system and the calculation that leads to the above strong coupling scales around the AdS scale does not take into account the interactions of the graviton with the other matter fields in the system.

There are several reasons why one should expect the absence of a strong coupling problem in our massive graviton theories. At the AdS scale, which is the dangerous scale for our naive estimates, new
degrees of freedom appear as bound states of the dual scalars (these are crucial for the Higgs phenomenon in gravity). The AdS scale is also a scale where our multi-string theories exhibit non-locality. At the same time, on the dual gauge theory side there is no indication for an intermediate strong coupling scale.

The usual ghost/tachyon instabilities also appear to be absent in our examples. On the gauge theory side there is no sign of an instability which one would interpret as a Boulware-Deser mode. This would fit nicely with the absence of intermediate strong coupling scales, because the two are, in general, related [9].

What we learned in this short review can be summarized in a few lines as follows. There are UV complete multi-graviton theories with light gravitons – masses of the order $O(1/N)$ – which have a dual description in terms of a deformed product of large-$N$ CFTs. In these theories

- When the diffeomorphism invariance breaking is small, there are no strong coupling problems and no Boulware-Deser instabilities.
- Stable massive graviton backgrounds are not generic, they require fine-tuning.
- Without fine-tuning the spacetime background backreacts either to ‘destroy’ the theory or to erase the effects of the graviton mass and trivialize the bi-gravity interactions. The ‘destruction’ of the theory occurs with a breakdown of the large-$N$ expansion or in 2D examples [33] with a breakdown of the double-scaling limit that defines it.

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References

[1] M. Fierz and W. Pauli, Proc. Roy. Soc. Lond. A 173 211 (1939).
[2] S. V. Babak and L. P. Grishchuk, Int. J. Mod. Phys. D 12 1905 (2003) [arXiv:gr-qc/0209006].
[3] E. Kiritsis, Fortsch. Phys. 52 200 (2004) [Phys. Rept. 421 (2005) ERRAT,429,121-122.2006] 105, [arXiv:hep-th/0310001].
[4] E. Kiritsis and V. Niarchos, doi:10.1016/j.nuclphysb.2008.12.010 [arXiv:0808.3410].
[5] C. Aragone and S. Deser, Nuov. Cim. 3A 709 (1971).
[6] C. Aragone and S. Deser, Phys. Lett. B 86 161 (1979).
[7] D. G. Boulware and S. Deser, Phys. Lett. B 40 227 (1972).
[8] S. Deser and A. Waldron, Nucl. Phys. B 631 369 (2002) [arXiv:hep-th/0112182].
[9] C. Deffayet and J. W. Rombouts, Phys. Rev. D 72, 044003 (2005) [arXiv:gr-qc/0505134].
[10] P. Creminelli, A. Nicolas, M. Papucci and E. Trincherini, JHEP 0509, 003 (2005) [arXiv:hep-th/0505147].
[11] G. Gabadadze and A. Gruzinov, Phys. Rev. D 72, 124007 (2005) [arXiv:hep-th/0312074].
[12] N. Arkani-Hamed, H. C. Cheng, M. A. Luty and S. Mukohyama, JHEP 0405 074 (2004) [arXiv:hep-th/0312099].
[13] V. A. Rubakov and P. G. Tinyakov, [arXiv:0802.4379].
[14] R. Gregory, Prog. Theor. Phys. Suppl. 172, 71-80 (2008) [arXiv:0801.1603].
[15] A. I. Vainshtein, Phys. Lett. B 39 393 (1972).
[16] N. Arkani-Hamed, H. Georgi and M. D. Schwartz, Annals Phys. 305 96 (2003) [arXiv:hep-th/0210184].
[17] C. J. Isham, A. Salam and J. Strathdee, Phys. Rev. D3, 867 (1971).
[18] A. Salam and J. Strathdee, Phys. Rev. D16, 2668 (1977).
[19] T. Damour and I. I. Kogan, Phys. Rev. D 66, 104024 (2002) [arXiv:hep-th/02020642].
[20] T. Damour, I. I. Kogan and A. Papazoglou, Phys. Rev. D 66 104025 (2002) [arXiv:hep-th/02020644].
[21] T. Damour, I. I. Kogan and A. Papazoglou, Phys. Rev. D 67 064009 (2003) [arXiv:hep-th/0212155].
[22] A. Karch and L. Randall, JHEP 0105 008 (2001) [arXiv:hep-th/0011156].
[23] E. Kiritsis, JHEP 0611 049 (2006) 049 [arXiv:hep-th/0608088].
[24] O. Aharony, A. B. Clark and A. Karch, Phys. Rev. D 74 086006 (2006) [arXiv:hep-th/0608089].
[25] E. Witten, [arXiv:hep-th/0112258].
[26] M. Berkooz, A. Sever and A. Shomer, JHEP 0205 034 (2002) [arXiv:hep-th/0112264].
[27] W. Muck, Phys. Lett. B 531 301 (2002) [arXiv:hep-th/0201100].
[28] P. Minces, Phys. Rev. D 68 024027 (2003) [arXiv:hep-th/0201172].
[29] A. C. Petkou, JHEP 0206 009 (2002) [arXiv:hep-th/0201258].
[30] M. Porrati, JHEP 0204 058 (2002) [arXiv:hep-th/0112166].
[31] M. Porrati, Mod. Phys. Lett. A 18 1793 (2003) [arXiv:hep-th/0306253].
[32] O. Aharony, M. Berkooz and E. Silverstein, JHEP 0108 006 (2001) [arXiv:hep-th/0105309].
[33] E. Kiritsis and V. Niarchos, JHEP 0808 044 (2008) [arXiv:0805.4234].
[34] O. Aharony, M. Berkooz and B. Katz, JHEP 0510 097 (2005) [arXiv:hep-th/0504177].
[35] H. van Dam and M. J. G. Veltman, Nucl. Phys. B 22 397 (1970).
[36] M. Porrati, Phys. Lett. B 498 92 (2001) [arXiv:hep-th/0011152].