

New Nonsupersymmetric Tachyon-Free Strings

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In four decades of string theory research, only a handful of nonsupersymmetric tachyon-free strings with only one neutral scalar at tree level were found. We construct new nonsupersymmetric tachyon-free string theories using asymmetric orbifolds that serve as the lower-dimensional counterparts to the $O(16) \times O(16)$ string in 4, 6, and 8 dimensions, each featuring only one neutral scalar at tree level, chiral matter and positive leading order potential. The 4D construction uses a quasicrystalline orbifold.

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Introduction—The goal of any theory of quantum gravity is to provide realistic models that could describe both our cosmological history and the standard model physics. In particular, one would aim for nonsupersymmetric vacua that have no moduli at low energies, positive dark energy, chiral matter, and the standard model gauge group.

A particularly difficult endeavor in string theory is finding nonsupersymmetric tachyon-free stable vacua. In the past decades of search, only three nonsupersymmetric tachyon-free strings have been found in 10 dimensions: the $O(16) \times O(16)$ heterotic string [1,2], the $USp(32)$ string [3] and the type 0'B string [4,5], with all having positive vacuum energy at leading order. The first two are almost rigid in the sense that their only massless scalar is the dilaton, but the $USp(32)$ string also has a gravitino and realizes non-linear supersymmetry. Furthermore, the 0'B theory has two neutral scalars. Therefore, the $O(16) \times O(16)$ string is the only tachyon-free nonsupersymmetric string with one neutral scalar and no gravitinos. This feature persists for the toroidal compactifications of the $O(16) \times O(16)$ string at special moduli with maximal non-Abelian enhancement. However, the additional toroidal moduli will in general develop a mass or become tachyonic at one loop [6]. There has been further study of the $O(16) \times O(16)$ vacua recently [7] with new constructions including flux compactifications to AdS_3 [8].

Finding more nonsupersymmetric string theories is crucial for diversifying our nonsupersymmetric string lab, which, until now, consisted of only three healthy models

and their compactifications. Such strings should be tachyon-free in order to be stable at tree level, and rigid, where the only neutral scalar is the string coupling, which is unavoidable in perturbative string theory. The $O(16) \times O(16)$ theory in particular is the only closed string example that has these features. In this Letter, we provide the sought-after analogs of the $O(16) \times O(16)$ string directly in 4, 6, and 8 dimensions with no supersymmetry, no tachyons, one neutral massless scalar, together with additional massless charged fields at tree level. At one loop the neutral scalar develops a running vacuum energy and the charged matter fields typically pick up mass or become tachyonic.

Interestingly, all nonsupersymmetric models with one neutral modulus have positive leading order potential. This is the case for all the models we present in this Letter as is for the other three 10D string theories, including the 10D $O(16) \times O(16)$ string and its compactification on a circle at the point that there are no tachyons [7]. This seems to suggest that any “rigid” tachyon-free string theory necessarily has positive cosmological constant and chiral fermions providing a potential naturalness explanation for our positive vacuum energy and the existence of the standard model fermions.

The Letter is organized as follows. First we review strings on orbifolds and set our notation. In particular, we will use asymmetric orbifolds [9] and the more special *quasicrystalline* orbifolds [10], recently studied in [11]. Details of the lattice constructions used are summarized in Appendix B of Supplemental Material [12]. Next we describe the new nonsupersymmetric tachyon-free rigid strings in 4, 6, and 8 dimensions obtained as asymmetric orbifolds of the heterotic string. In Appendix C of [12] we present the computation of the 1-loop potentials. All the models have chiral fermions and are subject to anomaly cancellation. In Appendix D of [12] (see also Refs. [13–25] therein) we summarize the anomaly cancellation methods used in each model.

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The construction we present in this Letter is particularly important, as finding such models is extremely rare. Achieving a tachyon-free and rigid spectrum requires significant fine-tuning and the systematic exploration of a large number of candidate models.

Orbifolds review—In this section we briefly review the important aspects of heterotic compactifications on tori and their orbifolds.

