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ON CONDENSATION OF CLOSED-STRING TACHYONS

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ABSTRACT

An F-theory dual of a nonsupersymmetric orientifold is considered. It is argued that the condensation of both open and closed string tachyons in the orientifold corresponds to the annihilation of branes and anti-branes in the F-theory dual. The end-point of tachyon condensation is thus expected to be the vacuum of Type-IIB superstring. Some speculations are presented about the F-theory dual of the bosonic string and tachyon condensation thereof.

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1. Introduction and Summary

Much has been learnt recently about the condensation of open-string tachyons (see [1] for an overview). One of the main ingredients that has made it possible to successfully analyze this very off-shell process is the fact that even before doing any computation there is a compelling *a priori* argument about the end-point of tachyon condensation. The open-string tachyon in a brane-antibrane system has an interpretation in the closed-string channel as the instability towards brane-antibrane annihilation [2]. Therefore, the end-point of tachyon condensation is expected to be nothing but the closed string vacuum. This physical picture underlies the 'Sen conjecture' about the height of the tachyon potential [3] and has been a very valuable guide for the explicit computations of the tachyon potential using various off-shell formalisms.

One of the difficulties in applying similar ideas to the more interesting case of the closed-string tachyon is that there is no analogous dynamical picture for the condensation of closed-string tachyons that is immediately obvious.

In this paper, we address this issue for a nonsupersymmetric orientifold model in eight dimensions. Before presenting the details, let us summarize the main argument. The orientifold theory that we consider contains orientifold 7-planes as well as anti-orientifold 7-planes. Tadpoles are canceled by including both $D7$-branes and $\bar{D}7$-branes. The spectrum contains tachyons in certain regions of the moduli space from both the open-string and the closed-string sectors. The open-string tachyons indicate as above the instability towards the annihilation of $D7$-$\bar{D}7$ branes and arise when the separation between branes and anti-branes is sufficiently small. However, there is no analogous interpretation of the closed string tachyons. Moreover, the $O7$ and the $\bar{O}7$ planes both have negative tension and hence are not dynamical objects. In the perturbative orientifold theory they cannot get annihilated and will stay around even after the annihilation of $D7$-$\bar{D}7$ branes. Non-perturbatively, it is natural to consider an F-theory dual of this orientifold theory in which each $O7$-plane is replaced by a pair of (p, q) 7-branes. The $\bar{O}7$ plane is similarly replaced by a charge-conjugate pair of (p, q) 7-branes. Now, in the F-theory picture a (p, q) 7-brane is a dynamical object locally not any different form a $D7$-brane and can certainly get annihilated along with its charge conjugate $(-p, -q)$ 7-brane. Thus, a collection of (p, q)

branes and antibranes is unstable towards annihilation. We therefore interpret the various open and closed string tachyons of the orientifold as a signal of this instability. In these models, there are equal number of (p, q) 7-branes and anti 7-branes so they can annihilate each other completely emitting various closed string modes such as gravitons and dilatons. This process will leave behind the vacuum of the Type-IIB superstring. We are thus led to the conjecture that the condensation of open and closed tachyons of the orientifold leads to the supersymmetric Type-IIB string as the endpoint.

We present the details of the orientifold in §2 and discuss the proposed F-theory dual in §3. We conclude in §4 with some speculative remarks about the closed-string tachyons in the bosonic string.

2. A Non-supersymmetric Orientifold

Our starting point will be an orientifold of the Type-IIB string compactified on a 2-torus to eight dimensions by a $\mathbf{Z_2} \times \mathbf{Z_2}$ symmetry $\{1, \Omega(-1)^{F_L} I_{89}\} \times \{1, \Omega(-1)^{F_R} I_{89}\sigma\}$ where Ω is the usual orientation reversal on the worldsheet, I_{89} is the reflection of the 89 coordinates, σ is a half-shift along the 9th direction, and F_L and F_R are the left and right-moving spacetime fermion numbers respectively [4,5].

