Kähler moduli stabilization in semirealistic magnetized orbifold models

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We study Kähler moduli stabilizations in semirealistic magnetized D-brane models based on $Z_2 \times Z_2'$ toroidal orbifolds. In type IIB compactifications, 3-form fluxes can stabilize the dilaton and complex structure moduli fields, but there remain some massless closed string moduli fields, Kähler moduli. The magnetic fluxes generate Fayet-Iliopoulos terms, which can fix ratios of Kähler moduli. On top of that, we consider D-brane instanton effects to stabilize them in concrete D-brane models and investigate the brane configurations to confirm that the moduli fields can be stabilized successfully. In this paper, we treat two types of D-brane models. One is based on D9-brane systems respecting the Pati-Salam model. The other is realized in a D7-brane system breaking the Pati-Salam gauge group. We find suitable configurations where the D-brane instantons can stabilize the moduli fields within both types of D-brane models, explaining an origin of a small constant term of the superpotential, which is a key ingredient for successful moduli stabilizations.

I. INTRODUCTION

Superstring theories are expected for an ultimate unified theory of particle physics including gravitational interactions. One of their remarkable features is that superstring theories are defined in ten-dimensional (10D) spacetime and predict the presence of extra dimensions of space for theoretical consistencies. We usually consider that the extra six-dimensional (6D) space is compactified in order to describe our Universe.

In such string compactifications, one of the challenging tasks is to realize a chiral spectrum in their four-dimensional (4D) effective theories, because they must be consistent with the standard model (SM) or some extensions such as the minimal supersymmetric standard model (MSSM). For this purpose, D-brane models are attractive because they can lead to various gauge groups with generations of chiral fermions [1–5], and several D-brane models were proposed realizing suitable 4D chiral spectra as zero modes of open strings on intersecting D-branes [6–9]. For the last decade, similar model building was actively attempted in the T-dual picture of the type IIA string theory, that is, in the framework of IIB strings with magnetized D-branes, and it was found that viable three-generation models can be obtained [10–12]. In particular, in the concrete model proposed in Ref. [13], a semirealistic flavor structure of the quarks and the leptons including their hierarchical masses and mixing angles was obtained, and furthermore, a spectrum of the supersymmetric particles and the Higgs bosons was calculated to verify its consistency with experimental results.

Another one of the key issues in the string compactifications is the stabilization of moduli fields that are massless

scalar modes originating from extra components of the higher-dimensional gravitational fields and n-form fields. Moduli stabilization is necessary to stabilize the extra compact space, and that is also significant in particle and cosmological phenomenologies. In these decades, several moduli stabilization mechanisms are proposed in the framework of superstring theories. We will discuss the moduli stabilization, concentrating on type IIB compactifications in this paper to associate them with magnetized D-brane models (moduli stabilizations with the magnetic fluxes were discussed in Refs. [14–16]). We find three types of dynamical variables to stabilize: the dilaton field, complex structure moduli, and Kähler moduli fields. Basically, in IIB string theories, we can introduce nontrivial fluxes for 3-form field strengths to stabilize the dilaton and complex structure moduli fields [17,18]. When the presence of the 3-form fluxes is turned on, however, the potential for the Kähler moduli stays flat at the tree level, and there remain some flat directions even when α'-corrections and string one-loop corrections are taken into account. We usually expect that those flat directions of Kähler moduli fields are stabilized by nonperturbative effects somehow.

In D-brane models, one of the computable nonperturbative effects is D-brane instantons [19–23], which we call Euclidean-branes (E-branes) in the present paper. These are D-branes that are localized at a point on 4D Minkowski spacetime but have a nonzero volume in the extra compact space. Thus, they can possibly yield a superpotential for the Kähler moduli and the dilaton field. Besides that, gaugino condensations of hidden D-branes also have computable nonperturbative effects to yield the superpotential of the moduli fields, but we will focus on the former in this paper.

In most of the previous works,[1] D-brane model building for the visible sector and for the moduli stabilization were discussed independently of each other. Such a scenario can be justified under the situation that the visible sector is irrelevant to the sector to stabilize moduli. For example, if the SM sector is localized at a certain point on the 6D compact space and the dynamics to stabilize moduli originates from the sector on cycles far away from the SM-localized point, those would be independent. However, if the SM sector and the moduli-stabilizing sector occupy a similar place in the 6D compact space, they would affect each other. Indeed, it is not trivial that the instanton effect yields a superpotential suitable for the moduli stabilization such as $W \sim Ae^{-aT}$, where T is the modulus, and that in practice depends on configurations of D-branes for the visible sector. This is due to the fact that one needs to integrate over the instanton zero modes to obtain nonperturbative superpotentials. We can realize the superpotential successfully when there is only a single E-brane wrapping $O(1)$-cycles without D-branes. On the other hand, in association with D-branes, there appear open string zero modes between the E-branes and the D-branes. When they cannot be soaked up by fermionic integration, the nonperturbative superpotential vanishes. Furthermore, even if zero modes are successfully soaked up, the superpotential including matter fields can be induced as $W \sim (\Phi_1 \Phi_2 \cdots)e^{-aT}$, but not the pure moduli term $W \sim Ae^{-aT}$. Such moduli-dependent terms with matter fields would be important in order to realize the right-handed Majorana neutrino masses and μ-terms of the Higgs fields in the MSSM [19–23,28,29]. However, such moduli-dependent terms with matter fields are not suitable for moduli stabilizations. We are thus required to study distributions of the zero modes for each brane configuration and confirm that no harmful fermionic zero modes remain to incorporate the moduli stabilizations with the D-brane models.

In this paper, we study the moduli stabilization due to the E-branes in association with concrete magnetized D-brane models for the visible sector in type IIB orientifolds. We assume the 3-form fluxes to stabilize the dilaton and the complex structure moduli fields preserving supersymmetry (SUSY), which allows us to concentrate on the Kähler moduli stabilization.[2] In those models, we will also turn on the "magnetic" fluxes for the world volume gauge field strength of the D-branes in order to realize the flavor structure of the SM. These magnetic fluxes classically produce moduli-dependent Fayet-Iliopoulos (FI) terms. We will find supersymmetric vacua with a certain ratio of the

vacuum expectation value (VEV) of the moduli fields, which means that the D-term potential can stabilize the Kähler moduli fields except for one flat direction. In order to stabilize the flat direction, we introduce E-branes and investigate the zero-mode structure in the D-brane models.