Torus compactifications: Heterotic string theory compactified on a d -dimensional torus T^d has ground state momenta $(p_L; p_R)$ taking values in an even self-dual lattice $\Gamma^{d+16;d} \subset \mathbb{R}^{d+16;d}$, called the *Narain lattice* [26,27]. Torus compactifications and background field data are completely characterized by the choice of the Narain lattice.

Lattice symmetries: At special points in the moduli space, some automorphisms of the Narain lattice act as rotations on the left and right movers, without mixing the two. They form the *Narain symmetry group*

$$\begin{aligned} \text{Sym}(\Gamma^{d+16;d}) &:= \text{Aut}(\Gamma^{d+16;d}) \\ &\cap (\text{O}(d+16, \mathbb{R}) \times \text{O}(d, \mathbb{R})). \end{aligned} \quad (1)$$

Such automorphisms act as symmetries on the worldsheet CFT, with T-duality at the self dual radius being the ur-example. Since these Narain symmetries θ are rotations, they have eigenvalues $e^{\pm 2\pi i \phi_L^i}$ on the left and $e^{\pm 2\pi i \phi_R^i}$ on the right and order $N \in \mathbb{Z}$. Therefore, we can equivalently characterize these rotations by their twist vectors

$$\begin{aligned} \phi &= \left(\phi_L^1, \dots, \phi_L^{\frac{d+8}{2}}; \phi_R^1, \dots, \phi_R^{\frac{d}{2}} \right), \quad N\phi_L^i \in \mathbb{Z}, \\ N\phi_R^i &\in 2\mathbb{Z}, \end{aligned} \quad (2)$$

where the $2\mathbb{Z}$ value specifies the $\text{Spin}(d)$ uplift on the right.

A Narain symmetry θ is a *quasicrystallographic symmetry* if it acts crystallographically on $\Gamma^{d+16;d}$ but not separately on the left and right [10]. In contrast, *crystallographic symmetries* have separate crystallographic actions on the left and on the right.

For example, there is no 2D lattice with a $\mathbb{Z}_5$ rotational symmetry due to the crystallographic restriction theorem. However, there is a $\mathbb{Z}_5$ action on the indefinite lattice $\Gamma_5^{2;2}$ constructed in Appendix B of Supplemental Material [12], whose left and right movers form 2D quasicrystals with $\mathbb{Z}_5$ rotational symmetry. We present an in depth study of quasicrystallographic symmetries in [11].

In addition to twists θ , one can also act by a shift vector $v = (v_L; v_R)$ in the span of the Narain lattice $\Gamma^{d+16;d}$. The shift vector on the left and right movers are denoted by v_L and v_R , respectively.

Orbifolds: String theory on orbifold backgrounds was introduced in [28]. Geometrically, *orbifolds* can be constructed from a torus T^d by quotienting by a cyclic group $\mathbb{Z}_N = \langle g \rangle$ generated by isometry g as $T^d/\mathbb{Z}_N$. On the

worldsheet, the orbifolding procedure amounts to relaxing the boundary conditions of the string

$$X^i(\tau, \sigma + 2\pi) = g^n \cdot X^i(\tau, \sigma) = \theta X^i(\tau, \sigma) + v^i, \quad (3)$$

and then projecting to the invariant subspace of $\hat{g}$

$$\hat{g} \cdot |\psi\rangle = |\psi\rangle. \quad (4)$$

The states with g^n -twisted boundary conditions make up the $\hat{g}^n$ -twisted sector. The $n = 0$ sector corresponds to the *untwisted sector*. Note that by CPT, $\hat{g}^n$ and $\hat{g}^{N-n}$ -twisted sectors contain antiparticles of each other.

For geometric orbifolds, the action of g on the Narain lattice is left-right symmetric as $\theta_L = \theta_R$ and $v_L = v_R$. Since the left and right degrees of freedom of strings are decoupled, one can generalize the orbifolding procedure to left-right asymmetric actions g with $\theta_L \neq \theta_R$ or $v_L \neq v_R$. The procedure is carried out on the worldsheet CFT in a similar fashion by relaxing the left and right boundary conditions by $g = (\theta_L, v_L; \theta_R, v_R)$

$$X_{L,R}^i(\tau, \sigma + 2\pi) = \theta_{L,R} \cdot X_{L,R}^i(\tau, \sigma) + v_{L,R}^i, \quad (5)$$

and projecting to the invariant subspace of $\hat{g}$. Such orbifolds have no target space interpretation and are called *asymmetric orbifolds* [9,29]. But they can be related to geometric models as discussed in [30,31]. Details on the consistency of orbifolds and computation of the spectrum can be found in Appendix A of Supplemental Material [12].

