Anomaly-free nonsupersymmetric large N gauge theories from orientifolds

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We construct anomaly-free nonsupersymmetric large N gauge theories from orientifolds of type IIB on $\mathbf{C}^3/\Gamma$ orbifolds. In particular, massless as well as tachyonic one-loop tadpoles are canceled in these models. This is achieved by starting with $\mathcal{N}=1,2$ supersymmetric orientifolds with a well defined world-sheet description and including discrete torsion (which breaks supersymmetry) in the orbifold action. In this way we obtain nontrivial nonchiral as well as anomaly-free chiral large N gauge theories. We point out certain subtleties arising in the chiral cases. Subject to certain assumptions, these theories are shown to have the property that computation of any M-point correlation function in these theories reduces to the corresponding computation in the parent $\mathcal{N}=4$ oriented theory. This generalizes the analogous results recently obtained in supersymmetric large N gauge theories from orientifolds, as well as in (non)supersymmetric large N gauge theories without orientifold planes. [S0556-2821(99)03804-7]

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I. INTRODUCTION

Recent developments in the AdS conformal field theory (CFT) correspondence (see, e.g., [1–5]) motivated a set of conjectures proposed in [6,7] which state that certain gauge theories with $\mathcal{N}=0,1$ supersymmetries are (super)conformal. These gauge theories are constructed by starting from a $U(N)$ gauge theory with $\mathcal{N}=4$ space-time supersymmetry in four dimensions, and orbifolding by a finite discrete subgroup Γ of the R-symmetry group Spin(6) [7]. These conjectures were shown at one-loop level for $\mathcal{N}=0$ theories [6,7], and to two loops for $\mathcal{N}=1$ theories using ordinary field theory techniques [7].

In the subsequent development [8] these conjectures were shown to be correct to all loop orders in the large N limit of 't Hooft [9] (also see [10][1]). In particular, in [8] the above gauge theories were obtained in the $\alpha'\to 0$ limit of type IIB with N parallel D3-branes imbedded in an orbifolded space-time. The work in [8] was generalized in [13] where the type IIB string theory included both D3-branes and orientifold 3-planes (with the transverse space being $\mathbf{C}^3/\Gamma$). In certain cases string consistency also requires presence of D7-branes and orientifold 7-planes. This corresponds to type IIB orientifolds. Introducing orientifold planes is necessary to obtain SO and Sp gauge groups (without orientifold planes the gauge group is always unitary), and also allows for additional variety in possible matter content.

In the presence of orientifold planes the string world-sheet topology is characterized by the numbers b of boundaries (corresponding to D-branes), c of cross-caps (corresponding to orientifold planes), and g of handles (corresponding to closed string loops). Such a world-sheet is weighted with

$$(N\lambda_s)^b\lambda_s^c\lambda_s^{2g-2}=\lambda^{2g-2+b+c}N^{-c-2g+2}. \qquad (1)$$

The 't Hooft's large N limit then corresponds to taking the limit $N\to\infty$ with $\lambda=N\lambda_s$ fixed, where λ_s is the type IIB string coupling. (Here we identify $\lambda_s=g_{YM}^2$, where g_{YM} is the Yang-Mills coupling of the D3-brane gauge theory.) Note that addition of a cross-cap or a handle results in a diagram suppressed by an additional power of N, so that in the large N limit the contributions of cross-caps and handles are subleading. In fact, in [13] it was shown that for string vacua which are perturbatively consistent (that is, the tadpoles cancel) calculations of correlation functions in $\mathcal{N}<4$ gauge theories reduce to the corresponding calculations in the parent $\mathcal{N}=4$ *oriented* theory. This holds not only for finite (in the large N limit) gauge theories but also for the gauge theories which are not conformal. (In the latter case the gauge coupling running was shown to be suppressed in the large N limit.)

One distinguishing feature of large N gauge theories obtained via orientifolds is that the number of possibilities which possess well defined world-sheet expansion (that is, are perturbative from the orientifold viewpoint) is rather limited. In particular, in [13–15] (also see [16,17]) $\mathcal{N}=1$ and $\mathcal{N}=2$ supersymmetric large N gauge theories from orientifolds were studied in detail. Construction of these models is accompanied by certain subtleties. Thus, naively it might appear that the corresponding orientifolds of type IIB on $\mathbf{R}^6/\Gamma$ should result in theories with well defined world-sheet description for any orbifold group $\Gamma\subset\mathrm{Spin}(6)$. This is, however, not the case. For supersymmetric models it was shown in [13–15] that only a small number of orbifold groups results in such theories. In all the other cases the corresponding orientifolds contain nonperturbative states (which can be viewed as arising from D-branes wrapping various collapsed two-cycles in the orbifold). Multiple independent checks [18,13–15] have confirmed these conclusions.

Although the present understanding of large N supersymmetric gauge theories from orientifolds appears to be rather complete [13,14,16,15], no nonsupersymmetric examples of such theories have been constructed. In fact, there are certain difficulties associated with such a construction. Two of the main reasons that make it nontrivial to construct nonsupersymmetric large N gauge theories from orientifolds are the

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[1]For other related works, see, e.g., [11,12].

following. First, just as in the supersymmetric case, *a priori* we expect nonperturbative contributions to the orientifold spectrum at least for some choices of the orbifold group. Checking whether such states are present in a given nonsupersymmetric orientifold is much more nontrivial than in the supersymmetric cases where one can use type I-heterotic duality [19] along the lines of [20,21,18,15], as well as F theory [22] considerations [18]. Moreover, in the cases without supersymmetry the corresponding orientifolds generically contain tachyons in (some of) the twisted closed string sectors. As we explain in the following, this by itself does not affect the consistency of the gauge theory in the large N limit (as the closed string sector decouples). However, one potentially has twisted as well as massless tadpoles, and generically the tadpole cancellation conditions are overconstrained.

In this paper we construct a class of non-supersymmetric large N gauge theories from orientifolds in which all the tadpoles (that is, both massless and tachyonic ones) are cancelled. This is achieved by starting with $\mathcal{N}=1,2$ supersymmetric orientifolds of type IIB on $\mathbf{C}^3/\Gamma$ with well-defined world-sheet expansion, and considering nonsupersymmetric orientifolds of type IIB on $\mathbf{C}^3/\Gamma'$, where Γ' is obtained via a modification of the action of Γ. This modification amounts to including *discrete torsion* such that a $\mathbf{Z}_2$ subgroup of Γ' acts differently on space-time bosonic and fermionic sectors of the orientifold. This way we obtain nontrivial anomaly free nonsupersymmetric large N gauge theories. In particular, we find both chiral and nonchiral gauge theories. In the former case there are certain subtleties (related to possible nonperturbative states) which we explain in detail in Sec. IV.

The remainder of this paper is organized as follows. In Sec. II we review some of the important points in the orientifold construction, and explain how nontrivial discrete torsion breaks supersymmetry. In Sec. III we construct nonchiral models. In Sec. IV we give a construction of chiral models. There we also point out certain subtleties that arise in the construction of these models. In Sec. V we give our conclusions.