This paper is organized as follows. In Sec. II, we first review the magnetized $T^6/Z_2 \times Z_2'$ orbifolds in 10D supersymmetric Yang-Mills (SYM) theories, which correspond to the low-energy effective field theory of D9-brane systems, which explains an essence of magnetized orbifold models. Consequently, we propose several concrete models based on the Pati-Salam gauge group. Two types of E-branes are possible to give stable brane configurations in association with D9-branes. In the rest of the section, we study both the instanton effects to find several brane configurations with which the instanton effects work successfully and the induced superpotential that stabilizes the moduli field. In Sec. III, we perform a similar analysis with D7-brane models of the visible sector where the Pati-Salam gauge group is broken by the magnetic fluxes to realize a more realistic spectrum. Section IV is devoted to conclusions and discussions. In the Appendix, we discuss the zero-mode structure in the T-dual picture.

II. D9-BRANE MODELS

We study mixed configurations of magnetized D-branes and E-branes to construct models with all the moduli fields stabilized. In this section, we focus on Pati-Salam models based on a stack of eight D9-branes as the SM sector. These are the simplest but still semirealistic magnetized orbifold models. First we briefly review the 10D SYM theories compactified on a magnetized orbifold, which are low-energy effective field theories of magnetized D9-branes. In the theories, we can find several semirealistic models based on the Pati-Salam gauge group. Finally, we will investigate an E-brane's effects on the D9-brane systems, which generate a nonperturbative superpotential and stabilize the moduli fields. Note that any configuration of E-branes can appear and we have to take into account all possible E-branes. Some of them have no effects in low-energy effective field theory, but a certain E-brane can have nonperturbative moduli terms such as $W \sim Ae^{-aT}$. We are interested in such E-brane effects.

A. Review of magnetized orbifolds in 10D SYM theories

We give an overview of magnetized orbifolds in 10D SYM theories. In this paper, we consider three 2-tori, $T^2 \times T^2 \times T^2$, as an extra compact space, denoting their coordinates by z_i and $\bar{z}_i$ ($i = 1, 2, 3$). The 10D SYM theories can be described in the formulation of 4D $\mathcal{N} = 1$ superspace, focusing on a 4D $\mathcal{N} = 1$ SUSY out of full $\mathcal{N} = 4$ SUSY of the 10D SYM theories [31]. This was developed in compactifications of $T^2 \times T^2 \times T^2$ with

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magnetic fluxes in Ref. [32]. 10D SYM theories consist of 10D vector and Majorana-Weyl spinor fields, which are decomposed into 4D vector, complex scalar, and Weyl spinor fields. These 4D fields form 4D $\mathcal{N} = 1$ super-multiplets. As a result, field contents of the theories are expressed by a vector superfield V and three chiral super-fields Φ_i. Note that they are in adjoint representations of gauge symmetry of the SYM theories. In the following, we consider $U(N)$ SYM theories as effective field theories of one stack of N D9-branes.

We introduce Abelian magnetic fluxes in the $U(N)$ theories, which are parametrized by $N \times N$ diagonal matrices as

$$M^{(i)} = \text{diag}(m_1^{(i)}, m_2^{(i)}, \ldots, m_N^{(i)}),$$

where i runs over 1,2,3 corresponding to three T^2. When $m_n^{(i)}$ takes nondegenerate values, the $U(N)$ gauge group is broken down. For example, suppose the simplest case as follows:

$$M^{(i)} = \text{diag}(m_a^{(i)}, \ldots, m_a^{(i)}, m_b^{(i)}, \ldots, m_b^{(i)}), \quad (2.1)$$

where $m_a^{(i)} \neq m_b^{(i)}$. Then, these magnetic fluxes break the gauge group as $U(N) \rightarrow U(N_a) \times U(N_b)$. In this gauge symmetry breaking, we express the superfields as

$$\Phi_i \rightarrow \begin{pmatrix} \Phi_i^{aa} & \Phi_i^{ab} \\ \Phi_i^{ba} & \Phi_i^{bb} \end{pmatrix}, \quad (2.2)$$

where diagonal and off-diagonal entries are in adjoint and bifundamental representations of the unbroken gauge group $U(N_a) \times U(N_b)$, respectively. On this magnetized background, zero-mode equations for Φ_j^{ab} on the ith T^2 are given by

$$\left[\bar{\partial}_{\bar{i}} + \frac{\pi}{2\tau_i} (m_a^{(i)} - m_b^{(i)}) z_i \right] \Phi_j^{ab} = 0 \quad (\text{for } i = j), \quad (2.3)$$

$$\left[\partial_i - \frac{\pi}{2\tau_i} (m_a^{(i)} - m_b^{(i)}) \bar{z}_{\bar{i}} \right] \Phi_j^{ab} = 0 \quad (\text{for } i \neq j), \quad (2.4)$$

where τ_i is a complex structure of the ith T^2. For $i = j$, that has $m_a^{(i)} - m_b^{(i)}$ degenerate zero modes when $m_a^{(i)} - m_b^{(i)}$ is positive, while its conjugate one Φ_j^{ba} has no zero modes because $m_b^{(i)} - m_a^{(i)} < 0$. Thus, the magnetic fluxes produce generations of chiral fermions in 4D effective theories. This is almost the same for $i \neq j$, except for a relative sign in Eq. (2.4), and $|m_a^{(i)} - m_b^{(i)}|$ degenerate zero modes are produced for Φ_j^{ab} when $m_a^{(i)} - m_b^{(i)}$ is negative.

Next we study Z_2 orbifolding in this magnetized SYM theories. Let us consider Z_2 orbifolding, which acts on the first and second T^2, that is,

TABLE I. The number of active zero modes on the magnetized orbifold, where M represents an effective magnetic flux [that corresponds to $m_a^{(i)} - m_b^{(i)}$ in Eqs. (2.3) and (2.4)].