Tachyon-free orbifolds—In this section, we construct the lower dimensional analogs of the $\text{O}(16) \times \text{O}(16)$ heterotic string.

We have two criteria: a nontachyonic spectrum, and the string coupling as the only massless neutral scalar at tree level. Note that the torus compactifications of the $\text{O}(16) \times \text{O}(16)$ string at enhanced gauge points fulfill these criteria—they are nontachyonic and the only massless neutral scalar at tree level is the dilaton [6]. Our main goal is to find heterotic orbifold examples beyond the $\text{O}(16) \times \text{O}(16)$ string paradigm.

There is no systematic way to ensure that a nonsupersymmetric orbifold will have a nontachyonic spectrum without carrying out the computations, but there is a way to ensure that there will be no neutral scalars. In particular, for geometric (i.e., symmetric) orbifolds $T^d/\mathbb{Z}_n$, there will always be neutral scalars associated with the moduli of the torus T^d . In contrast, asymmetric actions are only possible when the torus is string sized, therefore orbifolding at that point fixes the geometry and gets rid of neutral scalars.

To avoid tachyons on the other hand, one needs to get lucky—when there is a tachyonic left-mover state, there should be no level matching tachyonic right mover, and vice versa. We have found such an example in each even dimension. [The ground state energies $(E_0)_L, (E_0)_R$ in 2D

orbifolds are usually nontachyonic, so it is significantly easier to find such nonsupersymmetric models in 2D and we do not consider them here.] All the examples have chiral fermions, similar to the $O(16) \times O(16)$ string.

The reason we only have even dimensional examples is because a cyclic orbifold in odd dimensions would either have a fixed circle S^1 direction, hence an extra neutral scalar, or odd number of -1 eigenvalues in the twist, in which case the lift of the action to the worldsheet CFT would be more subtle [29,32]. It would be interesting if one can obtain similar strings in odd dimensions as well.

4D: In this section we present a 4D nonsupersymmetric heterotic orbifold with only one neutral complex scalar, given by the dilaton ϕ and the universal string axion $\star B$, i.e., the dual of the Kalb-Ramond field. All other scalars are charged and there are no tachyons in the twisted sectors.

In fact, this model has an anomalous $U(1)$ [14], which becomes massive at 2-loops [18] and “eats” the universal string axion to become its longitudinal mode. Details of the mechanism are summarized in Appendix D.1 of Supplemental Material [12]. As a result, the dilaton remains as the only neutral real scalar.

As is standard we present all Weyl fermions only as left handed, since a right-handed Weyl fermion can equivalently be presented in terms of its left handed CP partner. (A right-handed Weyl fermion in representation R is equivalent to a left-handed Weyl fermion in representation $\bar{R}$).

The orbifold construction is as follows: we choose the Narain lattice $\Gamma(E_8) \oplus L\Gamma_5^{2;2} \oplus \Gamma_5^{2;2}\Gamma_5^{2;2}$, where L is an eight-dimensional lattice with $\det L = 5$ and $\Gamma_5^{2;2}$ is a $\mathbb{Z}_5$ quasicrystal. Both of these lattices and their gluings are constructed explicitly in Appendix B of Supplemental Material [12].

For the orbifolding action, we twist by $\phi = (4, 4, 4, 0^8; 2, 2, 2)/5$ and shift in the invariant lattice $\Gamma(E_8) \oplus L$ by

$$v = (2, 0, 3, 0, 1, 4, 0, 0, 4, 3, 0, 3, 3, 3, 4, 0)/5. \quad (6)$$

The continuous gauge group is $SO(10) \times SU(5) \times SU(3) \times SU(2) \times U(1)^4$. (It is interesting to note that quasicrystalline compactifications are not gauge enhanced points, therefore the rank of the orbifold gauge group is lower than what one would have expected from the usual crystallographic orbifolds.) The massless spectrum is

$$G + 84 \times V + 561 \times F_c + 9 \times F_0 + 354 \times S_c + 1 \times S_0, \quad (7)$$

where G denotes the graviton, S denotes a complex scalar, and F denotes a left-handed Weyl fermion, and the subscripts 0 and c denote whether the state is neutral or charged with respect to the continuous gauge group. The representations for the scalars and fermions are given in Table I.

