Physics Letters B 668 (2008) 67–71

Contents lists available at ScienceDirect

Physics Letters B

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A note on quantum aspects of multiple membranes

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ARTICLE INFO

Article history:
Received 4 August 2008
Accepted 10 August 2008
Available online 19 August 2008
Editor: L. Alvarez-Gaumé

ABSTRACT

In this Letter we investigate quantum aspects of the newly proposed theory of multiple membranes put forward by Bagger and Lambert. In particular we analyse the possibility of a finite renormalisation of the coupling at one loop.

1. Introduction

The theory of multiple membranes has been something of a mystery for many years [1]. Now, it appears real progress has been made with a newly proposed theory by Bagger, Lambert and Gustavsson [2–5]. This breakthrough immediately attracted significant attention with a great deal of work exploring numerous avenues [6]. One of the key elements of the Bagger–Lambert theory is the presence of a coupling given by the inverse of the level, k of the Chern–Simons theory. At large k the theory becomes perturbative. Clearly, a tunable dimensionless parameter is not to be expected from a theory of membranes since by definition M-theory does not contain any free parameters. This parameter therefore encodes a property of a particular background. Its interpretation is that the background is formed with a particular modding out by a $\mathbb{Z}_k$ action and is discussed in [7–9] (see also [10]). We will happily accept the presence of a perturbative parameter in the theory and without looking the gift horse in the mouth proceed to use it to examine some quantum aspects of the theory.

An obvious question is whether the beta function vanishes to give quantum consistency. In fact, through work by Kapustin and Pronin [11][1] on properties of Chern–Simons theories coupled to matter this question may be immediately answered and indeed the beta function must vanish. (This has also been addressed directly at one loop by Gustavsson [12].) Here, we will be concerned with the possibility of a finite shift in the level, k at one loop. Pure Chern–Simons theories are known to produce such a shift in the coupling at one loop once a careful regularisation is used (see, e.g., [13,14]), while supersymmetric Chern–Simons theories have also been investigated [15] with the possible one loop shift explored in detail for a variety of different supersymmetries. (The fermion content is crucial since integrating out massive fermions is known to contribute to the shift through their effective action [16].)

One must interpret this carefully especially since the effect seems scheme dependent. We will follow the view espoused in [17] for pure Chern–Simons where the shift could be seen from a careful treatment of the phase of the partition function. In [17], the partition function of pure Chern–Simons was calculated non-perturbatively, where possible, and it was found to be a function of the shifted level. This indicated that although the shift may be derived at one loop, the calculation is picking up that the full non-perturbative result will be a function not of k but of the shifted k. In the present scenario, the most immediate physical effect caused by such a shift will be on the moduli space which depends critically on k [7–9].

2. Bagger–Lambert theory

We now describe the theory of Bagger, Lambert and Gustavsson [2–5] using the conventions of Van Raamsdonk [18]. There are eight scalars X^I, $I = 1, \dots, 8$ valued in SO(4) or equivalently the bifundamental of $SU(2) \times SU(2)$ as follows: $X^I = \frac{1}{2} X^I_a \sigma^a$ where $\sigma^a = (i\sigma^i, 1)$ and σ^i are the Pauli matrices. (In what follows $a, b, c = 1, \dots, 4$ and $i, j, k = 1, \dots, 3$.) There are eight fermions and their conjugates similarly valued in SO(4) and two gauge fields A_μ and $\hat{A}_\mu$ valued in SU(2), i.e., $A_\mu = A^+_{\mu i}\sigma^i$ and $\hat{A}_\mu = A^-_{\mu i}\sigma^i$, where A^+ and A^- are the self-dual and anti-self-dual parts of the SO(4) gauge field respectively. The gauge fields couple to matter through the covariant derivative:

$$D_\mu X^I = \partial_\mu X^I + i A_\mu X^I - i X^I \hat{A}_\mu, \tag{2.1}$$

and the action is given by:

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[1] We would like to thank Costis Papageorgakis for bringing [11] to our attention.