$|M|$	0	1	2	3	4	5	$2n$	$2n+1$
Even	1	1	2	2	3	3	$n+1$	$n+1$
Odd	0	0	0	1	1	2	$n-1$	n

$$(z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3).$$

On this orbifold, the superfields have to transform as

$$V(z_1, z_2, z_3) \rightarrow +PV(-z_1, -z_2, z_3)P^{-1},$$
$$\Phi_1(z_1, z_2, z_3) \rightarrow -P\Phi_1(-z_1, -z_2, z_3)P^{-1},$$
$$\Phi_2(z_1, z_2, z_3) \rightarrow -P\Phi_2(-z_1, -z_2, z_3)P^{-1},$$
$$\Phi_3(z_1, z_2, z_3) \rightarrow +P\Phi_3(-z_1, -z_2, z_3)P^{-1}, \quad (2.5)$$

where the projection operator P is an $N \times N$ matrix satisfying $P^2 = \mathbf{1}$. In accordance with these transformation laws, each entry of Eq. (2.2) is assigned to either a Z_2 even or odd mode. This Z_2 projection reduces the number of degenerate zero modes induced by the magnetic fluxes, as shown in Table I [33]. We can also introduce discrete Wilson lines, and the number of zero modes depends on the values of the discrete Wilson lines [34]. Here, for simplicity, we restrict ourselves to models with vanishing Wilson lines.

It is most important that the Abelian magnetic fluxes generically induce the FI term for trivial $U(1)$ parts of unbroken gauge subgroups. For instance, in the case of Eq. (2.1), there appear the FI terms with the following parameters in diagonal parts $U(1)_a \times U(1)_b$ of $U(N_a)$ and $U(N_b)$,

$$\xi_a = \frac{1}{\mathcal{A}^{(1)}} m_a^{(1)} + \frac{1}{\mathcal{A}^{(2)}} m_a^{(2)} + \frac{1}{\mathcal{A}^{(3)}} m_a^{(3)},$$
$$\xi_b = \frac{1}{\mathcal{A}^{(1)}} m_b^{(1)} + \frac{1}{\mathcal{A}^{(2)}} m_b^{(2)} + \frac{1}{\mathcal{A}^{(3)}} m_b^{(3)},$$

where $\mathcal{A}^{(i)}$ is the area of the ith T^2. When setting $\mathcal{A}^{(i)}$ for ξ_a and ξ_b to vanish, we can find a supersymmetric vacuum with unbroken $U(N_a)$ and $U(N_b)$ gauge symmetries.[3] This means that some of the Kähler moduli fields are stabilized by the D-term potential at the supersymmetric vacuum. In the present case, only the ratios of $\mathcal{A}^{(i)}$ are completely determined unless $m_a^{(1)} = m_a^{(2)} = m_a^{(3)} = 0$ and/or $m_b^{(1)} = m_b^{(2)} = m_b^{(3)} = 0$, and thus, only a linear combination of the

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[3]Magnetized supersymmetric vacua with broken $U(N_a)$ and $U(N_b)$ can also exist when charged fields develop their non-vanishing VEV in D-flat directions. This was discussed in Ref. [35].

three Kähler moduli remains massless. There exists one flat direction even when we consider more complicated configurations of the magnetic fluxes to get three or more unbroken gauge subgroups. The aim of this paper is to stabilize this remaining massless moduli field by a nonperturbative superpotential originating from E-branes.

B. Pati-Salam models based on D9-branes

We construct Pati-Salam models based on a stack of eight D9-branes, whose low-energy effective field theory is 10D $U(8)$ SYM theory. In the rest of this paper, we consider $Z_2 \times Z_2'$ orbifolding to eliminate harmful zero modes, which acts as

$$Z_2: (z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3),$$
$$Z_2': (z_1, z_2, z_3) \rightarrow (z_1, -z_2, -z_3). \quad (2.6)$$

Under these Z_2 and Z_2' symmetries, the superfields transform properly [see Eq. (2.5)] with projection operators P and P', respectively. For later convenience we define the following matrix,

$$P_{\alpha\beta\gamma} = \begin{pmatrix} \alpha \times \mathbf{1}_4 & 0 & 0 \\ 0 & \beta \times \mathbf{1}_2 & 0 \\ 0 & 0 & \gamma \times \mathbf{1}_2 \end{pmatrix}, \quad (2.7)$$

where α, β, and γ take $+1$ or -1 and $\mathbf{1}_n$ denotes an $(n \times n)$ unit matrix. Orbifolding with a projection operator of this form must respect the Pati-Salam gauge group.

In the $U(8)$ SYM theories, magnetic fluxes are represented by three 8×8 matrices. It is convenient to parametrize them as

$$M^{(1)} = \mathrm{diag}(0,0,0,0,X,X,-Y,-Y) + a \times \mathbf{1}_8,$$
$$M^{(2)} = \mathrm{diag}(0,0,0,0,-1,-1,0,0) + b \times \mathbf{1}_8,$$
$$M^{(3)} = \mathrm{diag}(0,0,0,0,0,0,1,1) + c \times \mathbf{1}_8,$$

where $a, b, c \in \mathbb{Z}$ and $X, Y \in \mathbb{N}$. Note that the 4D effective theories are independent of a, b, and c within the D9-brane sector except for the FI parameters. They will play a significant role in the association with E-branes. These magnetic fluxes break the $U(8)$ gauge group down to the Pati-Salam gauge group, $U(4)_C \times U(2)_L \times U(2)_R$ up to $U(1)$ factors, and produce the FI terms for diagonal parts of them as

$$\xi_C = \frac{1}{\mathcal{A}^{(1)}} a + \frac{1}{\mathcal{A}^{(2)}} b + \frac{1}{\mathcal{A}^{(3)}} c,$$
$$\xi_L = \frac{1}{\mathcal{A}^{(1)}} (a + X) + \frac{1}{\mathcal{A}^{(2)}} (b - 1) + \frac{1}{\mathcal{A}^{(3)}} c,$$
$$\xi_R = \frac{1}{\mathcal{A}^{(1)}} (a - Y) + \frac{1}{\mathcal{A}^{(2)}} b + \frac{1}{\mathcal{A}^{(3)}} (c + 1).$$

These FI parameters vanish when

$$\mathcal{A}^{(1)}/\mathcal{A}^{(2)} = X, \qquad \mathcal{A}^{(1)}/\mathcal{A}^{(3)} = Y, \qquad a + Xb + Yc = 0. \quad (2.8)$$

At supersymmetric vacua with the Pati-Salam gauge group, this implies that two of the three Kähler moduli are stabilized by the D term.