TABLE I. The charges of the scalars and fermions of the nonsupersymmetric 4D quasicrystalline $\mathbb{Z}_5$ orbifold. The spectrum is given in terms of complex scalars and left-handed Weyl fermions.

$SO(10) \times SU(5) \times SU(3) \times SU(2) \times U(1)^4$ reps		
Sector	Complex scalars	Left-handed Weyl fermions
Untwisted	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,0,0,0}$	$9(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,0,0,0}$
	$3(\mathbf{16}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,0,0,-15}$	$(\mathbf{16}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,0,0,-15}$
	$3(\mathbf{10}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{0,0,0,10}$	$(\mathbf{10}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{0,0,0,10}$
	$3(\mathbf{1}, \mathbf{5}, \mathbf{1}, \mathbf{1})_{-2,-10,-2,0}$	$(\mathbf{1}, \mathbf{5}, \mathbf{1}, \mathbf{1})_{-2,-10,-2,0}$
	$3(\mathbf{1}, \mathbf{\bar{10}}, \mathbf{1}, \mathbf{2})_{-1,5,-1,0}$	$(\mathbf{1}, \mathbf{\bar{10}}, \mathbf{1}, \mathbf{2})_{-1,5,-1,0}$
	$3(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{2})_{5,-5,5,0}$	$(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{2})_{5,-5,5,0}$
		$3(\mathbf{16}, \mathbf{1}, \mathbf{3}, \mathbf{1})_{0,0,0,5}$
		$3(\mathbf{1}, \mathbf{10}, \mathbf{1}, \mathbf{1})_{-4,0,-4,0}$
		$3(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{0,0,0,-20}$
		$3(\mathbf{1}, \mathbf{5}, \mathbf{1}, \mathbf{2})_{3,5,3,0}$
$\hat{g} + \hat{g}^4$		$3(\mathbf{1}, \mathbf{\bar{5}}, \mathbf{1}, \mathbf{1})_{2,-10,2,0}$
		$15(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{2})_{-1,-3,-1,12}$
		$15(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{-2,7,1,12}$
		$15(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,7,-3,12}$
		$5(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{2})_{1,-3,-5,12}$
		$5(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{2})_{-3,-3,3,12}$
		$5(\mathbf{1}, \mathbf{5}, \mathbf{1}, \mathbf{1})_{2,2,2,12}$
		$5(\mathbf{1}, \mathbf{\bar{5}}, \mathbf{1}, \mathbf{1})_{2,-3,-1,12}$
		$5(\mathbf{1}, \mathbf{\bar{5}}, \mathbf{1}, \mathbf{1})_{0,-3,3,12}$
		$5(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{2})_{-1,-3,-1,-8}$
$\hat{g}^2 + \hat{g}^3$		$5(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{0,7,-3,-8}$
		$5(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{-2,7,1,-8}$
	$15(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{4,-1,1,4}$	$5(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{4,-1,1,4}$
	$15(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{2,-1,5,4}$	$5(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{2,-1,5,4}$
	$15(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{-2,-6,-2,4}$	$5(\mathbf{1}, \mathbf{1}, \mathbf{\bar{3}}, \mathbf{1})_{-2,-6,-2,4}$

The gauge group is obtained as follows. The invariant lattice has two orthogonal components $\Gamma(E_8) \oplus L$, whose corresponding gauge group is broken by the shift as

$$E_8 \times [E_7 \times U(1)] \rightarrow [SO(10) \times SU(3) \times U(1)^3] \times [SU(5) \times SU(2) \times U(1)]. \quad (8)$$

The order of the $U(1)$ factors in Table I is the same as in (8). The Kac-Moody levels k_i of the $U(1)$ factors are 30, 100, 60, 300, respectively, which we use to normalize the $U(1)$ generators.