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doi:10.1016/j.physletb.2008.08.021

$$
\begin{aligned}
S = \int d^3x\, \mathrm{tr}\Big[& (D^\mu X^I)^\dagger D_\mu X^I + i\bar{\psi}^\dagger \Gamma^\mu D_\mu \psi \\
& - \frac{1}{3}\cdot\frac{4\pi}{k} i\bar{\psi}^\dagger \Gamma_{IJ}\big(X^I (X^J)^\dagger + X^J \psi^\dagger X^I + \psi (X^I)^\dagger X^J\big) \\
& - \frac{2}{3}\cdot\left(\frac{4\pi}{k}\right)^2 X^{[I} X^{J\dagger} X^{K]} X^{K\dagger} X^J X^{I\dagger} \\
& + \frac{k}{4\pi}\epsilon^{\mu\nu\lambda}\left(A_\mu \partial_\nu A_\lambda + \frac{2}{3} i A_\mu A_\nu A_\lambda\right) \\
& - \frac{k}{4\pi}\epsilon^{\mu\nu\lambda}\left(\hat{A}_\mu \partial_\nu \hat{A}_\lambda + \frac{2}{3} i \hat{A}_\mu \hat{A}_\nu \hat{A}_\lambda\right)\Big]. \qquad (2.2)
\end{aligned}
$$

The first line contains the usual kinetic terms, the second the Yukawa couplings, the third the sextic interaction and final line the two Chern–Simons actions for the vector potentials.

As described in [19] we may also consider a massive deformation of the theory that still preserves all the supersymmetries. This is given by adding the following term to the action[2]:

$$S_{\mathrm{mass}} = \int d^3x \big(-\mu^2\, \mathrm{tr}(X^I X^I) + i\mu\, \mathrm{tr}(\bar{\psi}\Gamma_{3456}\psi)\big). \qquad (2.3)$$

Naively one would expect the mass term for fermions to affect the one loop shift in k [15,16]. However, as we shall see the presence of Γ_{3456} will mean that the mass deformation will actually leave the one loop shift of k invariant. That is, the shift will be independent of the fermion mass deformation that preserves supersymmetry.

3. One loop shift in the level

We begin by reviewing the known arguments for a perturbative shift in the coupling in light of Bagger–Lambert theory and go on to elucidate the effects of the extra structure present. Primarily we follow the arguments in [14,15] where careful treatment of the one loop correction can be found. There is also an excellent description of how this shift arises when being careful with the phase of the path integral and the associated introduction of the eta invariant via the Atiyah Patodi Singer index theorem [17]. When the partition function can be calculated exactly it is then a function of the shifted k. This is similar to various $(1+1)d$ integrable models where a finite shift in parameters makes the WKB approximation exact [17]. This second approach provides insight into the nature of the shift and faith that the one loop correction is something physical, but is not directly extendible to the case of Chern–Simons coupled to matter that is required here.

To start with, we must identify an appropriate parameter with which to do perturbation theory. Fortunately, as mentioned previously, the Bagger–Lambert action does contain such a parameter; it is the Chern–Simons coefficient, k, which we will take to be large in order that there may be a perturbative regime. One might make the objection that since k must be integer-valued in order that gauge invariance is preserved, we should not really use it as a coupling constant since it cannot be continuously varied. However, we will take a leaf out of earlier work on perturbative Chern–Simons theory and brush this subtlety aside. Indeed, as we shall see, an important result is that quantum corrections impose that k can change by addition of an integer, thus preserving large gauge invariance.

Consider, then, one of the Chern–Simons terms in (2.2). To regulate divergences, we introduce a Yang–Mills term $-F^2/(2g_{\mathrm{YM}}^2)$. In three dimensions, g_{YM}^2 has dimensions of mass which makes the gauge fields topologically massive. In order to obtain physical results, however, we will want to take $g_{\mathrm{YM}}^2 \to \infty$ which decouples the Yang–Mills regulator. For this purpose, it is useful to define the dimensional parameter $m = g_{\mathrm{YM}}^2(k/4\pi)$ and consider the limit as $m \to \infty$. In order to perform calculations we must also fix a gauge, which we take to be $\partial_\mu A^\mu = 0$. This we do in the usual way by adding ghosts $(c, \bar{c})$ and the gauge-fixing term $-(\partial_\mu A^\mu)^2/\alpha$ to the action. Furthermore, we choose to work in Landau gauge where $\alpha = 0$, which has the dual advantages of radically simplifying the gluon propagator and taming infrared divergences (see, e.g., [14] and references therein).