On this magnetized orbifold with $P' = P_{+--}$ [see Eq. (2.7)], there remain the following zero modes,

$$\Phi_1 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & H \\ 0 & 0 & 0 \end{pmatrix}, \qquad \Phi_2 = \begin{pmatrix} 0 & Q_L & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix},$$

$$\Phi_1 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ Q_R & 0 & 0 \end{pmatrix},$$

where three rows and columns correspond to $U(4)_C$, $U(2)_L$, and $U(2)_R$. We can find degenerate zero modes in the bifundamental representation $(1, \mathbf{2}, \bar{\mathbf{2}})$; $(\mathbf{4}, \bar{\mathbf{2}}, 1)$; and $(\bar{\mathbf{4}}, 1, \mathbf{2})$, which can be identified with the Higgs fields H, the left-handed matter fields Q_L, and the right-handed matter fields Q_R, respectively. Their degeneracy, that is, the number of generations, is determined by X, Y, and the Z_2 projection operator P. Three-generation magnetized orbifold models based on the Pati-Salam gauge group were systematically studied in Ref. [11]. According to that, we summarize all possible *Ansätze* of (X, Y, P) for realizing the three generations of the quarks and the leptons in Table II. In these models, a reasonable mechanism to realize hierarchical masses and mixing angles works, which leads to a semirealistic spectrum without fine-tunings for the parameters [36]. Note that there are other configurations to realize the three generations, but they have a

TABLE II. They are all possible sets of X, Y, and P to realize the three-generation structure of the SM. One can exchange the values of X and Y in configurations 5 and 6. In configurations 7–10, we have to replace the projection operator by P_{+++} when exchanging X and Y.

	X	Y	P	Number of Higgs
Number 1	4	4	P_{+-+}	5
Number 2	5	5	P_{+-+}	6
Number 3	7	7	P_{++-}	8
Number 4	8	8	P_{++-}	9
Number 5	4	5	P_{+-+}	5
Number 6	7	8	P_{++-}	8
Number 7	4	7	P_{+--}	5
Number 8	4	8	P_{+--}	5
Number 9	5	7	P_{+--}	5
Number 10	5	8	P_{+--}	6

phenomenological difficulty in textures of Yukawa matrices and we have omitted them here. It is remarkable that zero modes cannot remain in diagonal entries of the above matrices which correspond to open string moduli fields. That is, open string moduli are completely stabilized. The idea of this open string moduli stabilization would be a T-dual picture to intersecting D-branes wrapping rigid cycles [24].

One may expect that magnetized backgrounds with more complicated gauge symmetry breaking, e.g., $U(8) \rightarrow U(3)_C \times U(1)_\ell \times U(2)_L \times U(2)_R$, lead to a new class of three-generation models. In that case, however, a non-vanishing FI term inevitably appears within the 10D $U(8)$ SYM theories [36]. We will propose such a model with all the vanishing FI parameters on the basis of D7-brane systems in Sec. III.

C. Nonperturbative superpotential: E1-branes

We study E-branes in D9-brane models. In general, there can be various E-branes generating superpotential. Here, we focus only on E-brane configurations that contribute to the moduli stabilization. In the presence of D9-branes, two types of E-branes can possibly lead to a stable brane system: E1-branes wrapping 2-cycles and E5-branes. E-branes generically have $O(N)$ or $USp(N)$ gauge groups, and only the $O(N)$-type instantons can generate the superpotential. In the present setup, we can choose discrete torsions to obtain the $O(N)$-type instantons [24], and we assume that the discrete torsions are tuned suitably in this paper.

These instantons can induce a superpotential of the form

$$W_{np} = \sum_i A_i e^{-a_i T_i} + A_S e^{-S}, \qquad (2.9)$$

where T_i and S are Kähler moduli and dilaton superfields, respectively. Coefficients A_i and A_S depend on complex moduli fields, which are supposed to be stabilized by the 3-form fluxes and replaced by their VEVs. In the present case, they are given by

$$T_i = e^{-\phi} \mathcal{A}^{(i)} + i \int_{T^2} C_2,$$

$$S = e^{-\phi} \mathcal{A}^{(1)} \mathcal{A}^{(2)} \mathcal{A}^{(3)} + i \int_{T^6} C_6,$$

where C_2 and C_6 are RR-forms and ϕ is the 10D dilaton field. The SUSY condition (2.8) stabilizes two directions of T_i. It is important for this superpotential to change or vanish if there exist open string zero modes between the D9-branes and the E-branes. We have to study configurations of these branes in order to eliminate such harmful zero modes.

First we study E1-branes, which wrap one of the three T^2 and are collapsed at a fixed point on the other T^2. A single

E1-brane has an $O(1)$ gauge symmetry and is to generate the superpotential for the Kähler moduli [the first term of Eq. (2.9)] as long as there is no extra zero mode. A zero mode configuration of E1/D9 systems is equivalent to that of a system consisting of D9-branes and an unfluxed D5-brane wrapping the ith T^2. Such a D-brane system contains a six-dimensional $\mathcal{N} = 1$ hypermultiplet as D5-D9 (or E1-D9) open strings. Naming these D9-branes "D9$_A$," we can represent the hypermultiplet by using two 4D $\mathcal{N} = 1$ superfields as $(\Phi_j^{AE}, \Phi_k^{EA})$ ($i \neq j \neq k \neq i$) in the superfield description (see Ref. [37]). Note that superscripts AE and EA reflect their gauge transformation laws, and they are an (anti)fundamental representation of the $U(N)$ gauge group of the D9-branes. They are affected by the magnetic fluxes of the D9-branes, and thus a chiral spectrum with generation structure is produced in this E1-D9 sector, in the same way as the D9-brane sector. The transformation law of these chiral superfields under the Z_2 and Z_2' orbifolding is given in a similar way to the D9-brane fields, e.g.,

$$\Phi_1^{AE} \rightarrow -P\Phi_1^{AE} P_E^{-1}, \qquad \Phi_1^{AE} \rightarrow +P'\Phi_1^{AE} P_E'^{-1},$$

where we can set P_E and P_E' to ± 1. Note that all of the E1-branes with $P_E = \pm 1$ and $P_E' = \pm 1$ can appear and we have to take into account all the possible E1-branes including projections, $P_E = \pm 1$ and $P_E' = \pm 1$. However, some of them do not induce nonperturbative terms and others induce nonperturbative terms such as (2.9), as well as nonperturbative terms with matter fields. We are interested in E1-branes with proper orbifold parities, P_E and P_E', which can induce (2.9). When the superfield has a different subscript, the overall signs can be changed. Their wave functions can be even or odd functions on the ith T^2. On the other T^2, however, they cannot survive the orbifold projection when they are assigned into an odd mode, because they are localized at a fixed point of the T^2 and their wave functions must be given by a delta function.