This model contains chiral charged fermions and hence anomalies can occur through nontrivial triangle diagrams with external gauge field or gravitons. All irreducible anomalies need to cancel identically but anomalies that can factorize as $X_2 Y_4$ can be canceled by the Green-Schwarz mechanism [13] with a counterterm of the form $B \wedge X_2$. This implies that the only fields that can be

anomalous are $U(1)$'s with $X_2 \propto F_{U(1)}$. Such $U(1)$'s are called “anomalous” and occur when $\text{Tr}(Q) \neq 0$ or when at least one of the external triangle legs contains that $U(1)$. All such possibilities including the details of the GS mechanism in 4 dimensions are summarized in Appendix D.1 of Supplemental Material [12]. The counterterm $B \wedge F_{U(1)}$ can be dualized to $\theta \partial_\mu A^\mu$ where θ is the universal axion dual to the Kalb-Ramond field. However, such a coupling necessarily includes a mass term for the gauge field with θ becoming its longitudinal mode. The mass term for the gauge field is produced at 2-loops in string theory [24]. Additionally, a potential for the charged scalar fields will be generated at 1-loop [18] and for the dilaton at 2-loops [25].

The leading order contribution to the canonically normalized dilaton potential in Einstein frame (The canonically normalized contribution from converting to Einstein frame is $e^{[(-D)/\sqrt{D-2}]\hat{\phi}}$ which exactly multiplies all 1-loop corrections.) is

$$V_{1\text{-loop}}(\hat{\phi}) \approx e^{-2\sqrt{2}\hat{\phi}}(3.13 \times 10^{-2})M_s^4, \quad (9)$$

which is positive.

Therefore, we have a very interesting example with standard model like gauge symmetry and chiral matter. The model has only one neutral scalar, the dilaton [The dual to the Kalb-Ramond 2-form is an axion and hence a scalar field which will also have a potential with a two-loop mass term and will become the longitudinal mode of the Abelian $U(1)$ (Appendix D.1 of Supplemental Material [12]).], for which there is a potential making it massive at higher order with a 2-loop contribution to the potential due to the massive $U(1)$, see Appendix D.1 of [12].

We can also compute the beta function for these gauge groups to understand their IR behavior. The computation of the one-loop beta function in Appendix D.1 of Supplemental Material [12] shows that the $SO(10)$ and $SU(5)$ factors are confined and all other gauge factors are IR free. All higher dimensional models are automatically IR free due to the gauge kinetic terms being irrelevant.

6D: We continue our search for tachyon-free models with one neutral scalar in higher dimensions. The next example is in 6 dimensions. All other scalars are complex and charged, and the spectrum is tachyon-free. In 6 dimensions, CP preserves the chirality and therefore we indicate the chirality when writing the spectrum.

Choose the Narain lattice $\Gamma(E_8 \times E_8) \oplus \Gamma^{4,4}(A_4)$. Twist by $\phi = (0^{10}; 2, 4)/5$ and shift in the invariant lattice $\Gamma(E_8 \times E_8) \oplus \Gamma(A_4)$ by

$$v = (3, 3, 1, 4, 4, 1, 2, 2, 4, 4, 1, 1, 2, 4, 2, 3, 3, 3, 2, 3)/5. \quad (10)$$

The gauge group corresponding to the invariant lattice $\Gamma(E_8 \times E_8) \times \Gamma(A_4)$ is broken by the shift as

$$E_8 \times E_8 \times SU(5) \rightarrow SU(5)^2 \times SU(5)^2 \times U(1)^4. \quad (11)$$

So the continuous gauge group is $SU(5)^4 \times U(1)^4$. The Kac-Moody levels of the $U(1)$ factors are 4, 60, 10, 30, which we use to normalize the $U(1)$'s to have level 1. The massless spectrum is

$$G + B + 100 \times V + 460 \times F_c + 460 \times \bar{F}_c \quad (12)$$

$$+ 460 \times S_c + 1 \times S_0. \quad (13)$$

Here, G and B denote the graviton and the Kalb-Ramond fields, S_0 denotes a real neutral scalar (the dilaton), S_c denotes a complex and charged scalar, and $F_c, \bar{F}_c$, respectively, denote charged left- or right-handed Weyl fermions. The representations for the charged spectrum are given in Table II.