Introducing renormalisation functions in the usual way ($Z_{\bar{c}Ac}$ for the ghost–gluon vertex, Z_c for the ghost kinetic term and Z_{a} for the antisymmetric part of the gluon kinetic term) and a Ward identity associated with the ghost–gluon vertex, we can then write the renormalised Chern–Simons coupling as

$$k' = k\frac{Z_{\bar{c}Ac}^2}{Z_{\mathrm{a}} Z_c^2}. \qquad (3.1)$$

So, in order to determine the effect of renormalisation on the Chern–Simons coefficient we need to examine the ghost self-energy, ghost–gluon interaction and the gluon self-energy.

3.1. Gluon self-energy

In a general gauge, the classical A^+ propagator is given by

$$\Delta_{\mu\nu} = -\frac{4\pi i}{k}\frac{m}{p^2(p^2-m^2)}\left\{ im\epsilon_{\mu\nu\lambda}p^\lambda + \eta_{\mu\nu}p^2 - p_\mu p_\nu\left(1 - \frac{\alpha k}{4\pi m}\frac{p^2-m^2}{p^2}\right)\right\}, \qquad (3.2)$$

which, in the $m \to \infty$ limit (and in Landau gauge) reduces simply to

$$\Delta_{\mu\nu} = -\frac{4\pi}{k}\frac{\epsilon_{\mu\nu\lambda}p^\lambda}{p^2} \qquad (3.3)$$

as expected. The propagator for A^- is simply minus this.

For the renormalisation functions here, the symmetries of the theory can be used to separate the gluon self-energy, $\Pi_{\mu\nu}^{(1)}$ (at one loop), into symmetric ($\Pi_{\mathrm{s}}^{(1)}$) and anti-symmetric ($\Pi_{\mathrm{a}}^{(1)}$) parts respectively:

$$\Pi_{\mu\nu}^{(1)} = \Pi_{\mathrm{a}}^{(1)}\epsilon_{\mu\nu\lambda}p^\lambda + \frac{1}{m}\Pi_{\mathrm{s}}^{(1)}\big(\eta_{\mu\nu}p^2 - p_\mu p_\nu\big), \qquad (3.4)$$

where $Z_{\mathrm{a}} = 1 - \Pi_{\mathrm{a}}^{(1)}$ at one loop.

3.1.1. Gluonic contributions

In the case of pure Chern–Simons theory one can only have gluons and ghosts running in the loop via diagrams (a)–(c) in Fig. 1 and this was covered in detail in [14]. Furthermore, the absence of an $A^+ - A^-$ propagator means that these arguments apply separately to the self-dual and anti-self-dual parts of the gauge field respectively. In this respect we can view the levels of the two Chern–Simons theories as being essentially independent and see what effect the renormalisation has on each in turn.

The calculations performed in [14] are therefore unchanged in this scenario and we refer the reader to that paper for the full details. Here we simply present the results: $\Pi_{\mathrm{a}}^{(1)}$ has a leading term proportional to $m/|m| = \mathrm{sgn}(m)$. This can be related to $\mathrm{sgn}(k)$, and when the contribution is evaluated in the $m \to \infty$ limit one obtains[3]

$$\Pi_{\mathrm{a}}^{(1)} = \frac{7}{3k}C_2\, \mathrm{sgn}(k). \qquad (3.5)$$

[2] In fact, one must also add a mass-dependent potential in order for supersymmetry to be preserved, but this will not play any role in our discussions.

[3] Note that in our conventions, $f^{acd}f^{bcd} = C_2\delta^{ab}$ so that $C_2(\mathrm{SU}(N)) = N$.