We study how to find the E1-brane configurations where all the harmful massless modes are eliminated, taking an E1-brane wrapping the third T^2 as an example. For this purpose, it is satisfactory to investigate a zero-mode configuration of Φ_1^{AE} and Φ_2^{EA}. They transform under the Z_2 symmetry as

$$\Phi_1^{AE} \rightarrow -P\Phi_1^{AE} P_E^{-1}, \qquad \Phi_2^{EA} \rightarrow -P_E \Phi_2^{EA} P^{-1}.$$

They cannot have zero modes when they are assigned into a Z_2 odd mode on the first and the second T^2 as discussed above. Thus, for $P = P_{+++}$, we can eliminate all the components of Φ_1^{AE} and Φ_2^{EA} by $P_E = +1$. Even when $P \neq P_{+++}$, it is possible to eliminate them as follows. In the Pati-Salam models, both of them have eight components, which are classified into three parts by their gauge representations, i.e., $U(4)_C$, $U(2)_L$, and $U(2)_R$. A proper

choice for P_E can forbid the charged zero modes in two of the three parts. Seen from Table I, we can eliminate the remaining ones when the absolute values of their effective magnetic fluxes are less than 3 and they are assigned into the Z_2' odd mode on the third T^2. We can always find P_E' and c that realize such a situation, satisfying Eq. (2.8). Thus, it is always possible for the E1-brane to generate the nonperturbative superpotential. One can easily confirm that E1-branes wrapping the first or the second T^2 can also induce nonperturbative terms to stabilize the moduli.

We examine the stabilization of the moduli field minimizing its potential. We expect to obtain the following nonperturbative superpotential,

$$W = Ae^{-2\pi T_3} + W_0. \qquad (2.10)$$

We assumed that the nonperturbative term due to the E1-brane wrapping the third T^2 is dominant. Even when other terms are dominant, the following discussion is the same. A constant term W_0 is also necessary for the moduli stabilization, and we will discuss its origin later. In toroidal compactifications of type IIB with O5/O9 planes, the Kähler potential for the moduli fields is given by

$$K_0 = -\log(S + \bar{S}) - \sum_{i=1}^{3}\log(T_i + \bar{T}_i) - \sum_{i=1}^{3}\log(U_i + \bar{U}_i).$$

Setting $\mathrm{Re}T_i = \tau_i$ and $\mathrm{Im}T_i = 0$, we get the F-term potential

$$V_F = \frac{\pi Ae^{-4\pi\tau_3}}{\tau_1\tau_2}(A + 2\pi A\tau_3 + W_0e^{2\pi\tau_3}).$$

Minimizing this potential, we find a supersymmetric minimum,

$$\frac{W_0}{A} = -(1 + 4\pi\tau_3)e^{-2\pi\tau_3},$$

where τ_3 is stabilized. In this case, one sees that a legitimate value of $\langle\tau_3\rangle$ implies a quite small value of W_0; indeed, $\langle\tau_3\rangle = 1$ requires $W_0/A \sim 10^{-2}$. The origin of such a small W_0 will be discussed in the next subsection.

D. Nonperturbative superpotential: E5-branes

We perform a study similar to the previous subsection for E5-branes. The number of zero modes in D9-E5 open strings can be counted in the same way as a mixed configuration of the magnetized D9-branes and an additional D9-brane with no magnetic fluxes.

Although it is difficult in D9/E5 systems to give a setup to eliminate all the harmful zero modes systematically, we show a reasonable setup to generate the nonperturbative superpotential to be incorporated in a wide class of the Pati-Salam models shown in Table II. First we set $b = -1$ and

$c = +1$, which implies $a = X - Y$ for the vanishing D terms [see Eq. (2.8)]. That is, the magnetic fluxes in the Pati-Salam sector are given by

$$M^{(1)} = \mathrm{diag}(X - Y, X - Y, X - Y, X - Y, 2X - Y,$$
$$\qquad 2X - Y, X - 2Y, X - 2Y),$$
$$M^{(2)} = \mathrm{diag}(-1, -1, -1, -1, -2, -2, -1, -1),$$
$$M^{(3)} = \mathrm{diag}(1, 1, 1, 1, 1, 1, 2, 2).$$

When $2X - Y \neq 0$ and $X - 2Y \neq 0$, an association of chirality projections due to the magnetic fluxes and Z_2' orbifold projections with $P_E' = -1$ eliminates open string zero modes charged under $U(2)_L$ and $U(2)_R$. The remaining ones, which are (anti)fundamentals in the $U(4)_C$ gauge group, can also be eliminated by Z_2 orbifolding with a suitable choice for P_E, if $0 < |X - Y| < 3$. Thus we can always provide configurations of D9/E5 systems to generate the nonperturbative superpotential for the dilaton superfield [the second term of Eq. (2.9)], when X and Y satisfy

$$2X - Y \neq 0, \qquad X - 2Y \neq 0, \qquad 0 < |X - Y| < 3.$$

That is, models 5, 6, and 9 shown in Table II are available (we can exchange the values of X and Y as discussed there). In particular, we find that some of these models can be associated with an E5-brane and an E1-brane simultaneously (e.g., $X = 7$, $Y = 5$, and $P = P_{+++}$). In this case, we can obtain the nonperturbative superpotential

$$W = A_E e^{-2\pi T_3} + A_S e^{-S}. \qquad (2.11)$$

We have assumed the presence of supersymmetric 3-form fluxes to stabilize the dilaton satisfying $\langle W_{3-\mathrm{form}}\rangle = 0$. The superfield S can be replaced by its VEV, and then the effective superpotential is equivalent to Eq. (2.10), that is,

$$W_0 = A_S e^{-\langle S\rangle}.$$

From this expression, it is found that a reasonable value of $\langle S\rangle$ induces a sufficiently small W_0 which is required for the above successful moduli stabilization. Thus, all the moduli fields can be stabilized in the framework of magnetized D-branes by an interplay of the two instanton effects.