Let us first look at the spectrum in the closed-string sector. The massless states from the untwisted sector that survive the orientifold projections are the metric, the dilaton and the 2-form RR field. The orientifold group $\{1, \Omega(-1)^{F_L} I_{89}, \Omega(-1)^{F_R} I_{89}\sigma, (-1)^F\sigma\}$ contains the element $(-1)^F\sigma$ where $F = F_L + F_R$ is the total fermion number; hence the spectrum will contain states twisted by this element. In the twisted sector, the ground state energy in the light-cone Green-Schwarz formalism is $-\frac{8}{24} - \frac{8}{48} = -\frac{1}{2}$ for both left-movers and right-movers because eight fermions are half-integer moded and eight bosons are integer moded. If $(-1)^F$ were not accompanied by the half-shift σ around the circle along the 9th direction, this would lead to a neutral closed-string tachyon as well as an additional massless R-R 2-form. However, because of the half-shift, all these states have large positive mass-squared when the radius of the circle along the 9th direction is large. Only when this radius becomes smaller than the string-length does one see a tachyon in the spectrum. All closed-string massless fermions are projected out.

In the open-string sector, there are four orientifold 7-planes at the fixed points of $\Omega(-1)^{F_L} I_{89}$. Moreover, there are four anti-orientifold 7-planes at the fixed points of $\Omega(-1)^{F_R} I_{89}\sigma$ For simplicity, let us consider a square torus with the radii of the circles in the 8 and 9 directions to be R. Let us write $z = X^8 + iX^9$ as the complex coordinate of the torus. Then the orientifold planes are located at the fixed points of I_{89}:$z_1 = 0$: $z_2 = \frac{1}{2}$, $z'_1 = 0 + \frac{i}{2}$, and $z'_2 = \frac{1}{2} + \frac{i}{2}$ in units of R. Similarly, the anti-orientifold planes are located at $z_3 = 0 + \frac{i}{4}$, $z_4 = \frac{1}{2} + \frac{i}{4}$, $z'_3 = 0 + \frac{3i}{4}$ and $z'_4 = \frac{1}{2} + \frac{3i}{4}$ which are the fixed points of $I_{89}\sigma$. Note that z'_1 and z'_2 respectively are the images of z_1 and z_2 under $I_{89}\sigma$ whereas z'_3 and z'_4 respectively are the images of z_3 and z_4 under I_{89}. (Fig 1)

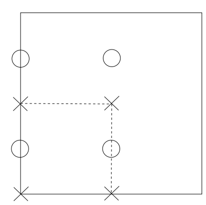

Fig. 1: The crosses indicate the orientifold planes and the circles indicate anti-orientifold planes. The fundamental region of the orientifold is the quarter of the original torus that is inside dashed lines.

The tadpoles can be canceled locally by putting eight elementary D7-branes at each of $O7$-planes and eight $\bar{D}7$-branes at each of the $\bar{O}7$ planes†. The spectrum contains $SO(8)^4$ gauge bosons and massless fermions in the adjoint reps of the gauge group. There are charged scalar multiplets that coming from the $7-\bar{7}$ open strings that become tachyonic when the separation between the branes and anti-branes is sufficiently small.

At tree-level, there are a number of massless moduli fields. We can talk about the classical moduli space even though these moduli will typically be lifted at loop level and moreover in the interacting theory the tachyons will in any case destabilize the vacuum. In the closed-string sector, there are two complex moduli that correspond to complex-structure and Kähler deformations and a real modulus that corresponds to the dilaton. The open-string sector contributes sixteen complex moduli corresponding to the positions

† By elementary D-branes we mean the D-branes of the original Type-IIB string before the orientifold projection. On the $\mathbf{Z_2} \times \mathbf{Z_2}$ orientifold we need four elementary branes to make a single dynamical D-brane that can move away from the orientifold planes.

of the sixteen dynamical D-branes on the torus. There are thus altogether eighteen complex moduli and one real modulus.

If we move the D-branes away from the orientifold planes then the gauge symmetry is broken to appropriate subgroups of $SO(8)^4$. However, in this case the tadpoles are not canceled locally even though the total charge on the torus is zero. As a result, the metric of the torus as well as the dilaton would vary as we move around on the torus.

Before discussing the F-theory dual, we would like to note that there are two other $\mathbf{Z_2} \times \mathbf{Z_2}$ orientifolds that have been considered in the literature in various contexts. The orientifold action is closely related and but the spectrum is distinct.