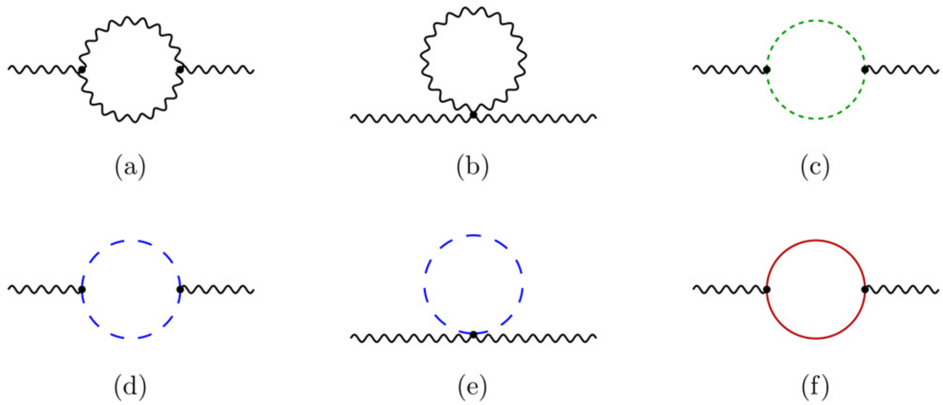

Fig. 1. Processes contributing to the gluon self-energy at one loop with F^2 regularisation. (a) and (b) are corrections from virtual gluons; (c) is due to ghosts; (d) and (e) are the exchange of virtual scalars, while (f) is the fermionic contribution.

3.1.2. Matter contributions

Of course, in Bagger–Lambert theory we also have contributions from matter fields as shown in (d)–(f) of Fig. 1 and it is important to ascertain how they contribute to the gluon self-energy. In this context it is important to note that, as is well known [14], F^2 regularisation does not regulate *all* the divergences in one loop diagrams. As such, one may use a supplementary regularisation scheme. Indeed, as far as the evaluation of integrals is concerned, dimensional continuation was already used to deal with those arising from the gluonic contributions. For the matter fields, we may be tempted to add a supersymmetrised F^2 term as a regulator in a similar spirit to the treatment of supersymmetric Yang–Mills–Chern–Simons theories in, e.g., [15,20]. However, this does not seem to make much sense as we already have standard kinetic terms for the fermions and scalars. Similarly, we could conceivably add mass-terms for the matter fields independently of the F^2 term. This is more like a regulator in the infrared and in any case would not give any contributions to the quantities of interest as we remove it. Thus, we will content ourselves with dealing with any remaining divergences using a dimensional continuation of the integrals where necessary.[4]

To begin with, scalars contribute via diagrams (d) and (e) of Fig. 1 and it is easily verified that they do not contribute to the antisymmetric part of $\Pi_{\mu\nu}$. A more generic way to view this is that scalars are not parity violating. Secondly, fermions run in the loop via the last diagram, (f) of Fig. 1, and we would like to look at their contribution to $\Pi_{\mathrm{a}}^{(1)}$. For these purposes it is enough to look at the momentum structure involved.

Each fermion vertex comes with a factor of Γ_μ, while the fermion propagators take the usual form $\Gamma_\mu p^\mu/p^2$. This means that after a little algebra

$$(\mathrm{f}) \propto \int \mathrm{d}^3 q \frac{q_\mu (q-p)_\nu + q_\nu (q-p)_\mu - \eta_{\mu\nu} q\cdot(q-p)}{q^2 (q-p)^2}, \tag{3.6}$$

which is manifestly symmetric under the interchange of μ and ν. The massless fermions in Bagger–Lambert thus do not contribute to the antisymmetric part of $\Pi_{\mu\nu}$ at one loop.

Massive fermions, on the other hand, *are* known to contribute to a shift in k, so it is interesting to note that even if we were to add a standard mass term such as $m_\psi \bar\psi\psi$ to help regulate the fermions, then these contributions vanish as we remove the regulator $m_\psi \to 0$. This is due to the fact that the antisymmetric part of (f) generated by this mass is schematically $m_\psi + \mathcal{O}(m_\psi^2)$, which vanishes as $m_\psi \to 0$. This is in accordance with expectations [16].

In conclusion, we can see that the matter fields do not contribute to the antisymmetric part of the gluon self-energy at one loop.

3.2. Ghost corrections

Now we must look at the contributions to the ghost self-energy and the ghost–gluon interaction. Since the ghosts do not couple to the matter fields it is obvious that the matter fields do not contribute to the ghost self-energy or the ghost–gluon interaction at all at one loop. Thus the one loop contributions to these quantities are precisely the same as in pure Chern–Simons theory (with F^2 regularisation), a case which has previously been investigated in detail in [14]. Thus we again spare the reader the details (referring instead to [14]) and present the results here.