II. PRELIMINARIES

In this section we review the setup in [13] which leads to supersymmetric large N gauge theories from orientifolds. We then discuss how to obtain nonsupersymmetric large N gauge theories from orientifolds that are free of (both massless and tachyonic) tadpoles (and, consequently, of space-time anomalies).

A. Setup

Consider type IIB string theory on $\mathbf{C}^3/\Gamma$ where $\Gamma\subset SU(3)[SU(2)]$ so that the resulting theory has $\mathcal{N}=2(4)$ supersymmetry in four dimensions. In the following we will use the following notations: $\Gamma=\{g_a|a=1,\ldots,|\Gamma|\}$ ($g_1=1$). Consider the ΩJ orientifold of this theory, where Ω is the world-sheet parity reversal, and J is a $\mathbf{Z}_2$ element ($J^2=1$) acting on the complex coordinates z_i ($i=1,2,3$) on $\mathbf{C}^3$ as follows: $Jz_i=-Jz_i$. The resulting theory has $\mathcal{N}=1(2)$ supersymmetry in four dimensions.

Note that we have an orientifold 3-plane corresponding to the ΩJ element of the orientifold group. If Γ has a $\mathbf{Z}_2$ subgroup, then we also have an orientifold 7-plane. If we have an orientifold 7-plane we must introduce 8 of the corresponding D7-branes to cancel the R-R charge appropriately. (The number 8 of D7-branes is required by the corresponding tadpole cancellation conditions.) Note, however, that the number of D3-branes is not constrained (for the corresponding untwisted tadpoles automatically vanish in the noncompact case).

We need to specify the action of Γ on the Chan-Paton factors corresponding to the D3 and D7 branes. These are given by Chan-Paton matrices which we collectively refer to as $n^\mu\times n^\mu$ matrices γ^μ_a, where the superscript μ refers to the corresponding D3 or D7 branes. Note that $\mathrm{Tr}(\gamma^\mu_1)=n^\mu$ where n^μ is the number of D branes labelled by μ.

At one-loop level there are three different sources for massless tadpoles: the Klein bottle, annulus, and Möbius strip amplitudes. The factorization property of string theory implies that the tadpole cancellation conditions read (see, e.g., [13] for a more detailed discussion):

$$B_a+\sum_\mu C^\mu_a\mathrm{Tr}(\gamma^\mu_a)=0. \tag{2}$$

Here B_a and C^μ_a are (model dependent) numerical coefficients of order 1.

In the world-volume of D3 branes there lives a four dimensional $\mathcal{N}=1(2)$ supersymmetric gauge theory (which is obtained in the low energy, that is, $\alpha'\to 0$ limit). Since the number of D3 branes is unconstrained, we can consider the large N limit of this gauge theory. In [13] (generalizing the work in [8]) it was shown that, if for a given choice of the orbifold group Γ the world-sheet description for the orientifold is adequate, then in the large N limit (with $\lambda=N\lambda_s$ fixed, where λ_s is the type IIB string coupling) computation of any correlation function in this gauge theory is reduced to the corresponding computation in the parent $\mathcal{N}=4$ supersymmetric *oriented* gauge theory before orbifolding and orientifolding. In particular, the running of the gauge coupling is suppressed in the large N limit. Moreover, if

$$\mathrm{Tr}(\gamma^\mu_a)=0\,\forall a\neq 1 \tag{3}$$

(that is, $B_a=0\,\forall a\neq 1$), then the one-loop β-function coefficients b_0 for non-Abelian gauge theories living in world-volumes of the D3 branes vanish.

B. Perturbative orientifolds

The arguments of [13] that imply the above properties of D3-brane gauge theories are intrinsically perturbative. In particular, a consistent world-sheet expansion is crucial for their validity. It is therefore important to understand the conditions for the perturbative orientifold description to be adequate.

Naively, one might expect that any choice of the orbifold group $\Gamma\subset\mathrm{Spin}(6)$ [note that Spin(6) is the R-symmetry group of $\mathcal{N}=4$ gauge theory] should lead to an orientifold with well defined world-sheet expansion in terms of boundaries (corresponding to D branes), cross-caps (corresponding to orientifold planes) and handles (corresponding to closed string loops). This is, however, not the case [18,14,15]. In fact, the number of choices of Γ for which such a world-sheet expansion is adequate is rather constrained. In particular, in [18,14,15] it was argued that for $\mathcal{N}=1$ there are only seven choices of the orbifold group leading to consistent perturbative orientifolds: $\mathbf{Z}_2\otimes\mathbf{Z}_2$ [23], $\mathbf{Z}_3$ [24], $\mathbf{Z}_7$, $\mathbf{Z}_3\otimes\mathbf{Z}_3$ and $\mathbf{Z}_6$ [21], $\mathbf{Z}_2\otimes\mathbf{Z}_2\otimes\mathbf{Z}_3$ [25], and $\Delta(3\cdot 3^2)$ (the latter group is non-Abelian) [15]. All the other orbifold groups (including those considered in [26,27]) lead to orientifolds containing sectors which are nonperturbative from the orientifold viewpoint (that is, these sectors have no world-sheet description). These sectors can be thought of as arising from D branes wrapping various (collapsed) 2-cycles in the orbifold. The above restrictions on the orbifold group will be important in the construction of consistent non-supersymmetric large N gauge theories from orientifolds.

C. Discrete torsion and nonsupersymmetric models

The question we are going to address next is if we can obtain *nonsupersymmetric* large N gauge theories from orientifolds via generalizing the above construction for supersymmetric theories. Such a generalization might naively seem to be straightforward. However, there are certain subtleties here. Thus, in nonsupersymmetric theories we generically have tachyons (which, in particular, will be the case in the models constructed in this paper) in the twisted closed string sectors. As we will point out in a moment, the presence of tachyons by itself does not pose a problem for the consistency of the corresponding large N theories. However, if tachyons are present in the physical spectrum of the corresponding orientifold model, *a priori* they too contribute into the tadpoles. Moreover, the cancellation conditions for the tachyonic and massless tadpoles generically are rather different. That is, the numerical coefficients B_a and C_a in Eq. (2) corresponding to the massless and tachyonic tadpoles are generically different. This typically overconstrains the tadpole cancellation conditions, which makes it rather difficult to find tadpole free nonsupersymmetric orientifolds.