The anomaly polynomial is given by

$$(2\pi)^4 I_8 = \frac{1}{12} X_4 Y_4 \quad (14)$$

with

$$X_4 = -15 \sum_{a=1,2} \text{Tr}(F_{SU(5)_a}^2) + 15 \sum_{a=3,4} \text{Tr}(F_{SU(5)_a}^2) + \sum_{i \geq j} \frac{1}{4} b_{ij} \frac{\text{Tr}(F_{U(1)_i})}{\sqrt{k_i}} \frac{\text{Tr}(F_{U(1)_j})}{\sqrt{k_j}}, \quad (15)$$

where

$$b_{ij} = \begin{pmatrix} 0 & 120 & 60 & 60 \\ 120 & 60 & 60 & -420 \\ 60 & 60 & -15 & 120 \\ 60 & -420 & 120 & 15 \end{pmatrix} \quad (16)$$

and

$$Y_4 = \sum_i \text{Tr}(F_{U(1)_i}^2) + 2 \sum_a \text{Tr}(F_a^2) - \text{Tr}(R^2), \quad (17)$$

where we have rescaled the $U(1)$'s to bring them in the canonical form Y_4 . The I_8 anomaly is cancelled with the GS counterterm $B \wedge X_4$.

The string frame cosmological constant for this theory is calculated in Appendix C of Supplemental Material [12] and gives the leading term in the dilaton potential in Einstein frame

$$V(\hat{\phi})_{1\text{-loop}} \approx e^{-3\hat{\phi}}(2.89 \times 10^{-3})M_s^6, \quad (18)$$

which is positive at leading order.

8D: Next we give a nonsupersymmetric and tachyon-free 8D orbifold model. Similarly, the only neutral scalar is

TABLE II. The charged spectrum of the nonsupersymmetric 6D $\mathbb{Z}_5$ orbifold. For each representation listed here, there is a complex scalar and a symplectic Majorana-Weyl fermion of the indicated chirality.

Sector	$SU(5) \times SU(5) \times SU(5) \times SU(5) \times U(1)^4$ reps	
Untwisted	$R(\mathbf{10}, \mathbf{5}, \mathbf{1}, \mathbf{1})_{0,0,0,0}$	$R(\bar{\mathbf{5}}, \mathbf{10}, \mathbf{1}, \mathbf{1})_{0,0,0,0}$
	$L(\mathbf{1}, \mathbf{1}, \mathbf{10}, \mathbf{5})_{0,0,0,0}$	$L(\mathbf{1}, \mathbf{1}, \mathbf{5}, \bar{\mathbf{10}})_{0,0,0,0}$
	$R(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{-1,3,2,6}$	$R(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,8,-3,1}$
	$R(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{1,-7,-3,1}$	$R(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{-2,-2,2,4}$
	$R(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{2,-2,2,-4}$	$L(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{1,1,4,2}$
	$L(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{0,-4,-1,7}$	$L(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{2,6,-1,-3}$
$\hat{g} + \hat{g}^4$	$L(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{-1,-9,-1,-3}$	$L(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{1})_{-2,6,-1,-3}$
	$R(\bar{\mathbf{5}}, \mathbf{1}, \mathbf{1}, \mathbf{5})_{0,-4,0,-2}$	
	$R(\bar{\mathbf{5}}, \mathbf{1}, \mathbf{1}, \mathbf{5})_{0,0,2,0}$	$R(\bar{\mathbf{5}}, \mathbf{1}, \mathbf{1}, \mathbf{5})_{1,1,-1,-1}$
$\hat{g}^2 + \hat{g}^3$	$R(\bar{\mathbf{5}}, \mathbf{1}, \mathbf{1}, \mathbf{5})_{-1,3,0,0}$	$R(\bar{\mathbf{5}}, \mathbf{1}, \mathbf{1}, \mathbf{5})_{0,0,-1,3}$
	$L(\mathbf{1}, \mathbf{5}, \mathbf{5}, \mathbf{1})_{1,1,1,-1}$	
	$L(\mathbf{1}, \mathbf{5}, \mathbf{5}, \mathbf{1})_{0,-4,-1,1}$	$L(\mathbf{1}, \mathbf{5}, \mathbf{5}, \mathbf{1})_{-1,-1,0,-2}$
	$L(\mathbf{1}, \mathbf{5}, \mathbf{5}, \mathbf{1})_{0,0,1,3}$	$L(\mathbf{1}, \mathbf{5}, \mathbf{5}, \mathbf{1})_{0,4,-1,-1}$