III. D7-BRANE MODELS

In this section, we consider another model based on D7-branes, instead of the Pati-Salam models based on D9-branes.[4]

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[4]The model discussed in this section was proposed in Ref. [38].

TABLE III.   The configuration of two stacks of D7-branes is shown. A symbol "✓" means that D-branes wrap T^2, and another one "×" expresses that D-branes are localized at a fixed point on T^2.

	T^2	T^2	T^2
D7$_A$	✓	×	✓
D7$_B$	✓	✓	×

A. MSSM-like model

We consider a MSSM-like model on the basis of two stacks of four D7-branes, which we denote by D7$_A$-branes and D7$_B$-branes with a configuration shown in Table III. An effective field theory of D7-branes is derived from a 10D SYM theory, and the superfield description of that was formulated in Ref. [37]. One of the three chiral superfields Φ_i contained in 10D SYM theories turns into a position moduli field there. In the present case of the mixed D7-brane system, there also appears a hypermultiplet corresponding to the open string modes between the D7$_A$- and D7$_B$-branes, which is denoted by two chiral superfields Φ_2^{AB} and Φ_3^{BA}. Thus, this system consists of the following chiral superfields,

$$\Phi_1^A, \quad \tilde{\Phi}_2^A, \quad \Phi_3^A, \quad \Phi_1^B, \quad \Phi_2^B, \quad \tilde{\Phi}_3^B, \quad \Phi_2^{AB}, \quad \Phi_3^{BA}.$$

The first three superfields are in the $U(4)_A$ adjoint representation, and the next three are in the $U(4)_B$ adjoint one. The last two are a bifundamental representation of $U(4)_A \times U(4)_B$. The tilde represents that the superfield turns out to be a position moduli of the corresponding D7-branes.

In this section, we again consider $T^2 \times T^2 \times T^2$ as the extra compact space with $Z_2 \times Z_2'$ orbifolding. These $Z_2 \times Z_2'$ act on the three T^2 in the same way as in the previous section, and the transformation laws of the superfields are determined by their subscript and four 4×4 projection matrices, P_A, P_B, P_A', and P_B'. Note that active D7-brane fields must be assigned into an even mode on T^2 where the D7-brane is localized as a point because such a pointlike localization implies a wave function of the delta function. In particular, D7-D7 open strings, Φ_2^{AB} and Φ_3^{BA}, have to be assigned into an even mode on the second and the third T^2 in order to survive the orbifold projections.

We introduce the magnetic fluxes in this D7$_A$/D7$_B$ brane system as follows:

$$M_A^{(1)} = \begin{pmatrix} -5 \times \mathbf{1}_3 & 0 \\ 0 & -4 \times \mathbf{1}_1 \end{pmatrix}, \quad M_A^{(3)} = \begin{pmatrix} 5 \times \mathbf{1}_3 & 0 \\ 0 & 4 \times \mathbf{1}_1 \end{pmatrix},$$

$$M_B^{(1)} = \begin{pmatrix} 0 \times \mathbf{1}_3 & 0 \\ 0 & -12 \times \mathbf{1}_2 \end{pmatrix}, \quad M_B^{(2)} = \begin{pmatrix} 0 \times \mathbf{1}_2 & 0 \\ 0 & 1 \times \mathbf{1}_2 \end{pmatrix}.$$

These magnetic fluxes break $U(4)_A \times U(4)_B \to U(3)_C \times U(1)_l \times U(2)_L \times U(2)_R$. One remarkable feature of this model is the breaking of the $U(4)_C$ gauge symmetry of the Pati-Salam models. This means that the quarks and the leptons can have a distinguished difference in their flavor structure. The flux-induced FI terms vanish in all the unbroken gauge subgroups when

$$\mathcal{A}^{(1)}/\mathcal{A}^{(2)} = 12 \quad \text{and} \quad \mathcal{A}^{(1)}/\mathcal{A}^{(3)} = 1. \qquad (3.1)$$

Setting the projection operators as $P_A = P_B = P_A' = \mathbf{1}_4$ and $P_B' = -\mathbf{1}_4$, we find the following zero-mode structure,

$$\Phi_1^B = \begin{pmatrix} 0 & H \\ 0 & 0 \end{pmatrix}, \quad \Phi_2^{AB} = \begin{pmatrix} Q_L & 0 \\ L_L & 0 \end{pmatrix}, \quad \Phi_3^{BA} = \begin{pmatrix} 0 & 0 \\ Q_R & L_R \end{pmatrix},$$

and Φ_1^A, Φ_2^A, Φ_3^A, Φ_2^B, and Φ_3^B have no zero mode. We can identify H, Q_L, Q_R, L_L, and L_R with the Higgs fields, the left-handed quarks, the right-handed quarks, the left-handed leptons, and the right-handed leptons of the MSSM, respectively. All of the position and Wilson-line moduli fields are stabilized in this model as well as in the D9-models.

B. Nonperturbative superpotential: E3-branes

In the present D7-brane system, there are two types of E-branes keeping the whole brane system stable: E3-branes and E(−1)-branes. When there are no open string zero modes interplaying the D-branes and the E-branes, these instantons generate the nonperturbative superpotential [note again that we have assumed discrete torsions tuned to obtain $O(N)$-type E-branes],

$$W_{\mathrm{np}} = \sum_i A_i e^{-a_i T_i} + A_S e^{-S}. \qquad (3.2)$$

In the IIB orientifold with $O3/O7$-planes, T_i and S are given by ($i \neq j \neq k \neq i$)

$$T_i = e^{-\phi} \mathcal{A}^{(j)} \mathcal{A}^{(k)} + i \int_{T^4} C_4, \qquad S = e^{-\phi} + iC_0,$$

and again, we see that two of the Kähler moduli fields are stabilized by Eq. (3.1).