There is a way around the above difficulties, however. Let $\Gamma\subset SU(3)$ be an orbifold group such that it contains a $\mathbf{Z}_2$ subgroup. Let the generator of this $\mathbf{Z}_2$ subgroup be R. Consider now the following orbifold group: $\Gamma'=\{g'_a|a=1,\dots,|\Gamma|\}$, where the elements g'_a are the same as g_a except that R is replaced everywhere by $R'=RT$, where T is the generator of a $\mathbf{Z}_2$ group corresponding to the *discrete torsion*.[2] The action of T is defined as follows: it acts as identity in the bosonic sectors [that is, in the Nevev-Schwarz–Nevev-Schwarz (NS-NS) and Ramond-Ramond (R-R) closed string sectors, and in the NS open string sector], and it acts as -1 in the fermionic sectors (that is, in the NS-R and R-NS closed string sectors, and in the R open string sector) when, say, acting on the ground states. Now consider type IIB on $\mathbf{C}^3/\Gamma'$. Generically, in this theory all supersymmetries are broken. There are certain ''exceptions,'' however. Thus, if $\Gamma\approx\mathbf{Z}_2$ such that $\Gamma\subset SU(2)$, then inclusion of the discrete torsion does not break supersymmetry. Similarly, if $\Gamma\approx\mathbf{Z}_2\otimes\mathbf{Z}_2$ such that $\Gamma\subset SU(3)$, the number of unbroken supersymmetries is not affected by the discrete torsion. The basic reason for this is that the $\mathbf{Z}_2$ twist is self-conjugate. On the other hand, in certain cases we can include the discrete torsion in ways slightly different from the one just described. In particular, let $\Gamma\approx\mathbf{Z}_4$ such that $\Gamma\subset SU(2)$. Let g be the generator of this $\mathbf{Z}_4$. Next, consider the following orbifold group: $\Gamma'=\{1,gT,g^2,g^3T\}$. In other words, the $\mathbf{Z}_4$ twists g and g^3 (but not the corresponding $\mathbf{Z}_2$ twist g^2) are accompanied by the discrete torsion T. In this case we also have no unbroken supersymmetries. More generally, we can include the discrete torsion if the orbifold group Γ contains a $\mathbf{Z}_{2^n}$ subgroup. However, the corresponding orbifold group Γ' does not always lead to nonsupersymmetric theories.

Next, suppose using the above construction we have found an orbifold group Γ' such that Type IIB on $\mathbf{C}^3/\Gamma'$ is nonsupersymmetric. Then it is not difficult to show that the following statement holds. If Γ is such that the ΩJ orientifold of Type IIB on $\mathbf{C}^3/\Gamma$ is a perturbatively well defined $\mathcal{N}=1$ or $\mathcal{N}=2$ theory (that is, all the massless tadpoles cancel, and there are no nonperturbative contributions to the massless spectrum), then in the ΩJ orientifold of type IIB on $\mathbf{C}^3/\Gamma'$ (which is nonsupersymmetric) all the tachyonic and massless one-loop tadpoles (and, consequently, all the anomalies) automatically cancel. This fact is the key observation in the construction of anomaly free nonsupersymmetric large N gauge theories from orientifolds which we give in the subsequent sections.

D. Large N limit

As we already mentioned above, in all the nonsupersymmetric models constructed in this paper there are tachyons in some of the twisted closed string sectors. [These are Gliozzi-Scherk-Olive (GSO) projected out in orbifolds without the discrete torsion but are kept if the discrete torsion is nontrivial.] This immediately raises a question of whether the corresponding orientifolds are meaningful. In particular, in the presence of tachyons we expect vacuum instability. However, there is a subtlety here which saves the day. The point is that here we are interested in large N gauge theories in the 't Hooft limit $N\lambda_s=$fixed, which implies that the closed string coupling constant $\lambda_s\rightarrow 0$ as we take N to infinity. In particular, all the world-sheets with handles (corresponding to closed string loops) as well as cross-caps (corresponding to orientifold planes) are suppressed in this limit. That is, the closed string sector decouples from the open string sector in this limit, and after taking $\alpha'\rightarrow 0$ we can ignore the closed string states (regardless of whether they are tachyonic or not)

[2]Note that ''discrete torsion'' T only acts on the space-time fermionic sectors, and should not be confused with discrete torsion as used in orbifolds such as $\mathbf{Z}_2\otimes\mathbf{Z}_2$. Thus, the action of T in the closed string sectors can be written as $(-1)^{F_L+F_R}$, where F_L and F_R are the space-time fermion numbers in left- and right-moving closed string sectors, respectively.

altogether. In other words, the string construction here is simply an efficient and fast way of obtaining a field theory result, and at the end of the day we are going to throw out all the irrelevant ingredients (such as complications in the closed string sectors due to the presence of tachyons) and keep only those relevant for the field theory discussion. Note that we would not be able to do the same had we considered a *compact* model (with a finite number of D3 branes). In this case the closed string sector does not decouple and the theory is sick due to the presence of tachyons.

Thus, in the large N limit of 't Hooft (accompanied by taking α' to zero) the nonsupersymmetric gauge theory living in the world volumes of the D3 branes is completely well defined, and following [13] we conclude that in the nonsupersymmetric gauge theories arising from the orientifolds with well defined world-sheet expansions computation of any correlation function reduces to the corresponding computation in the parent $\mathcal{N}=4$ oriented gauge theory (before orbifolding and orientifolding). The only remaining question is which of these nonsupersymmetric orientifolds have well defined world-sheet expansion. The answer to this question should be clear from our previous discussions: as long as the supersymmetric ΩJ orientifold of type IIB on $\mathbf{C}^3/\Gamma$ is perturbatively well defined, we expect the corresponding nonsupersymmetric ΩJ orientifold of type IIB on $\mathbf{C}^3/\Gamma'$ to also possess a well defined world-sheet expansion. This will be our guiding principle in constructing consistent nonsupersymmetric large N gauge theories from orientifolds in the following sections. The fact that non-Abelian gauge anomalies cancel nontrivially in the models discussed in this paper indicates self-consistency of this assumption. (However, as we discuss in Sec. IV, such an expectation might not hold in some cases.)

III. NONCHIRAL $\mathcal{N}=0$ GAUGE THEORIES

In this section we construct nonchiral large N gauge theories from nonsupersymmetric type IIB orientifolds. The idea here is the following. Consider type IIB on $\mathbf{C}\otimes(\mathbf{C}^2/\Gamma)$ where $\Gamma\subset SU(2)$. Suppose now Γ contains a $\mathbf{Z}_2$ subgroup. If there is no discrete torsion accompanying its generator in the fermionic sectors, then the theory will be $\mathcal{N}=2$ supersymmetric (after orientifolding). However, if we include nontrivial discrete torsion, the supersymmetry is going to be broken. Actually, if $\Gamma\approx\mathbf{Z}_2$ [such that $\Gamma\subset SU(2)$], then (as we already mentioned in Sec. II) including the discrete torsion cannot break supersymmetry. On the other hand, as was shown in [13] in detail, only $\Gamma\approx\mathbf{Z}_M$, $M=2,3,4,6$ orbifold groups lead to perturbatively well defined ΩJ orientifolds of type IIB on $\mathbf{C}\otimes(\mathbf{C}^2/\Gamma)$ without discrete torsion (and such orientifolds have $\mathcal{N}=2$ supersymmetry). In all the other cases some of the massless tadpoles cannot be cancelled. We will therefore concentrate on these orbifold groups. The cases $M=2,3$ are of no interest to us: for $M=2$, as we just discussed, inclusion of the discrete torsion does not break supersymmetry; in the $M=3$ case we have no $\mathbf{Z}_2$ subgroup, hence we cannot have discrete torsion. The only cases left then are the $\mathbf{Z}_4$ and $\mathbf{Z}_6$ cases.

We are now ready to give an explicit construction of large N gauge theories from the above nonsupersymmetric orientifolds.