The relevant diagrams are given in Fig. 2, and there is only one contribution to the ghost self-energy—that of Fig. 2(a). This evaluates to give

$$\Pi_c^{(1)} = -\frac{2}{3k} C_2 \,\mathrm{sgn}(k). \tag{3.7}$$

In terms of the $\bar c A c$ vertex corrections, many parts of the diagrams in Fig. 2 (b) and (c) cancel against each other and the remainder vanishes as $m \to \infty$ giving $Z_{\bar c A c} = 1$. This is in any case expected from general arguments [21].

As we can see, the matter fields in the game do not contribute to the finite renormalisation of k and with $Z_c = 1 - \Pi_c^{(1)}$ at one loop we get

$$\begin{aligned} k' &= k\big(1 + \Pi_{\mathrm{a}}^{(1)} + 2\Pi_c^{(1)}\big) \\ &= k + C_2 \,\mathrm{sgn}(k) \\ &= k + 2\,\mathrm{sgn}(k), \end{aligned} \tag{3.8}$$

where we have used that $C_2(\mathrm{SU}(2)) = 2$. Note the crucial presence of $\mathrm{sgn}(k)$. In Bagger–Lambert theory we have two SU(2) theories with opposite levels, k and $-k$. The fact that the correction depends on the sign of k means that the one loop corrections preserve this structure, i.e., the new levels are k' and $-k'$. Without this the SO(4) structure would be anomalous.

4. Mass deformation

We may also consider the mass-deformed version of Bagger–Lambert theory described in Section 2 which (when taken together with a mass-dependent potential) preserves $\mathcal{N} = 8$ supersymmetry [19]. As far as the quadratic terms go, this involves the addition of a term S_{mass} given in (2.3) to the action. The mass-

[4] In fact, none of the integrals involved pose any problems and we can happily evaluate them in d dimensions and simply set $d = 3$.

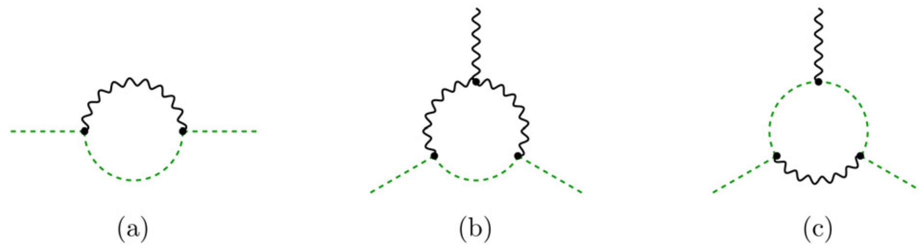

Fig. 2. Processes contributing to the ghost self-energy and the ghost–gluon vertex at one loop.

dependent potential is of fourth order in the scalars and therefore does not contribute to the quantities of interest to us at one loop.

It is again clear that the scalar mass term cannot change the contributions to $\Pi_{\mathrm{a}}^{(1)}$ since the scalars can still only give symmetric contributions and do not affect the ghosts at this level. However, now that we have a mass term for the fermions, they *can* in principle contribute. The reason is that the fermion propagator is now schematically of the form

$$\frac{\Gamma_\mu p^\mu + \mu\Gamma_{3456}}{p^2-\mu^2}, \tag{4.1}$$

and the terms arising from Fig. 1(f) which are proportional to μ contain an odd number of Γ_μ[5] and can therefore contribute to $\Pi_{\mathrm{a}}^{(1)}$. Nonetheless, as we see below, the contribution does in fact vanish.

A simple way to see this is that Γ_{3456} squares to one and is traceless. Its eigenvalues are thus equal numbers of ± 1 and so the mass deformation of (2.3) is like adding equal numbers of fermions with mass μ as with mass $-\mu$. Since their antisymmetric contribution to the gluon self-energy is proportional to their mass, these contributions cancel out. It is reassuring to note that this is in accordance with general arguments which give contributions to k' proportional to $\mathrm{sgn}(\mu)$ [16]. The presence of equal numbers of oppositely-signed massive fermions is thus expected to give no overall contribution.