A. The orientifold group is $\{1, \Omega(-1)^{F_L} I_{89}\} \times \{1, \Omega(-1)^{F_R} I_{89}\}$. Here there is no half-shift in the orientifold action and as a result the there are four pairs of coincident $O7$-$\bar{O}7$ planes. The gauge symmetry $SO(8)^8$ is twice as large and there are tachyons at all points in the moduli-space. This is T-dual to the orientifold $\{1, \Omega\} \times \{1, \Omega(-1)^F\}$ which has been conjectured in [6] to be S-dual to the 26-dimensional bosonic string compactified to ten dimensions on the $Spin(32)$ lattice. We will comment upon this orientifold in §4.

B. The orientifold group is $\{1, \Omega(-1)^{F_L} I_{89}\} \times \{1, \Omega(-1)^{F_R} I_{89}\tilde{\sigma}\}$. The half-shift $\tilde{\sigma}$ is T-dual to σ[†]. The spinors no longer transform in the adjoints but rather in the bi-vector representations of $SO(8)^4$. Under T-duality along both the 8 and 9 directions $\Omega(-1)^{F_L} I_{89}$ goes to Ω and $\tilde{\sigma}$ goes to σ. The T-dual orientifold symmetry is equivalently $\{1, \Omega\} \times \{1, \Omega(-1)^F \sigma\}$. This is an orientifold of a variant of the 0B string considered in [7] and is conjectured in [8] to be dual to the $SO(16) \times SO(16)$ heterotic string [9,10,11].

3. Nonsupersymmetric F-theory Dual

The F-theory dual of our orientifold can be established reliably at the special points the moduli space with $SO(8)^4$ symmetry where the tadpoles are canceled locally. Our reasoning will be analogous to that used for establishing the supersymmetric duality between

† On a state $|m, n\rangle$ with quantized momentum m and winding n along the 9th direction, σ acts with a phase $(-1)^m$ whereas $\tilde{\sigma}$ acts with the phase $(-1)^n$.

Type-I and F-theory [12]. However, at a generic point in the moduli space as the branes move around, the F-theory dual will be valid only in the approximation that the branes and anti-branes are far apart from each other such that their interaction with each other can be ignored. The approximation and the regime of its validity will become clear below.

Our starting point will be an elliptic curve in the Weierstrass form

$$y^2 = x^3 + xf(z,\bar{z}) + g(z,\bar{z}) \tag{3.1}$$

where $x, y, z \in \mathbf{CP^1}$, $f(z,\bar{z})$ is a polynomial of combined degree eight and $g(z,\bar{z})$ is a function of combined degree twelve in $z, \bar{z}$. Unlike in the supersymmetric case [13,14], the curve depends on both the holomorphic and anti-holomorphic coordinates $(z,\bar{z})$ of the base. An explicit parametrization of these functions suitable for our purposes is given later. This equation defines an elliptic fibration over an $\mathbf{S^2}$ base coordinatized by $(z,\bar{z})$. The τ parameter of the elliptic fiber is determined by

$$j(\tau(z,\bar{z})) = \frac{4.(24f)^3}{27g^2 + 4f^3} \tag{3.2}$$

where $j(\tau)$ is the usual modular invariant function normalized so that $j(i) = (24)^3$. For large values of imaginary part of τ we have,

$$j(\tau) \sim e^{-2\pi i\tau}. \tag{3.3}$$

Now, we define the nonsupersymmetric F-theory as the compactification of Type-IIB theory on the non-Ricci-flat $\mathbf{S^2}$ base coordinatized by $(z,\bar{z})$. The axion-dilaton field as a function of the coordinates of the base is given by the modular parameter of the above curve:

$$a(z,\bar{z}) + ie^{-\phi(z,\bar{z})} \equiv \tau(z,\bar{z}), \tag{3.4}$$

where a is the R-R scalar and ϕ is the dilaton.