A. The $\mathbf{Z}_6$ orbifold

Let g and R be the generators of the $\mathbf{Z}_3$ and $\mathbf{Z}_2$ subgroups of the orbifold group $\Gamma\approx\mathbf{Z}_6\approx\mathbf{Z}_3\otimes\mathbf{Z}_2$. The action of g and R on the complex coordinates z_s is given by

$$gz_1=z_1,\quad gz_2=\omega z_2,\quad gz_3=\omega^{-1}z_3,\quad \omega=(2\pi i/3), \tag{4}$$

$$Rz_1=z_1,\quad Rz_2=-z_2,\quad Rz_3=-z_3. \tag{5}$$

Now consider the orbifold group Γ' where the $\mathbf{Z}_2$ twist is accompanied by nontrivial discrete torsion, that is, the generator R is replaced by RT. Supersymmetry is broken completely in this case.

In this model we have n_3 D3 branes, and 8 D7 branes. The world volumes of the D3 branes fill the noncompact space $\mathbf{R}^4$ transverse to the coordinates z_s. The world volumes of the D7 branes fill the noncompact space transverse to the z_1 coordinate. The solution to the twisted tadpole cancellation conditions is given by $[N=(n_3-2)/6]$

$$\gamma_{g,3}=\mathrm{diag}(\omega\mathbf{I}_{2N},\omega^{-1}\mathbf{I}_{2N},\mathbf{I}_{2N+2}), \tag{6}$$

$$\gamma_{R,3}=\mathrm{diag}(i,-i)\otimes\mathbf{I}_{3N+1}, \tag{7}$$

$$\gamma_{g,7}=\mathrm{diag}(\omega\mathbf{I}_2,\omega^{-1}\mathbf{I}_2,\mathbf{I}_4), \tag{8}$$

$$\gamma_{R,7}=\mathrm{diag}(i,-i)\otimes\mathbf{I}_4. \tag{9}$$

The massless spectrum of this model is given in Table I. The gauge group is $[U(N)^2\otimes U(N+1)]_{33}\otimes[U(1)^2\otimes U(2)]_{77}$. Note that the one-loop β-function coefficients $b_0(N)$ and $b_0(N+1)$ for the $SU(N)$ and $SU(N+1)$ subgroups of the 33 sector gauge group are independent of N:

$$b_0(N)=-3/2, \tag{10}$$

$$b_0(N+1)=+3. \tag{11}$$

Note that in some of the twisted closed string sectors we have tachyons, namely, the tachyons arise in the gR and $g^{-1}R$ twisted sectors. All the other closed string sectors are tachyon free.

B. The $\mathbf{Z}_4$ orbifold

Let g be the generator of the orbifold group $\Gamma\approx\mathbf{Z}_4$. The action of g on the complex coordinates z_s is given by

$$gz_1=z_1,\quad gz_2=iz_2,\quad gz_3=-iz_3. \tag{12}$$

Now consider the orbifold group Γ' where the $\mathbf{Z}_4$ twists g and g^3 are accompanied by nontrivial discrete torsion, that is, g and g^3 are replaced by gT and g^3T, respectively. (Note that this discrete torsion is such that the g^2 twist is torsion free.) Supersymmetry is broken completely in this case.

In this model we have n_3 D3 branes, and 8 D7 branes. The world volumes of the D3 branes fill the noncompact

TABLE I. The massless open string spectra of the $\mathcal{N}=0$ orientifolds of type IIB on $\mathbf{C}\otimes(\mathbf{C}^2/\mathbf{Z}_6)$ and type IIB on $\mathbf{C}\otimes(\mathbf{C}^2/\mathbf{Z}_4)$. The subscript "$b$" indicates that the corresponding field consists of the bosonic content of a hypermultiplet. (Thus, $\frac{1}{2}$ of this content corresponds to a complex scalar.) The subscript "f" indicates that the corresponding field consists of the fermionic content of a hypermultiplet (that is, of one left-handed and one right-handed chiral fermion in the corresponding representation of the gauge group). The notation **A** stands for the two-index antisymmetric representation of the corresponding unitary group, whereas **Adj** stands for the adjoint representation. For the sake of simplicity we have suppressed the $U(1)$ charges (which are not difficult to restore).

Model	Gauge group	Charged bosons	Charged fermions
$\mathbf{Z}_6$	$[U(N)^2\otimes U(N+1)]_{33}$	$3\times\frac{1}{2}[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1})_b]_{33}$	$[(\mathbf{N},\overline{\mathbf{N}},\mathbf{1};\mathbf{1})_f]_{33}$
	$\otimes[U(1)^2\otimes U(2)]_{77}$	$\frac{1}{2}[(\mathbf{Adj},\mathbf{1},\mathbf{1};\mathbf{1})_b]_{33}$	$[(\overline{\mathbf{N}},\mathbf{N},\mathbf{1};\mathbf{1})_f]_{33}$
		$\frac{1}{2}[(\mathbf{1},\mathbf{Adj},\mathbf{1};\mathbf{1})_b]_{33}$	$[(\mathbf{1},\mathbf{1},\mathbf{A};\mathbf{1})_f]_{33}$
		$\frac{1}{2}[(\mathbf{1},\mathbf{1},\mathbf{Adj};\mathbf{1})_b]_{33}$	$[(\mathbf{1},\mathbf{1},\overline{\mathbf{A}};\mathbf{1})_f]_{33}$
		$[(\mathbf{A},\mathbf{1},\mathbf{1};\mathbf{1})_b]_{33}$	$[(\mathbf{N},\mathbf{N},\mathbf{1};\mathbf{1})_f]_{33}$
		$[(\mathbf{1},\mathbf{A},\mathbf{1};\mathbf{1})_b]_{33}$	$[(\mathbf{N},\overline{\mathbf{1},\mathbf{N+1};\mathbf{1})}_f]_{33}$
		$[(\mathbf{N},\mathbf{1},\mathbf{N+1};\mathbf{1})_b]_{33}$	$[(\mathbf{1},\mathbf{N},\mathbf{N+1};\mathbf{1})_f]_{33}$
		$[(\mathbf{1},\overline{\mathbf{N},\mathbf{N+1};\mathbf{1})}_b]_{33}$	
		$\frac{1}{2}[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{3})_b]_{77}$	
		$3\times\frac{1}{2}[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1})_b]_{77}$	$5\times[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1})_f]_{77}$
		$2\times[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{2})_b]_{77}$	$2\times[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{2})_f]_{77}$
		$\frac{1}{2}[(\mathbf{N},\mathbf{1},\mathbf{1};\mathbf{1})_b]_{37}$	$[(\mathbf{N},\mathbf{1},\mathbf{1};\mathbf{1})_f]_{37}$
		$\frac{1}{2}[(\mathbf{N},\mathbf{1},\mathbf{1};\mathbf{2})_b]_{37}$	$[(\mathbf{1},\mathbf{N},\mathbf{1};\mathbf{1})_f]_{37}$
		$\frac{1}{2}[(\mathbf{1},\mathbf{N},\mathbf{1};\mathbf{1})_b]_{37}$	$[(\mathbf{1},\mathbf{1},\mathbf{N+1};\mathbf{2})_f]_{37}$
		$\frac{1}{2}[(\mathbf{1},\mathbf{N},\mathbf{1};\mathbf{2})_b]_{37}$	
		$2\times\frac{1}{2}[(\mathbf{1},\mathbf{1},\mathbf{N+1};\mathbf{1})_b]_{37}$	
$\mathbf{Z}_4$	$[U(N)^2]_{33}\otimes[U(2)^2]_{77}$	$[(\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})_b]_{33}$	
		$\frac{1}{2}[(\mathbf{Adj},\mathbf{1};\mathbf{1},\mathbf{1})_b]_{33}$	$[(\mathbf{N},\overline{\mathbf{N}};\mathbf{1},\mathbf{1})_f]_{33}$
		$\frac{1}{2}[(\mathbf{1},\mathbf{Adj};\mathbf{1},\mathbf{1})_b]_{33}$	$[(\overline{\mathbf{N}},\mathbf{N};\mathbf{1},\mathbf{1})_f]_{33}$
		$[(\mathbf{A},\mathbf{1};\mathbf{1},\mathbf{1})_b]_{33}$	$[(\mathbf{A},\mathbf{1};\mathbf{1},\mathbf{1})_f]_{33}$
		$[(\mathbf{1},\mathbf{A};\mathbf{1},\mathbf{1})_b]_{33}$	$[(\mathbf{1},\mathbf{A};\mathbf{1},\mathbf{1})_f]_{33}$
		$[(\mathbf{N},\mathbf{N};\mathbf{1},\mathbf{1})_b]_{33}$	$[(\mathbf{N},\mathbf{N};\mathbf{1},\mathbf{1})_f]_{33}$
		$3\times[(\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})_b]_{77}$	$2\times[(\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})_f]_{77}$
		$\frac{1}{2}[(\mathbf{1},\mathbf{1};\mathbf{3},\mathbf{1})_b]_{77}$	
		$\frac{1}{2}[(\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{3})_b]_{77}$	
		$[(\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{2})_b]_{77}$	$3\times[(\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{2})_f]_{77}$
		$[(\mathbf{N},\mathbf{1};\mathbf{2},\mathbf{1})_b]_{37}$	$[(\mathbf{N},\mathbf{1};\mathbf{2},\mathbf{1})_f]_{37}$
		$[(\mathbf{1},\mathbf{N};\mathbf{1},\mathbf{2})_b]_{37}$	$[(\mathbf{1},\mathbf{N};\mathbf{1},\mathbf{2})_f]_{37}$