the dilaton with all other scalars charged. The situation of the fermions is similar to that in 4 dimensions, so we will present the spectrum only in terms of left-handed Weyl fermions.

Consider the Narain lattice $\Gamma(E_8 \times E_8) \oplus \Gamma^{2,2}(A_2)$ with a twist of the form $\phi = (0^9; 2/3)$ and a shift in the invariant lattice $\Gamma(E_8 \times E_8) \oplus \Gamma(A_2)$ given by

$$v = (2, 1, 0, 2, 2, 2, 1, 2, 0, 2, 0, 2, 2, 0, 2, 2, 1, 0)/3. \quad (19)$$

The gauge group corresponding to the invariant lattice $\Gamma(E_8 \times E_8) \oplus \Gamma(A_2)$ is broken by the shift as

$$E_8 \times E_8 \times SU(3) \rightarrow SU(9) \times SU(9) \times U(1)^2. \quad (20)$$

The full continuous gauge group is $SU(9)^2 \times U(1)^2$. The levels of the $U(1)$ factors are 6, 18. The total massless spectrum is

$$G + B + 162 \times V + 414 \times F_c + 414 \times S_c + 1 \times S_0. \quad (21)$$

Here, S_0 denotes a real neutral scalar, S_c denotes a complex and charged scalar, and F_c denotes a charged left-handed Weyl fermion. The representations for the charged spectrum are given in Table III.

Anomalies for this theory can be computed by 5-leg 1-loop diagrams with external gauge fields or gravitons. Such anomalies can be cancelled using the Green-Schwarz mechanism which is summarized in Appendix D.3 of Supplemental Material [12]. The anomaly polynomial is given by

 TABLE III. The charged spectrum of the nonsupersymmetric 8D $\mathbb{Z}_3$ orbifold. For each representation listed here, there is a complex scalar and a left-handed Weyl fermion (together with its CP conjugate).

Sector	$SU(9) \times SU(9) \times U(1)^2$ reps	
Untwisted	$(\mathbf{84}, \mathbf{1})_{0,0}$	$(\mathbf{1}, \mathbf{1})_{0,-6}$
	$(\mathbf{1}, \mathbf{84})_{0,00}$	$(\mathbf{1}, \mathbf{1})_{3,3}$
	$(\mathbf{1}, \mathbf{1})_{-3,3}$	
$\hat{g} + \hat{g}^2$	$(\mathbf{9}, \mathbf{9})_{-1,1}$	$(\mathbf{9}, \mathbf{9})_{0,-2}$
	$(\mathbf{9}, \mathbf{9})_{1,1}$	

$$(2\pi)^5 I_{10} = \frac{1}{12} X_6 Y_4 \quad (22)$$

with

$$X_6 = \frac{-1}{4} \text{Tr}(F_{U(1)_2}^3) + 3 \text{Tr}(F_a^3) + \frac{3}{4} \text{Tr}(F_{U(1)_2}) \text{Tr}(F_{U(1)_1}^2) \quad (23)$$

and

$$Y_4 = \text{Tr}(F_{U(1)_2}^2) + 2 \text{Tr}(F_a^2) - \text{Tr}(R^2) + \text{Tr}(F_{U(1)_1}^2), \quad (24)$$

where we have rescaled the $U(1)$'s to bring them in the canonical form Y_4 . The I_{10} anomaly is canceled with the GS counterterm $B \wedge X_6$.