The fiber degenerates generically at 24 points where the discriminant

$$\Delta \equiv 4f^3 + 27g^2 \tag{3.5}$$

vanishes. There are twelve 'zeroes' and twelve 'anti-zeroes'. Near a zero $z = z_i$,

$$j(\tau(z,\bar{z}) \sim \frac{1}{z - z_i} \tag{3.6}$$

and therefore the axion-dilaton field is given by

$$\tau(z,\bar{z}) \sim \frac{1}{2\pi i} \ln(z - z_i). \tag{3.7}$$

When z goes around z_i, the monodromy of the axion-dilaton field is given by $\tau \to \tau + 1$, which means that there is a single 7-brane located at the zero that carries one unit of magnetic charge with respect to the axion field. The metric has a conical singularity with deficit angle $\pi/6$ as for a stringy cosmic string [15]. Similarly near an 'anti-zero' $\bar{z} = \bar{z}_j$, we have,

$$j(\tau(z,\bar{z}) \sim \frac{1}{\bar{z} - \bar{z}_j} \tag{3.8}$$

and

$$\tau(z,\bar{z}) \sim \frac{1}{2\pi i} \ln(\bar{z} - \bar{z}_j). \tag{3.9}$$

Thus, inclusion of anti-branes on the base necessarily requires that the axion-dilaton field has the anti-holomorphic dependence on $\bar{z}$[†]. We expect that to the F-theory defined using the equation (3.1)would be a good approximation so long as the parameters are chosen such that the zeroes and anti-zeroes are well separated. In this approximation, the total deficit angle of twelve branes and twelve anti-branes adds up to 4π, so the base is a sphere with Euler character two. Here, because we are dealing with a non-supersymmetric situation, we talk about a real $\mathbf{S^2}$ and the Euler character instead of $\mathbf{CP^1}$ and first Chern class.

To make contact with the orientifold, we consider the theory at the special point in the moduli space with $SO(8)^4$ symmetry, where the tadpoles are canceled locally. In the Type-IIB language, this implies that the axion-dilaton field, τ, is a constant on the torus

[†] Note that $\tau(z,\bar{z}) \sim -\frac{1}{2\pi i} \ln(z - z_i)$ also has the monodromy appropriate for an anti-brane, however, in this case the imaginary part of τ is no longer positive definite as required physically from the definition of τ.

and the metric is flat except for four conical singularities at the orientifold fixed points with conical deficit angle of π at each singularity. To obtain the F-theory dual we choose

$$\begin{aligned} f(z,\bar{z}) &= \alpha\Phi(z,\bar{z})^2, \quad g(z,\bar{z}) = \Phi(z,\bar{z})^3, \\ \Phi(z,\bar{z}) &= (z-z_1)(z-z_2)(\bar{z}-\bar{z}_3)(\bar{z}-\bar{z}_4) \end{aligned} \tag{3.10}$$

where z_i's are the locations of the fixed points as defined in the previous section and α is a constant. The discriminant now (3.5) has two zeroes of order six at the points z_1, z_2 implying that there are six coincident 7-branes sitting at each of these points. Similarly, there are two anti-zeroes of order six at the points $\bar{z}_3, \bar{z}_4$ signifying six anti 7-branes at each of these points. Now, since $f^3/g^2 = \alpha^3$ is a constant, the j-function is a constant from (3.3), and therefore the τ field is a constant over the base; however, there is an $SL(2,\mathbf{Z})$ monodromy around each of the points $z_1, \ldots, z_4$

$$\begin{pmatrix} -1 & 0 \\ 0 & -1 \end{pmatrix}. \tag{3.11}$$

The metric on the base is given, in this limit, by

$$ds^2 = \frac{dz d\bar{z}}{\prod_i |(z-z_i)|} \tag{3.12}$$

up to over all normalization [15].

The metric is thus flat everywhere except at four points. At each of these points there is a deficit angle of π corresponding to a bunch of coincident six branes or six anti-branes. Note that the monodromy around two bunches of six branes is the square of the monodromy matrix in (3.11) which is identity. Thus, there is no net monodromy and as a result the configuration of twelve branes and twelve anti-branes can be patched together smoothly.

It is interesting to note that there is *no net force* between a bunch of twelve branes and a bunch of twelve anti-branes in this particular non-BPS configuration even though there is an attractive force between a single brane and a single anti-brane. This behavior is much better than what one might naively expect for an assembly of oppositely charged objects such as electrons and positrons which would experience an attractive force at any distance. It is possible because the charges of 7-branes in Type-IIB are non-abelian and as a result this nonsupersymmetric configuration where tadpoles cancel is locally stable. In particular, there is no tachyon all the way up to small values of the radius R. The F-theory description in terms of Type-IIB supergravity is of course valid only when the base of the compactification is large compared to the string scale.