space $\mathbf{R}^4$ transverse to the coordinates z_s. The world volumes of the D7 branes fill the noncompact space transverse to the z_1 coordinate. The solution to the twisted tadpole cancellation conditions is given by ($N=n_3/4$):

$$\gamma_{g,3}=\mathrm{diag}(\omega\mathbf{I}_N,\omega^{-1}\mathbf{I}_N,\omega^3\mathbf{I}_N,\omega^{-3}\mathbf{I}_N), \quad (13)$$

$$\gamma_{g,7}=\mathrm{diag}(\omega\mathbf{I}_2,\omega^{-1}\mathbf{I}_2,\omega^3\mathbf{I}_2,\omega^{-3}\mathbf{I}_2). \quad (14)$$

Here $\omega=\exp(\pi i/4)$. The massless spectrum of this model is given in Table I. The gauge group is $[U(N)^2]_{33}\otimes[U(2)^2]_{77}$. Note that the one-loop β-function coefficient $b_0(N)$ for each of the $SU(N)$ subgroups of the 33 sector gauge group is independent of N. In fact, this coefficient vanishes:

$$b_0(N)=0. \quad (15)$$

Thus, at the one-loop order the theory is conformal (even for finite N). According to [13], this property persists to all loop orders in 't Hooft's large N limit.

Note that in some of the twisted closed string sectors we have tachyons, namely, the tachyons arise in the g and g^3 twisted sectors. All the other closed string sectors are tachyon free.

C. Comments

Here some remarks are in order. The first comment is regarding the fact that in the $\mathbf{Z}_6$ case the one-loop β-function coefficients of the non-Abelian subgroups of the 33 open string sector gauge group are nonzero, whereas in the $\mathbf{Z}_4$ case they vanish. This is in accord with the observation of [13] that $b_0 = 0$ if all the twisted Chan-Paton matrices γ_a ($a \neq 1$) are traceless. Moreover, generically we do not expect b_0 coefficients to vanish unless all $\mathrm{Tr}(\gamma_a)$ ($a \neq 1$) are traceless. However, there can be ''accidental'' cancellations in some models such that all $b_0 = 0$ despite some of the twisted Chan-Paton matrices not being traceless. Such ''accidental'' cancellations, in particular, occur in the $\mathcal{N}=2$ supersymmetric $\mathbf{Z}_3$ and $\mathbf{Z}_6$ models discussed in [13]. These cancellations were explained in [13] using the results obtained in [28]. On the other hand, such an ''accidental'' cancellation does not occur in the above nonsupersymmetric $\mathbf{Z}_6$ model.

The second remark is related to the following. As discussed in [29,18], the orientifold projection Ω in the above cases maps the g_a twisted sector to its inverse g_a^{-1} twisted sector. As discussed at length in [18], such a projection implies that there are no nonperturbative (from the orientifold viewpoint) states arising in sectors corresponding to the orientifold group elements Ωg_a and Ωg_a^{-1} for $g_a^2 \neq 1$. This property of such orientifolds is independent of the space-time supersymmetry. This is not a trivial statement as it need not hold generically, namely, it is far from being obvious in the cases we discuss in the next section.

IV. CHIRAL $\mathcal{N}=0$ GAUGE THEORIES

In this section we construct chiral large N gauge theories from nonsupersymmetric type IIB orientifolds. We start with type IIB on $\mathbf{C}^3/\Gamma$, where Γ is one of the $SU(3)$ subgroups (discussed in subsection B of Sec. II) leading to perturbative orientifolds. For us to be able to include nontrivial discrete torsion, Γ must contain a $\mathbf{Z}_2$ subgroup. As we already mentioned in Sec. II, including discrete torsion in the $\mathbf{Z}_2 \otimes \mathbf{Z}_2$ case does not break supersymmetry. We are therefore led to consider the $\mathbf{Z}_6$ and $\mathbf{Z}_2 \otimes \mathbf{Z}_2 \otimes \mathbf{Z}_3$ cases only.

A. The $\mathbf{Z}_6$ orbifold

Let g and R be the generators of the $\mathbf{Z}_3$ and $\mathbf{Z}_2$ subgroups of the orbifold group $\Gamma \approx \mathbf{Z}_6 \approx \mathbf{Z}_3 \otimes \mathbf{Z}_2$. The action of g and R on the complex coordinates z_s is given by

$$g z_s = \omega z_s, \quad \omega = \exp(2\pi i/3), \tag{16}$$

$$R z_1 = -z_1, \quad R z_2 = -z_2, \quad R z_3 = z_3. \tag{17}$$

Now consider the orbifold group Γ' where the $\mathbf{Z}_2$ twist R is accompanied by nontrivial discrete torsion, that is, R is replaced by RT. Supersymmetry is broken completely in this case.