We first discuss an E3-brane wrapping two of three T^2 and localized at a fixed point on the other one, which has an $O(1)$ gauge symmetry and generates the first term superpotential of Eq. (3.2), without extra zero modes. Generic E3/D7 systems are classified into two cases. One is the case when the E-branes and the D-branes wrap the same $T^4 = T^2 \times T^2$ and are localized at fixed points on the last T^2. In this case, it is easy to eliminate the E3-D7 open string zero modes, because the two stacks of the branes can be sequestered spatially when the two stacks are localized at different fixed points. In the other case, we have to study

the zero-mode distribution in detail for each model. Recall again that any E3-brane including all the possible positions and orbifold parities can appear and we have to take into account all the possibilities. However, we are interested only in E3-brane configurations to lead moduli-dependent superpotential terms.

In the present D7-brane system, there are two stacks of D7-branes which wrap the four different directions of extra compact space. An additional E-brane can be sequestered from one stack by a localization at different fixed points, but there exist massless open strings between the E-brane and the other stack of D7-branes to be eliminated by the orbifold projection. Let us consider an E3-brane which wraps the first and the second T^2 and is localized at a "vacant" fixed point on the third T^2. E3-D7$_B$ zero modes cannot appear, but there are E3-D7$_A$ open strings denoted by Φ_2^{AE} and Φ_3^{EA}. Fortunately, we can eliminate them easily as follows. They transform under the Z_2' symmetry $(P_A' = +\mathbf{1}_4)$ as

$$\Phi_2^{AE} \to -\Phi_2^{AE} P_E', \qquad \Phi_3^{EA} \to -P_E'^{-1}\Phi_3.$$

Thus they all can be assigned into a Z_2' odd mode by $P_E' = +1$ and are eliminated as we wanted, because wave functions of these open strings must be an even function on the second and the third T^2 as explained above. As a result we obtain the superpotential

$$W = A_3 e^{-2\pi T_3}.$$

We will see that this stabilizes the moduli in association with an additional E(-1)-brane in the following subsection.

C. Nonperturbative superpotential: E($-$1)-branes

It is much easier to find an E(-1)-brane configuration generating the superpotential for the dilaton superfield. The E(-1)-brane is an instanton localized completely at a point on the whole compact space. Thus we can trivially sequester the E(-1)-brane from the D7-brane system in order not to produce the harmful zero modes, unless four fixed points of a $T^2/Z_2(Z_2')$ are occupied by multiple stacks of D7-branes.

One can straightforwardly see that the D7-brane system admits an E(-1)-brane and an E3-brane simultaneously and superpotential (2.11) is generated. Similarly to the D9-brane systems, the second term of Eq. (2.11) produces a small constant term in the superpotential, and the Kähler moduli field is stabilized with a moderate value of the VEV.

IV. CONCLUSION AND DISCUSSION

We have studied the nonperturbative superpotential induced by E-branes in semirealistic D-brane models based on the toroidal orbifolds.

We have considered two types of D-brane models for the visible sector. One is based on a stack of eight D9-branes, where the magnetic fluxes and the orbifold projection yield the Pati-Salam gauge group with the three generations of the quarks and the leptons. In the models, magnetic fluxes generate FI terms, which depend on the Kähler moduli, and those fix the ratio among three Kähler moduli. Furthermore, we have found that an E1-brane and an E5-brane generate the superpotential for the dilaton and the Kähler moduli, respectively. The dilaton is replaced by its VEV in the nonperturbative superpotential because we have assumed the 3-form fluxes to stabilize that. That gives rise to the sufficient small constant term, and as a result, the Kähler moduli field is stabilized with a moderately large value of the VEV. The other D-brane model is derived from the two stacks of the D7-branes. In this model, the moduli-dependent FI terms can fix the ratio of three Kähler moduli. On top of that, we have found that an E3-brane and an E(-1)-brane successfully generate the superpotential and stabilize the moduli. In our study, we have found some constraints on the magnetic fluxes and the orbifold parities for realizing the moduli stabilization, and it is quite nontrivial that there exists a successful configuration of D-branes and E-branes.

In this paper, we have studied moduli stabilizations with only the visible sector. Since the vacuum we have studied is supersymmetric and has the negative energy, we need SUSY breaking and uplifting of the vacuum energy to almost zero energy. Thus, towards more realistic models, we should also consider a hidden sector for SUSY breaking.[5] A further additional hidden sector might be requested for the RR tadpole cancellations and/or other reasons. In these cases, we have to care about open string zero modes between the E-branes and the hidden D-branes, because the moduli stabilizing superpotential vanishes if there appears an extra zero mode. Besides the open string zero modes, we expect that there are several important interplays between the SUSY breaking and the moduli stabilization. It seems that such an extension to accommodate the SUSY breaking sector is a very challenging task towards realistic D-brane models.

We have also assumed that the orientifold projections are well defined and we can obtain the desired gauge groups and zero modes. Although it might be challenging to embed the models shown in this paper into a full system of the string theory holding those assumptions, that would be an important task towards getting realistic string models and we would study that elsewhere.

Nonperturbative effects due to E-branes are applied to other phenomenological issues than the moduli stabilization. Another challenging task of D-brane models is to obtain Majorana mass terms and the supersymmetric Higgs mass term (μ term). We are able to consider additional

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[5]See, e.g., for explicit construction of the SUSY breaking sector [39].

E-branes to generate these mass terms [19–23,28,29]. It is also an attractive prospect to try that in the D-brane models shown in this paper.

APPENDIX: T-DUAL PICTURE

Magnetized D-brane systems are T-dual to intersecting D-branes. Although they are physically equivalent to each other, one may more easily be able to investigate the remaining zero modes in intersecting D-brane systems than in magnetized D-brane systems. In this appendix, we introduce an instrument to count the active zero modes in the T-dual picture, i.e., intersecting D6-branes wrapping rigid 3-cycles on $T^6/Z_2 \times Z_2'$ with discrete torsion (see Ref. [24] for reference).

1. T-dual to D9-brane models

E1-branes and E5-branes discussed in Sec. II are both equivalent to E2-branes wrapping different extra-dimensional directions in the T-dual side. That is schematically depicted in Fig. 1. Note that this figure does not take into account orbifolding for simplicity.