The local analysis near the singularity is identical to the supersymmetric case. For example, near the point, $\bar{z} = \bar{z}_3$, the curve takes the form of a D_4 singularity

$$\tilde{y}^2 \sim \tilde{x}^3 + \tilde{x}\, \alpha \tilde{z}^2 + \tilde{z}^3 \tag{3.13}$$

after the rescaling

$$\tilde{y} = y\, \phi^{3/2}, \quad \tilde{x} = x\, \phi, \quad \tilde{z} = (z - z_1) \tag{3.14}$$

with

$$\phi = (z_1 - z_2)(\bar{z}_1 - \bar{z}_3)(\bar{z}_1 - \bar{z}_4). \tag{3.15}$$

This D_4 singularity corresponds to enhanced $SO(8)$ symmetry as in the supersymmetric case and we get $SO(8)^4$ symmetry from the four singular points $z_1, \ldots, z_4$.

Within F-theory, we can also consider more general configurations in which the D-branes are moved away from the orientifold plane. For example, locally, the deformation of (3.13) away from the D_4 singularity is given by the equation

$$y^2 = x^3 + x\, f(z) + g(z) \tag{3.16}$$

where f and g are now polynomials of degree two and three respectively. To begin with, the equation will depend on the seven complex parameters of the polynomials of which one can be removed by a shift of $\bar{z}$ and one more by rescaling x and y. The remaining five parameters are related, as explained in [12], to the parameters of Seiberg-Witten curve for the $N = 2$ supersymmetric $SU(2)$ gauge theory with four massive quarks [16]. The five parameters in the Seiberg-Witten theory are the four quark masses $\{m_i\}$ and the gauge coupling constant τ_0. The Seiberg-Witten theory corresponding the F-theory singularity has a physical interpretation as the world-volume theory of a D3-brane probe near the singularity [17]. In this correspondence, the positions of the four D-branes near the orientifold plane are given by the squares of the four masses and the asymptotic value

τ_0 of the axion-dilaton field equals the gauge coupling constant. Using this parametrization in terms of the Seiberg-Witten curve [16], the equation (3.16) can be recast as

$$y^2 = x^3 + x\alpha(\tau_0)\tilde{f}(z, m_i, \tau_0) + \tilde{g}(z, m_i, \tau_0), \tag{3.17}$$

where $\tilde{f}$ and $\tilde{g}$ are as before polynomial of degree two and three respectively but the coefficients of the leading power of z are chosen equal to one. In this picture, the orientifold plane splits into two planes which corresponds on the probe world-volume theory to the splitting of the origin with the $SU(2)$ classical symmetry into a 'monopole' and 'dyon' point. This splitting is non-perturbative in the string coupling and is of order $\exp(i\pi\tau_0/2)$.

The analysis near an anti-zero is the charge conjugate of the above. For example, near $\bar{z} = \bar{z}_3$ we will have four anti D-branes and the anti-orientifold plane will split into a charge-conjugate pair corresponding to a 'anti-monopole' and 'anti-dyon' point on world-volume theory of the D3-brane probe.

In summary, locally we have two copies of the Seiberg-Witten theory and two copies of anti-Seiberg-Witten theory. These four can be patched together to write the full equation that describes the most general deformation:

$$y^2 = x^3 + x\,\alpha(\tau_0)F + G, \tag{3.18}$$

where

$$\begin{aligned} F &= \tilde{f}(z - z_1, m_i^1, \tau_0)\tilde{f}(z - z_2, m_i^2, \tau_0)\tilde{f}(\bar{z} - \bar{z}_3, m_i^3, \tau_0)\tilde{f}(\bar{z} - \bar{z}_4, m_i^4, \tau_0) \\ G &= \tilde{g}(z - z_1, m_i^1, \tau_0)\tilde{g}(z - z_2, m_i^2, \tau_0)\tilde{g}(\bar{z} - \bar{z}_3, m_i^3, \tau_0)\tilde{g}(\bar{z} - \bar{z}_4, m_i^4, \tau_0) \end{aligned} \tag{3.19}$$

where $\tilde{f}$ and $\tilde{g}$ are the function appearing in equation (3.17). The overall scale of the base ρ is a real parameter that is arbitrary and not determined by this equation. Therefore, by conformal transformation of the $\mathbf{S^2}$ base, the solution depends only on the cross-ratio

$$\omega = \frac{(z_1 - z_2)(z_3 - z_4)}{(z_1 - z_3)(z_2 - z_4)}. \tag{3.20}$$

The solution is thus characterized by the sixteen complex mass parameters $\{m_i^s\}$ ($i = 1, \ldots, 4; s = 1, \ldots, 4$), the asymptotic value of coupling τ_0, the cross ratio ω and the scale

ρ of the base. This matches with the moduli of the orientifold theory that we have described in §2.