In this model we have n_3 D3 branes, and 8 D7 branes. The world volumes of the D3 branes fill the noncompact space $\mathbf{R}^4$ transverse to the coordinates z_s. The world-volumes of the D7 branes fill the noncompact space transverse to the coordinate z_3. The solution to the twisted tadpole cancellation conditions is given by $[N = (n_3 + 4)/6]$:

$$\gamma_{g,3} = \mathrm{diag}(\omega \mathbf{I}_{2N}, \omega^{-1} \mathbf{I}_{2N}, \mathbf{I}_{2N-4}), \tag{18}$$

$$\gamma_{R,3} = \mathrm{diag}(i, -i) \otimes \mathbf{I}_{3N-2}, \tag{19}$$

$$\gamma_{g,7} = \mathrm{diag}(\omega \mathbf{I}_4, \omega^{-1} \mathbf{I}_4), \tag{20}$$

$$\gamma_{R,7} = \mathrm{diag}(i, -i) \otimes \mathbf{I}_4. \tag{21}$$

The massless spectrum of this model is given in Table II. The gauge group is $[U(N)^2 \otimes U(N-2)]_{33} \otimes [U(2)^2]_{77}$. Note that the non-Abelian gauge anomaly is cancelled in this model. Moreover, the one-loop β-function coefficients $b_0(N)$ and $b_0(N-2)$ for the $SU(N)$ and $SU(N-2)$ subgroups of the 33 sector gauge group are independent of N:

$$b_0(N) = +3, \tag{22}$$

$$b_0(N-2) = -6. \tag{23}$$

Note that in the gR and g^2R twisted closed string sectors there are physical tachyons. All the other closed string sectors are tachyon free.

B. The $\mathbf{Z}_2 \otimes \mathbf{Z}_2 \otimes \mathbf{Z}_3$ orbifold

Let g, R_1 and R_2 be the generators of the $\mathbf{Z}_3$ and the two $\mathbf{Z}_2$ subgroups of the orbifold group $\Gamma \approx \mathbf{Z}_2 \otimes \mathbf{Z}_2 \otimes \mathbf{Z}_3$. The action of g and R_s ($R_3 = R_1 R_2$) on the complex coordinates $z_{s'}$ is given by (there is no summation over the repeated indices here):

$$g z_s = \omega z_s, \quad \omega = \exp(2\pi i/3), \tag{24}$$

$$R_s z_{s'} = -(-1)^{\delta_{ss'}} z_{s'}. \tag{25}$$

Without loss of generality we can consider the orbifold group Γ' where R_1 has no discrete torsion, whereas R_2 (and, therefore, R_3) is accompanied by nontrivial discrete torsion. (That is, R_2 and R_3 are replaced by R_2T and R_3T, respectively.)

In this model we have n_3 D3 branes, and three sets of D7 branes, which we refer to as D7_s branes, with 8 D7 branes in each set. The world volumes of the D3 branes fill the noncompact space $\mathbf{R}^4$ transverse to the coordinates z_s. The world volumes of the D7_s branes fill the noncompact space transverse to the coordinate z_s. The solution to the twisted tadpole cancellation conditions is given by $[N = (n_3 + 4)/6]$

$$\gamma_{g,3} = \mathrm{diag}(\mathbf{W} \otimes \mathbf{I}_N, \mathbf{I}_{2N-4}), \tag{26}$$

$$\gamma_{R_s,3} = i \sigma_s \otimes \mathbf{I}_{3N-2}. \tag{27}$$

Here $\mathbf{W} = \mathrm{diag}(\omega, \omega, \omega^{-1}, \omega^{-1})$. (The action on the D$7_s$ branes is similar.) The massless spectrum of this model is given in Table III. The gauge group is $[U(N) \otimes Sp(N-2)]_{33} \otimes_{s=1}^{3} [U(2)]_{7_s7_s}$. [Here we are using the convention where the rank of $Sp(2M)$ is M.] Note that the non-Abelian gauge anomaly is cancelled in this model. Moreover, the

TABLE II. The massless open string spectrum of the $\mathcal{N}=0$ orientifold of type IIB on $\mathbf{C}^3/\mathbf{Z}_6$. The subscript ‘‘$c$’’ indicates that the corresponding field is a complex boson. The subscript ‘‘L’’ indicates that the corresponding field is a left-handed chiral fermion. The notation **A** stands for the two-index antisymmetric representation of the corresponding unitary group. The $U(1)$ charges are given in parentheses.

Model	Gauge group	Charged complex bosons
$\mathbf{Z}_6$	$[U(N)\otimes U(N)\otimes U(N-2)]_{33}$	$2\times[(\mathbf{A},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})(+2,0,0;0,0)_c]_{33}$
	$\otimes[U(2)\otimes U(2)]_{77}$	$2\times[(\mathbf{1},\overline{\mathbf{A}},\mathbf{1};\mathbf{1},\mathbf{1})(0,-2,0;0,0)_c]_{33}$
		$2\times[(\overline{\mathbf{N}},\overline{\mathbf{1},\mathbf{N-2};\mathbf{1},\mathbf{1}})(-1,0,-1;0,0)_c]_{33}$
		$2\times[(\mathbf{1},\mathbf{N},\mathbf{N-2};\mathbf{1},\mathbf{1})(0,+1,+1;0,0)_c]_{33}$
		$[(\overline{\mathbf{N}},\mathbf{1},\mathbf{N-2};\mathbf{1},\mathbf{1})(-1,0,+1;0,0)_c]_{33}$
		$[(\mathbf{1},\overline{\mathbf{N},\mathbf{N-2};\mathbf{1},\mathbf{1}})(0,+1,-1;0,0)_c]_{33}$
		$[(\mathbf{N},\overline{\mathbf{N}},\mathbf{1};\mathbf{1},\mathbf{1})(+1,-1,0;0,0)_c]_{33}$
		$2\times[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})(0,0,0;+2,0)_c]_{77}$
		$2\times[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})(0,0,0;0,-2)_c]_{77}$
		$[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{2})(0,0,0;+1,-1)_c]_{77}$
		$[(\mathbf{N},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{2})(+1,0,0;0,+1)_c]_{37}$
		$[(\mathbf{1},\mathbf{N},\mathbf{1};\mathbf{2},\mathbf{1})(0,+1,0;+1,0)_c]_{37}$
		$[(\overline{\mathbf{N}},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{2})(-1,0,0;0,-1)_c]_{37}$
		$[(\mathbf{1},\overline{\mathbf{N}},\mathbf{1};\mathbf{2},\mathbf{1})(0,-1,0;-1,0)_c]_{37}$
		Charged chiral fermions
		$[(\mathbf{A},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})(+2,0,0;0,0)_L]_{33}$
		$[(\mathbf{1},\overline{\mathbf{A}},\mathbf{1};\mathbf{1},\mathbf{1})(0,-2,0;0,0)_L]_{33}$
		$[(\overline{\mathbf{N}},\overline{\mathbf{1},\mathbf{N-2};\mathbf{1},\mathbf{1}})(-1,0,-1;0,0)_L]_{33}$
		$[(\mathbf{1},\mathbf{N},\mathbf{N-2};\mathbf{1},\mathbf{1})(0,+1,+1;0,0)_L]_{33}$
		$2\times[(\overline{\mathbf{N}},\mathbf{1},\mathbf{N-2};\mathbf{1},\mathbf{1})(-1,0,+1;0,0)_L]_{33}$
		$2\times[(\mathbf{1},\overline{\mathbf{N},\mathbf{N-2};\mathbf{1},\mathbf{1}})(0,+1,-1;0,0)_L]_{33}$
		$2\times[(\mathbf{N},\overline{\mathbf{N}},\mathbf{1};\mathbf{1},\mathbf{1})(+1,-1,0;0,0)_L]_{33}$
		$[(\mathbf{N},\mathbf{N},\mathbf{1};\mathbf{1},\mathbf{1})(+1,+1,0;0,0)_L]_{33}$
		$[(\overline{\mathbf{N}},\overline{\mathbf{N}},\mathbf{1};\mathbf{1},\mathbf{1})(-1,-1,0;0,0)_L]_{33}$
		$[(\mathbf{1},\mathbf{1},\mathbf{A};\mathbf{1},\mathbf{1})(0,0,+2;0,0)_L]_{33}$
		$[(\mathbf{1},\mathbf{1},\overline{\mathbf{A}};\mathbf{1},\mathbf{1})(0,0,-2;0,0)_L]_{33}$
		$[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})(0,0,0;+2,0)_L]_{77}$
		$[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})(0,0,0;0,-2)_L]_{77}$
		$2\times[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{2})(0,0,0;+1,-1)_L]_{77}$
		$[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{2})(0,0,0;+1,+1)_L]_{77}$
		$[(\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{2})(0,0,0;-1,-1)_L]_{77}$
		$[(\overline{\mathbf{N}},\mathbf{1},\mathbf{1};\mathbf{2},\mathbf{1})(-1,0,0;-1,0)_L]_{37}$
		$[(\mathbf{1},\overline{\mathbf{1},\mathbf{N-2};\mathbf{1},\mathbf{2}})(0,0,-1;0,-1)_L]_{37}$
		$[(\mathbf{1},\mathbf{N},\mathbf{1};\mathbf{1},\mathbf{2})(0,+1,0;0,+1)_L]_{37}$
		$[(\mathbf{1},\mathbf{1},\mathbf{N-2};\mathbf{2},\mathbf{1})(0,0,+1;+1,0)_L]_{37}$