The number of zero modes between an E2-brane and a D6-brane is counted by the topological intersection number [24],

$$\frac{1}{4}\prod_{i=1}^{3}(n_E^{(i)}\tilde{m}_D^{(i)} - \tilde{m}_E^{(i)}n_D^{(i)})$$
$$+ \frac{1}{4}\sum_{g\in G}\sum_{i,j\in J_g^E}\sum_{k,l\in J_g^D}\epsilon_{E,ij}^g\epsilon_{D,kl}^g\delta_{ik}\delta_{jl}(n_E^{(I_g)}\tilde{m}_D^{(I_g)} - \tilde{m}_E^{(I_g)}n_D^{(I_g)}).$$
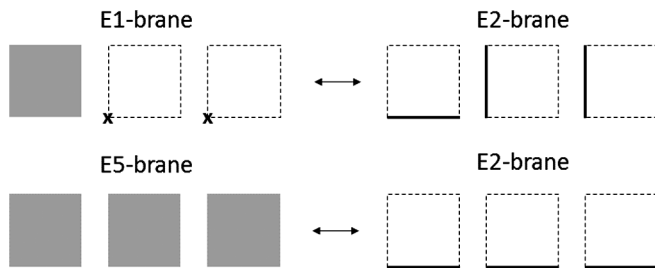$$(A1)$$

FIG. 1. A set of three squares expresses a fundamental region of $T^2 \times T^2 \times T^2$. The left and right sides correspond to each other in T-duality along three vertical axes.

In this expression, subscripts E and D express the E-brane and the D-brane. When we denote nontrivial elements of Z_2 and Z_2' by θ and θ', respectively, G is a set of θ, θ', and $\theta\theta'$. For each g, sets of fixed points where the E-branes and the D-branes live are given by J_E^g and J_D^g, respectively. There are two possible orientations on 2-cycles collapsed at a fixed point contained in J_E^g or J_D^g. This degree of freedom is defined by $\epsilon_{a,ij}^g = \pm 1$. In the magnetized D9-brane models, that corresponds to the discrete Wilson lines and parities P and P'. (We have not considered the Wilson lines in this paper, and then we get $\epsilon_{E,ij}^g = \epsilon_E^g$ and $\epsilon_{D,kl}^g = \epsilon_D^g$.) A set of $(n_D^{(i)}, m_D^{(i)})$ represents winding numbers along two fundamental cycles of the ith T^2, and $(n_a^{(I_g)}, m_a^{(I_g)})$ denotes winding numbers on a T^2 invariant under $g \in Z_2 \times Z_2'$. That is, in the present case (2.6), we see $(I_\theta, I_{\theta'}, I_{\theta\theta'}) = (3, 1, 2)$. The tilde on the winding number is a reflection of nontrivial complex structure, e.g., $\tilde{m}_D^{(i)} \equiv m_D^{(i)} + \frac{1}{2}n_D^{(i)}$ when the torus is tilted, and $\tilde{m}_D^{(i)} = m_D^{(i)}$ when the torus is rectangular. In the following, we take $\tilde{m}^{(i)} = m^{(i)}$ for simplicity which is satisfactory for the aim of this section.

In the upper case of Fig. 1, the winding numbers of the corresponding E2-brane are given by

$$(n_E^{(1)}, m_E^{(1)}) = (-1, 0), \qquad (n_E^{(2)}, m_E^{(2)}) = (0, 1),$$
$$(n_E^{(3)}, m_E^{(3)}) = (0, 1). \qquad\qquad (A2)$$

In the lower case, the winding numbers are

$$(n_E^{(i)}, m_E^{(i)}) = (1, 0) \quad \forall\, i. \qquad (A3)$$

For example, we will count the number of E2-D6 open string zero modes with winding numbers (A2). The corresponding intersection number is given by

$$I_{AE} = \frac{1}{4}m_D^{(1)} + \frac{1}{4}\epsilon_D^{\theta'}\epsilon_E^{\theta'}m_D^{(1)}$$
$$+ \frac{1}{4}\sum_{i,j\in S_{\theta'}^a}\sum_{k,l\in S_{\theta'}^b}\epsilon_D^\theta\epsilon_E^\theta\delta_{ik}\delta_{jl}(1 + \epsilon_D^{\theta'}\epsilon_E^{\theta'}).$$

One can see that this intersection number vanishes when $\epsilon_D^{\theta'}\epsilon_E^{\theta'} = -1$. Thus we can add an E2-brane into the Pati-Salam models based on the eight D6-branes to generate the moduli stabilizing superpotential, if $\epsilon_{D_a}^{\theta'}\epsilon_E^{\theta'} = -1$ is held for all of the eight D6-branes ($a = 1, 2, …, 8$). This implies $\epsilon_{D_1}^{\theta'} = \epsilon_{D_1}^{\theta'} = \cdots = \epsilon_{D_8}^{\theta'}$. The same result has been obtained in the magnetized D-brane systems; that is, we have shown that the zero modes of E1/D9 open strings are completely eliminated for $P = P_{+++}$ in Sec. II C. Similarly one can regain the result obtained in Sec. II D by using the winding numbers (A3).

2. T-dual to D7-brane models

We study the T-dual picture of the D7-brane model with E3- and E(-1)-branes. They correspond to two types of E2-branes in the T-duality as shown in Fig. 2. In the upper case, the winding number of the E2-brane is given by

$$(n_E^{(1)}, m_E^{(1)}) = (1, 0), \qquad (n_E^{(2)}, m_E^{(2)}) = (1, 0),$$
$$(n_E^{(3)}, m_E^{(3)}) = (0, 1).$$

In the other case one sees

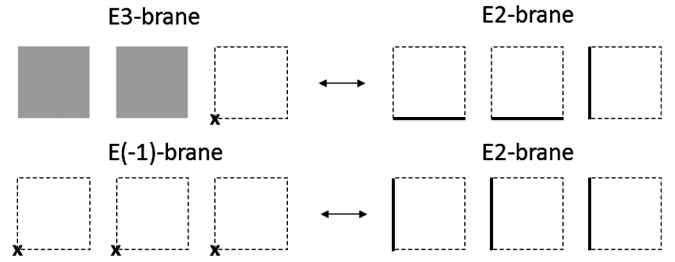$$(n_E^{(i)}, m_E^{(i)}) = (0, 1). \quad \forall\, i$$

FIG. 2.   A set of three squares expresses a fundamental region of $T^2 \times T^2 \times T^2$. The left and right sides correspond to each other in T-duality along three vertical axes.

Substituting them in Eq. (A1), one is able to confirm the result obtained in the previous section.

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