Let us recapitulate our logic behind using the F-theory description. We have established that at the special point in the classical moduli space with $SO(8)^4$ gauge symmetry, the orientifold theory is equivalent to the proposed F-theory. This equivalence will hold even if we change the shape of the torus or the string coupling in the orientifold theory as long as the tadpoles cancel locally. The duality thus holds on a submanifold of the moduli space with five real dimensions where the theory has $SO(8)^4$ symmetry. In this case, the Type-IIB supergravity equations are exactly satisfied everywhere on the base except at the four points where the fiber degenerates giving us a a reliable F-theory solution. One can now turn on massless deformations to move around the moduli space away from the submanifold and establish the duality in a large patch of the moduli space where the tadpoles no longer cancel locally. The branes are allowed to wander far away from the orientifold planes. The solutions described in this paper will be a good approximation as long as the branes are well-separated from the anti-branes. When the branes come close to the anti-branes, various tachyonic instabilities will kick in and the time-independent solution that we have been discussing in terms of the elliptic fibration will no longer be a good approximation.

Even within this approximation, the F-theory approach is more powerful than the orientifold for studying certain aspects of the theory. For example, as in [14] we can choose the parameters in such a way that the solution develops two E_8 singularities—one from coincident branes and the other from coincident anti-branes. We can take the singularities to be well-separated to work within our approximation. Thus, the theory will have $E_8 \times E_8$ symmetry—something that is difficult to see in the orientifold picture.

As we move some of the branes towards the anti-branes they will get annihilated. Even within the orientifold, all the D-branes can be annihilated along with the $\bar{D}7$-branes. We also expect that the split $O7$ planes will annihilate the split $\bar{O}7$-planes and the final configuration would have no charge left behind. The early stages of this annihilation process could be analyzed systematically up to a point using time-dependent solutions of the Type-IIB supergravity describing motion of branes and anti-branes. However, the final stage of this process is expected to be quite violent and will excite a lot of closed string modes. If we maintain translational invariance along the branes as they are annihilated, then the energy of annihilation will get deposited onto the base. Because the base is compact, the energy cannot be dissipated to asymptotic infinity. However, this will generate a potential for the moduli that control the size of the base and we expect that the theory will decompactify to restore the supersymmetric vacuum of Type-IIB superstring.

As in the supersymmetric case, we expect that the F-theory compactified further on a 2-torus should give a Type-IIA compactification. To see this explicitly, we compactify the 6th and 7th direction on circles and perform two T-dualities[†]. This gives an orientifold of Type-IIB with the orientifold symmetry $\{1, I_{6789}\Omega\} \times \{1, I_{6789}\Omega(-1)^F\sigma\}$ which contains $D5$-branes, $O5$-planes and similarly and $\bar{O}5$-branes and $\bar{O}5$-planes. Now, we can perform an S-duality transformation which turns Ω into $(-1)^{F_L}$ and we get an orbifold $\{1, I_{6789}(-1)^{F_L}\} \times \{1, I_{6789}(-1)^{F_R}\sigma\}$. Under this duality the $D5$-branes turns into $NS5$-branes and the orientifold planes would turn into orbifold planes with NS charge. Tachyons in the open string sector of the orientifold when a $D5$ brane comes close to $\bar{D}5$ brane correspond in the S-dual orbifold to a *closed* string tachyon when an $NS5$ to to an anti-$NS5$ brane approaching each-other. A further T-duality along the 6th direction will give us a non-supersymmetric Type-IIA compactification.