In 8 dimensions there can also be global anomalies as discussed in [14] but we have an even number of fermions and hence the model is anomaly free. Similarly, to the Swampland observation in [33] the rank is even too and it would be interesting to investigate if similar arguments can be applied for these cases. Similar, generalized anomalies have been computed in [34] for nonsupersymmetric strings which could directly be applied in our models.

The global gauge anomaly [35,36] is also absent since $\pi_8[SU(9)] = 0$ and the [37] is absent since we do not have the adjoint of $SU(9)$.

The string frame cosmological constant for this theory is computed by the 1-loop dilaton tadpole computed in Appendix C of Supplemental Material [12] which in the Einstein frame gives the leading contribution to the dilaton potential:

$$V_{1\text{-loop}}(\hat{\phi}) \approx e^{\frac{-8}{\sqrt{6}}\hat{\phi}} (1.26 \times 10^{-4}) M_s^8, \quad (25)$$

which is positive. Therefore, all three models we have constructed have positive potential at leading order in 4, 6, and 8 dimensions.

Conclusion—We have extended the landscape of non-supersymmetric, tachyon-free string theories by introducing

novel theories in 4, 6, and 8 dimensions that mirror the characteristics of the $O(16) \times O(16)$ heterotic string. The only neutral scalar in each theory is the string coupling and each theory has a positive potential at leading order. It would be worthwhile to search for more such examples including in odd dimensions. Other nonsupersymmetric orbifold model searches include [23,38–42].

The theories we have constructed all come from orbifold models of the heterotic string. This implies that in 4 dimensions we only have one universal axion which sets the mass of the “anomalous” $U(1)$. In previous works [22] more general models have been constructed that can lower the mass scale of the $U(1)$ which can be more phenomenologically plausible but contain more moduli. It would be interesting to study minimal extensions of the $\mathcal{N} = 0$ standard model including a massive $U(1)$, as was done for the MSSM [43], to address a potential photon mass [44], relation to the $B - L$ symmetry [45], the dark sector [46,47] or other phenomenological implications [48,49]. Other attempts include [40,50,51].

Additionally, the rank of the three theories is even similar to the observations in supersymmetric theories [33] and in the rigid theories [31,52]. Additionally, we noted that the rank of the 4D theory did not increase as a consequence of the quasicrystalline compactification while the 8D and 6D theories had the usual toroidal rank increase as those exist at special points of gauge symmetry enhancement.

As has been discussed throughout this Letter, string theory does give rise to nonsupersymmetric vacua which can be related to supersymmetric vacua through a web of dualities [53–56]. It would be interesting to further explore such dualities.

The theories we present are free of tachyons at tree level. However, there are two different one-loop effects that can give mass to tree level massless states. Firstly, the massless charged scalars can pick up a tachyonic or nontachyonic mass at one-loop. This is the case for the $O(16) \times O(16)$ string on a circle S^1 at points with maximal gauge group enhancement [6,7]. But also the 1-loop scalar potential discussed in Appendix D.1 of Supplemental Material [12] will always add a contribution to the scalars charged under the $U(1)$ at 1-loop. The dilaton also acquires a mass term at 2-loops that is positive.

There is a way to cure these two instabilities for $O(16) \times O(16)$ and obtain a perturbatively stable nonsupersymmetric AdS [8]. By tuning to weak coupling, the tachyons can be made light enough to be uplifted by the gravitational potential of AdS [57], and by turning on flux, the dilaton gets stabilized. It would be worthwhile to attempt the same procedure for the new nonsupersymmetric strings presented here. However, such constructions generally will have non-perturbative instabilities [7,58].

An interesting observation is that all leading order potentials are positive for our models as well as for the 10D non-susy tachyon-free theories. This seems to suggest

that more rigid tachyon free string theories always have positive cosmological constant. This is very reminiscent of our expectation from the real world. It would be very interesting to construct a counterexample or find a general argument why this has to be the case. There are potential examples with more moduli which have either negative cosmological constant [50] or vanishing cosmological constant at one loop [59–61].

A more in-depth analysis is required for our models to determine their 1-loop fate in more general backgrounds. Especially in 4 dimensions, the potential term due to the anomalous $U(1)$ competes with the possibly tachyonic effects at 1-loop, which is interesting to study.

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