one-loop β-function coefficients $b_0(N)$ and $b_0(N-2)$ for the $SU(N)$ and $Sp(N-2)$ subgroups of the 33 sector gauge group are independent of N:

$$b_0(N)=+1, \tag{28}$$

$$b_0(N-2)=-2. \tag{29}$$

Note that in the gR_2, g^2R_2, gR_3 and g^2R_3 twisted closed string sectors there are physical tachyons. All the other closed string sectors are tachyon free.

C. Comments

As we already mentioned, not all choices of the orbifold group $\Gamma\subset SU(3)$ lead to perturbatively well defined $\mathcal{N}=1$ supersymmetric ΩJ orientifolds of type IIB on $\mathbf{C}^3/\Gamma$. Let us review the reasons responsible for such a limited number of perturbative orientifolds. Thus, consider the ΩJ orientifold of type IIB on $\mathbf{C}^3/\Gamma$. The orientifold group is given by $\mathcal{O}=\{g_a,\Omega J g_a \mid a=1,\ldots,|\Gamma|\}$. The sectors labeled by g_a correspond to the unoriented closed twisted plus untwisted sectors. The sectors labeled by $\Omega J g_a$ with $(Jg_a)^2=1$ correspond to open strings stretched between D-branes. (In

TABLE III. The massless open string spectrum of the $\mathcal{N}=0$ orientifold of type IIB on $\mathbf{C}^3/\mathbf{Z}_2\otimes\mathbf{Z}_2\otimes\mathbf{Z}_3$. The subscript "$c$" indicates that the corresponding field is a complex boson. The subscript "L" indicates that the corresponding field is a left-handed chiral fermion. The notation $\mathbf{A}$ stands for the two-index antisymmetric representation of $SU(N)$, whereas $\mathbf{S}$ stands for the two-index symmetric representation of $SU(N)$. The notation $\mathbf{Adj}$ stands for the N^2-1 dimensional adjoint representation of $SU(N)$, whereas $\mathbf{a}$ stands for the $N(N-1)/2-1$ dimensional *traceless* antisymmetric representation of $Sp(N)$. Also, $s=1,2,3$, and $k=2,3$. The $U(1)$ charges are given in parentheses.

Model	Gauge group	Charged complex bosons
$\mathbf{Z}_2\otimes\mathbf{Z}_2\otimes\mathbf{Z}_3$	$[U(N)\otimes Sp(N-2)]_{33}$	$3\times[(\mathbf{A},\mathbf{1})(+2)_c]_{33}$
	$\otimes_{s=1}^{3}[U(2)]_{7_s7_s}$	$3\times[(\overline{\mathbf{N}},\mathbf{N-2})(-1)_c]_{33}$
		$3\times[(\mathbf{1}_s)(+2_s)_c]_{7_s7_s}$
		$[(\mathbf{N},\mathbf{1};\mathbf{2}_1)(+1;+1_1)_c]_{37_1}$
		$[(\mathbf{1},\mathbf{N-2};\mathbf{2}_1)(0;-1_1)_c]_{37_1}$
		$[(\mathbf{2}_2;\mathbf{2}_3)(+1_2;+1_3)_c]_{7_27_3}$
		$[(\mathbf{N},\mathbf{1};\mathbf{2}_k)(+1;-1_k)_c]_{37_k}$
		$[(\overline{\mathbf{N}},\mathbf{1};\mathbf{2}_k)(-1;+1_k)_c]_{37_k}$
		$[(\mathbf{2}_1;\mathbf{2}_k)(+1_1;-1_k)_c]_{7_17_k}$
		Charged chiral fermions
		$[(\mathbf{Adj},\mathbf{1})(0)_L]_{33}$
		$[(\mathbf{1},\mathbf{a})(0)_L]_{33}$
		$2\times[(\mathbf{1},\mathbf{1})(0)_L]_{33}$
		$[(\mathbf{S},\mathbf{1})(+2)_L]_{33}$
		$2\times[(\mathbf{A},\mathbf{1})(+2)_L]_{33}$
		$3\times[(\overline{\mathbf{N}},\mathbf{N-2})(-1)_L]_{33}$
		$[(\mathbf{3}_s)(+2_s)_L]_{7_s7_s}$
		$2\times[(\mathbf{1}_s)(+2_s)_L]_{7_s7_s}$
		$[(\mathbf{3}_s)(0_s)_L]_{7_s7_s}$
		$2\times[(\mathbf{1}_s)(0)_L]_{7_s7_s}$
		$[(\mathbf{N},\mathbf{1};\mathbf{2}_1)(+1;+1_1)_L]_{37_1}$
		$[(\mathbf{1},\mathbf{N-2};\mathbf{2}_1)(0;-1_1)_L]_{37_1}$
		$[(\mathbf{2}_2,\mathbf{1}_2;\mathbf{2}_3,\mathbf{1}_3)(+1_2;+1_3)_L]_{7_27_3}$
		$[(\overline{\mathbf{N}},\mathbf{1};\mathbf{2}_k)(-1;-1_k)_L]_{37_k}$
		$[(\mathbf{1},\mathbf{N-2};\mathbf{2}_k)(0;+1_k)_L]_{37_k}$
		$[(\mathbf{2}_1,\mathbf{2}_k)(-1_1;-1_k)_L]_{7_17_k}$

particular, if the set of points fixed under Jg_a has dimension 0 then these are D3-branes. If this set has real dimension 4, then these are D7-branes.) However, the sectors labeled by ΩJg_a with $(Jg_a)^2\neq 1$ do not have an interpretation in terms of open strings starting and ending on perturbative D-branes (i.e., they do not have an interpretation in terms of open strings with purely Dirichlet or Neumann boundary conditions in all directions) [18]. Instead, if viewed as open strings they would have mixed (that is, neither Dirichlet nor Neumann) boundary conditions. These states do not have world-sheet description. They can be viewed as arising from D-branes wrapping (collapsed) two cycles in the orbifold [18]. These states are clearly nonperturbative from the orientifold viewpoint.