4. Conclusions and Speculations

The main conclusion of this paper is that certain nonsupersymmetric orientifolds with tachyons can be viewed as an assembly of dynamical charged objects with vanishing net charge. This orientifold can be analyzed reliably as an F-theory compactification in certain regions of the moduli space and can be viewed as a non-supersymmetric unstable state in the Type-IIB superstring sitting at an extremum of energy in the configuration space. There are a number of unstable directions at this extremum that correspond to closed and open string tachyons. From general principles, we expect this assembly to pair-annihilate completely to relax to the vacuum containing no charges. A detailed study

[†] See for example [18] for the T-duality transformation of Ω etc.

of this dynamical process is of course quite complicated but fortunately not required for determining the final state. Thus, the endpoint of tachyon condensation is expected to be the vacuum of Type-IIB superstring. This is not any different from a system of equal number of electrons and positrons—the detailed dynamics is complicated to analyze even in QED but one can be sure that the system would decay to the vacuum after complete pair-annihilation. One difference for a system of 7-branes is that the charges are non-abelian, which makes it possible to have configurations where the net force is zero between bunches of charges and anti-charges, but this does not alter the conclusion about the final state.

Let us compare this situation with the analysis of open-string tachyons in the D-$\bar{D}$ system. Within the open-string theory, there is no dynamical picture of what happens when the tachyons condense. It is only from the viewpoint of the closed-string channel that the D-branes are dynamical objects. This interpretation leads to the conjecture about the endpoint of tachyon condensation which can then be tested using open-string field theory. Similarly, here, within the orientifold theory, there is no obvious dynamical interpretation of the closed string tachyons. It is only in the F-theory dual that orientifold planes are dynamical objects and this interpretation suggests a conjecture about the endpoint of tachyon condensation in the orientifold theory. The analysis of the F-theory dual is based on the supergravity equations which are the low energy limit of the full closed string field theory. To study the dynamical process of tachyon condensation would require a time-dependent F-theory like solution of the Type-IIB supergravity equations and moreover in the regions of interest the α' corrections will become important. It would be interesting to see if closed string field theory or a variant of the methods utilized in [19] for studying twisted sector tachyons can shed more light on this process.

We end with some speculative remarks about the A-orientifold mentioned in §2 which was part of the motivation for the present work. The action of our orientifold was cleverly chosen so that the $O7$-planes were separated from the $\bar{O}7$-planes. We could reliably study the F-theory dual when this separation was large and establish that the orientifold plane would split into two dynamical 7-branes using local analysis. In this regime there are no tachyons in the spectrum. By contrast, for the A-orientifold with orientifold symmetry

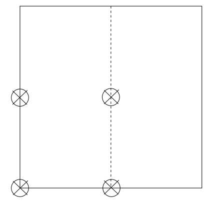

Fig. 2: The crosses indicating the orientifold planes and the circles indicating the anti-orientifold planes now coincide. The fundamental region of the orientifold is half of the original torus that is inside dashed lines.

$\{1, \Omega(-1)^{F_L} I_{89}\} \times \{1, \Omega(-1)^{F_R} I_{89}\}$, the $O7$ planes are coincident with the $\bar{O}7$-planes and the closed-string tachyons are always present (Fig 2).

There is no systematic way to analyze this situation in F-theory without dealing with tachyons but it is natural to speculate that even in this case the $O7$ and the $\bar{O}7$ planes would split nonperturbatively and would have the interpretation of dynamical objects. It is then reasonable to assume that this system will also dynamically pair-annihilate completely into the Type-II vacuum. Indeed, a comparison of Figs 1 and 2 suggests that we may be able to view the A-orientifold as a double cover of our F-theory with half-branes stuck at the fixed points. Note that the T-dual of the A-orientifold has been conjectured to be S-dual to the bosonic string. There are also intriguing relations between the bosonic string and the superstring [20] that underlie the 'bosonic map'[21]. If the A-orientifold is indeed dual to the bosonic string, then it leads to a speculation that the 26-dimensional bosonic string decays to the Type-II superstring. At this stage, however, we do not have more definitive evidence to offer.

There are a number of other recent conjectures about the fate of related non-supersymmetric theories such as the 0B theory, its orientifolds, and heterotic duals [22,23,24]. These are based on the interpretation of the 0B theory in terms of the Melvin solution of the Type-IIB string. Thus, also from this very different point of view, the 0B theory and possibly its orientifolds appear as excited configurations of the Type-IIB string which are then expected to decay to the supersymmetric vacuum. It would be interesting to see how these ideas connect with the those put forward in this paper.

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