From a more covariant point of view, where the fermions are still packaged into the single spinor ψ, the antisymmetric part of Fig. 1(f) is easily calculated to be

$$(\mathrm{f})_{\mathrm{a}} \propto \mathrm{tr}(\Gamma_\mu\Gamma_\nu\Gamma_\lambda\Gamma_{3456}). \tag{4.2}$$

By Lorentz invariance this can only be proportional to $\epsilon_{\mu\nu\lambda}$ and in order to fix the constant of proportionality we can consider the case of $\mu=0, \nu=1, \lambda=2$. By considering the relation $\Gamma_{0123456789(10)} = -\mathbb{1}$ we can see that

$$\Gamma_{0123456} = \pm\Gamma_{789(10)}, \tag{4.3}$$

depending on the signature of spacetime, and since $\mathrm{tr}\,\Gamma_{789(10)} = 0$ it is clear that the constant of proportionality is just zero. Thus these massive fermions do not contribute to a shift in k.[6] It is quite satisfying that the supersymmetry preserving mass deformation leaves k invariant even though a canonical mass deformation would certainly lead to a different shift in the one loop correction to k.

In conclusion, in both the original Bagger–Lambert theory and the deformed version, one expects one loop quantum corrections to shift the coupling by two:

$$k \to k + 2\,\mathrm{sgn}(k). \tag{4.4}$$

5. Discussion

In the proposed more general theory of [9] (see also [22]) with only $\mathcal{N}=6$ supersymmetry manifest there is $\mathrm{U}(N)\times\mathrm{U}(N)$ bifundamental matter coupled to Chern–Simons. In that theory, the link to Bagger–Lambert is that for $N=2$ the theory is expected to have extra symmetries which promote the $\mathcal{N}=6$ supersymmetry to $\mathcal{N}=8$. However, as a starting point for the $\mathcal{N}=6$ theory one may take $\mathcal{N}=4$ super-Yang–Mills plus Chern–Simons terms and integrate out the massive fields. In doing so, as discussed in [9] the fermions would cause a shift in k. This shift would then in that theory be subsequently cancelled by the shift in k due to the one loop correction from the Chern–Simons field as described above. Thus, overall there would be no shift in k. From this perspective the Bagger–Lambert theory as written above would be an effective theory where one loop effects have already been included. One cannot say a priori whether this is correct though there is now more evidence as to the success of [9] with [23].

This Letter took the approach of looking at the possible shift in k from a perturbative point of view using regularisation by addition of regulator terms such as the Yang–Mills term. In this context it does not seem to be possible to regularise the theory in the UV in a supersymmetric way. However, there may be a sense in which it *is* possible to regulate the theory while preserving at least some of the supersymmetry: One could consider regulating by replacing the entire Bagger–Lambert action with a supersymmetric YM–CS action of the sort encountered in [15].[7] As long as the specific form of this action preserves $\mathcal{N}\geqslant 2$ supersymmetries[8] then performing a one loop renormalisation would lead to a cancellation between bosons and fermions such that there is no overall shift in k [15]. Removing the UV regulator should be equivalent to integrating out the massive fields and thus one would expect to recover Bagger–Lambert theory in this limit à la [9] and without any shift in k.[9] On the other hand, it seems somewhat drastic to regulate by replacing the entire action with something new.

Thus, although the two approaches of [2–5] and of [9] are classically equivalent, they may not be so at the quantum level. A standard regularisation by addition of regulator terms as we have examined would seem to break the quantum equivalence and lead to a one loop shift in k for BL. If one additionally makes the assumption that the calculated shift is true for all k (which we

[5] Recall that the gamma matrices are split into Γ_μ for $\mu=0,1,2$ and Γ_A for $A=3,\dots,10$.

[6] It is interesting to note here that we could have used this mass deformation as a supersymmetric regulator for the matter fields and taken $\mu\to 0$ at the end of the day. Of course this would give the same results as previously found in Section 3.

[7] We would like to thank Seok Kim for illuminating discussions on this point.

[8] Though it does not seem likely that we could preserve the full $\mathcal{N}=8$ SUSY, at least explicitly.

[9] Note that if it is possible to do this procedure with a YM–CS action preserving only $\mathcal{N}=1$ supersymmetry, which is not entirely clear, then one would still see a shift in k.

imagine to be the case but certainly cannot derive given the perturbative nature of these calculations), then the moduli space at $k = 1$ would become corrected which seems unphysical.

Acknowledgements

We would like to thank Nick Dorey, Seok Kim, Costis Papageorgakis, Riccardo Ricci, Dan Thompson, Dave Tong and Gabriele Travaglini for useful discussions. This work was in part supported by the EC Marie Curie Research Training Network, MRTN-CT-2004-512194. J.B. is supported by an STFC Research Fellowship.

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