This difficulty is a generic feature in most of the orientifolds of type IIB compactified on toroidal orbifolds, as well as the corresponding noncompact cases such as the ΩJ orientifolds of type IIB on $\mathbf{C}^3/\Gamma$. However, there is a (rather limited) class of cases where the would-be nonperturbative states are massive (and decouple in the low energy effective field theory) if we consider compactifications on blown up orbifolds [18]. In fact, these blow-ups are forced by the orientifold consistency. The point is that the orientifold projection Ω must be chosen to be the same as in the case of type IIB on a smooth Calabi-Yau threefold. The reasons why this choice of the orientifold projection is forced have been recently discussed at length in [18]. In particular, we do not have an option of choosing the orientifold projection analogous to that in the six dimensional models of [30]. (Instead, the orientifold projection must be analogous to that in the $\mathbf{Z}_2$ models of [31].) On the other hand, the above Ω orientifold projection is not a symmetry of type IIB on $\mathbf{C}^3/\Gamma$ at the orbifold conformal field theory point [18]. The reason for this is that Ω correctly reverses the world-sheet orientation of world-sheet bosonic and fermionic oscillators and left- and right-moving momenta, but fails to do the same with the *twisted* ground states. (Such a reversal would involve mapping the g_a twisted ground states to the g_a^{-1} twisted ground states. In [18] such an orientation reversal was shown to be inconsistent.) This difficulty is circumvented by noting that the orientifold projection Ω is consistent for *smooth* Calabi-Yau threefolds, and, in particular, for a blown up version of the $\mathbf{C}^3/\Gamma$ orbifold. Thus, once the appropriate blow-ups are performed, the orientifold procedure is well defined.

In some cases the blow-ups result in decoupling of the would-be massless nonperturbative states, which is due to the presence of an appropriate superpotential (that couples the blow-up modes to the nonperturbative states). This feature, however, is not generic and is only present in a handful of cases. This was shown to be the case in Ω orientifolds of type IIB on T^6/Γ with $\Gamma \approx \mathbf{Z}_3$, $\mathbf{Z}_7$, $\mathbf{Z}_3 \otimes \mathbf{Z}_3$, $\Delta(3\cdot 3^2)$ in [20,21,15]. These orientifolds correspond to type I compactifications on the corresponding orbifolds, which in turn have perturbative heterotic duals [the corresponding heterotic compactifications are perturbative as there are no D5 branes (which would map to heterotic NS5 branes) in these models]. The nonperturbative (from the orientifold viewpoint) states were shown to correspond to twisted sector states on the heterotic side. The perturbative superpotentials for these states can be readily computed, and are precisely such that after the appropriate blow-ups (those needed for the orientifold consistency) the twisted sector states decouple.

These arguments were generalized in [18] to the $\mathbf{Z}_6$ model of [21] and the $\mathbf{Z}_2 \otimes \mathbf{Z}_2 \otimes \mathbf{Z}_3$ model of [25]. In all the other cases (except for the $\mathbf{Z}_2 \otimes \mathbf{Z}_2$ model of [23] which is obviously perturbative from the above viewpoint) it was argued in [18] that nonperturbative states do not decouple. Various checks of these statements were performed in [18] using the web of dualities between type IIB orientifolds, F-theory, and type I and heterotic compactifications on orbifolds. These statements, however, only depend on local properties of D branes and orientifold planes near orbifold singularities and should persist in noncompact cases such as the ΩJ orientifolds of type IIB on $\mathbf{C}^3/\Gamma$. In [14] it was shown that only for the above seven choices of the orbifold group do the perturbative (from the orientifold viewpoint) tadpole cancellation conditions have appropriate solutions for the ΩJ orientifolds of Type IIB on $\mathbf{C}^3/\Gamma$.

The reason we have reviewed the above facts is the following. In the nonsupersymmetric cases we discussed in this section *a priori* we also might expect nonperturbative states arising in various sectors of the orientifold. Unlike in the supersymmetric cases, however, it is unclear whether such states would decouple once the orbifold singularities are resolved. In fact, it is not even clear if the corresponding ''blow-ups'' are marginal deformation (since supersymmetry is broken). In particular, we do not have the dual heterotic picture in these cases which we would use to check the decoupling of nonperturbative (from the orientifold viewpoint) states: such heterotic duals would be intrinsically nonperturbative (as they would contain NS 5 branes), and it is not clear how to proceed in these cases. Thus, we do not really have an independent check in the nonsupersymmetric cases (in contrast to the $\mathcal{N}=1$ supersymmetric cases) for the perturbative consistency of the noncompact orientifolds we constructed in this section. However, non-Abelian anomaly cancellation in these models is rather nontrivial, so it is reasonable to believe that these models are indeed perturbatively consistent. Yet, the above discussion points to a possible caveat in the above construction.

V. CONCLUSIONS

In this paper we have constructed nonsupersymmetric large *N* gauge theories from orientifolds. The construction is similar to that of the supersymmetric models but involves nontrivial discrete torsion which is the source of supersymmetry breaking. The nonchiral models we have constructed in this paper are consistent as they are obtained by including nontrivial discrete torsion in $\mathcal{N}=2$ theories (in which we do not expect nonperturbative states due to a peculiar orientifold projection). However, the situation with the chiral models is less clear: they are obtained by including nontrivial discrete torsion in $\mathcal{N}=1$ theories where *a priori* we do expect nonperturbative states. Unlike in the supersymmetric cases, in the nonsupersymmetric case decoupling of such states is far from being obvious. It would be interesting to understand this issue in more detail. However, assuming that such states do decouple, the chiral models (along with the nonchiral ones) we have constructed in this paper provide examples of nonsupersymmetric large *N* gauge theories from orientifolds with well defined world-sheet expansion which is in one-to-one correspondence with 't Hooft's large *N* expansion (and results in rather nontrivial statements about the corresponding gauge theories in the large *N* limit).

In conclusion we would like to stress that if we attempted to construct *compact* models using the above techniques, we would get tachyonic models in which the closed string sector (unlike in the large *N* limit of the noncompact cases) does not decouple from the gauge theory, and these models would be sick due to tachyonic instabilities. It remains an open question whether it is possible to construct chiral tachyon and tadpole (and, therefore, anomaly) free compact type I models. It would be interesting to understand this issue in more detail.

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