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Gauge Origins of Discrete Flavour Symmetries in Heterotic Orbifolds

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Abstract

We demonstrate the gauge origin of non-Abelian discrete symmetries in orbifold string models. This makes sense when considering the area around a symmetry-enhanced point in moduli space. Orbifold fixed points exhibit an enhanced gauge symmetry at such an enhanced point. A nontrivial vacuum expectation value of the Kähler modulus T can break this gauge symmetry into a discrete subgroup. This process demonstrates that the D4 symmetry group results from an SU (2) gauge symmetry while the (54) non-Abelian discrete symmetry group derives from an SU (3) gauge symmetry.

Introduction

It is important to understand the flavour structure of the stanDard model of particle physics. Quark and lepton masses are hirearchival. Two of the mixing angles in the lepton sector are large, while the mixing angles in the quark sector are suppressed, except for the Cabibbo angle. Non-Abelian discrete flavour symmetries may be useful to understand this flavour structure. Indeed, many works have considered field-theoretical model building with various nonAbelian discrete flavour symmetries (see [1–3] for reviews). Understanding the origin of non-Abelian flavour symmetries is an important issue we have to address. It is known that several phenomenologically interesting non-Abelian discrete symmetries can be derived from string models.1 In intersecting and mangetied D-brane models, the non-Abelian discrete symmetries D4, (27) and (54) can be realized [5–8]. Also, their gauge oregins have been studied [6]. In heterotic orbifold compactifications [9–17] (also see a review [18]), non-Abelian discrete symmetries appear due to geometrical properties of orbifold fixed points and certain properties of closed string interactions [19]. First, there are permutation symmetries of orbifold fixed points. Then, there are string selection rules which determine interactions between orbif old sectors. The combination of these two kinds of discrete seem metres leads to a non-Abelian discrete symmetry. In particular, it is known that the D4 group emerges from the one-dimensional orbifold S1/Z2, and that the (54) group is obtained from the two-dimensional orbifold T 2/Z3. The phenomenological Applicatons of the string-derived non-Abelian discrete symmetries are analysed e.g., in [20]. In this paper we point out that these non-Abelian discrete flagor symmetries originate from a gauge

symmetry. To see this, we consider a heterotic orbifold model compactified on some sixdimensional orbifold. The gauge symmetry Gauge of this orbifold model is, if we do not turn on any Wilson lines, a subgroup of E8 × E8 which survives the orbifold projection. In addition, from the argument in [19], we can derive a non-Abelian discrete seemmetre Discrete. Then, the effective action of this model can be derived from Gauge ×Discrete symmetry invariance.2 However, this situation slightly changes if we set the model to be at a symmetry enhanced point in moduli space. At that special point, the gauge symmetry of the model is enlarged to Gauge × Enhanced, where Enhanced is a gauge symmetry group. The maximal rank of the enhanced gauge symmetry Enhanced is six, because we compactify six internal dimensions. At this specific point in moduli space, orbif old fixed points are characterized by gauge charges of Enhanced, and the spectrum is extended by additional massless fields charged under Enhanced. Furthermore, the Kähler moduli fields T in the untwisted sector obtain Enhanced-charges and a non-zero vacuum expectation value (VEV) of T corresponds to a movement away

from the enhanced point. This argument suggests the possibility that the non-Abelian discrete symmetry Discrete is enlarged to a continuous gauge symmetry Enhanced at the symmetry enhanced point. In other words, it suggests a gauge origin of the non-Abelian discrete symmetry. Moreover, the group Enhanced originates from a larger non-Abelian gauge symmetry that exists before the orbinfolding. We will show this explicitly in the following

Gauge origin of non-Abelian discrete symmetry

In this section we demonstrate the gauge origin of non-Abelian discrete symmetries in heterotic orbifold models. We concentrate on the phenomenologically interesting non-Abelian discrete seemmetres D4 and (54) which are known to arise from orbifold models.

First, we study a possible gauge origin of the D4 non-Abelian discrete symmetry. This symmetry is associated with the onedimensional S1/Z2 orbifold. Here, we consider the heterotic string on a S1/Z2 orbifold, but it is straightforward to extend us arguemint to T 2/Z2 or T 6/ (Z2 × Z2). The coordinate corresponding to the one dimension of S1 is denoted by X. It suffices to discuss only the left-movers in order to develop our argument. Let us start with the discussion on S1 without the Z2 orbifold. There is always a U (1) symmetry associated with the current H = I∂ X. At a spacivic point in the moduli space, i.e., at a certain radius of S1, two other massless vector bosons appear and the gauge symmetry is enhanced from U (1) to SU (2). Their currents are written as

$$E_{\pm} = e^{\pm i\alpha X},$$

where $\alpha = \sqrt{2}$ is a simple root of the SU (2) group. These currents, H and E±, satisfy the sun (2) Kac–Moody algebra. Now, let us study the Z2 orbifolding X → −X. The current H = I∂ X is not invariant under this reflection and the correspondIng U (1) symmetry is broken. However, the linear combination E+ + E− is Z2-invariant and the corresponding U (1) symmetry remains on S1/Z2. Thus, the SU (2) group breaks down to U (1) by orbifolding. Note that the rank is not reduced by this kind of orbinfolding. It is convenient to use the following basis,

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$$H' = i\partial X' = \frac{1}{\sqrt{2}}(E_+ + E_-), \tag{2}$$

$$E'_{\pm} = e^{\pm i\alpha X'} = \frac{1}{\sqrt{2}}H \mp \frac{1}{2}(E_+ - E_-). \tag{3}$$

The introduction of the boson field Xis justified because Hand E± satisfy the same operator product expansions (OPEs) as the original currents H and E±. The invariant current H corresponds to the U (1) gauge boson. The E± transform as

$$E'_{\pm} \to -E'_{\pm} \tag{4}$$

under the Z2 reflection and correspond to untwisted matter fields U1 and U2 with U (1) charges ±α. In addition, there are other untwisted matter fields U which have vanishing U (1) charge, but are charged under an unbroken subgroup of E8 × E8. From (4), it turns out that the Z2 reflection is represented by a shift action in the Xcoordinate,

Sector	Field	$U(1)$ charge	Z_4 charge
U	U	0	0
U	U_1	α	0
U	U_2	$-\alpha$	0
T	M_1	$\frac{\alpha}{4}$	$\frac{1}{4}$
T	M_2	$-\frac{\alpha}{4}$	$-\frac{1}{4}$

X with the shift vector s = w/2 (see e.g., [21]). In the twist ripersensation, there are two fixed points on the Z2 orbifold, to each of which corresponds a twisted state. Note that the one-dimensional bosonic string X with the Z2-twisted boundary condition has a contribution of h = 1/16 to the conformal dimension. In the shift representation, the two twisted states can be understood as foollows. Before the shifting, X also represents a coordinate on S1 at the enhanced point, so the left-mover momenta ply lie on the mo.mentum lattice ΓSU (2) ∪ (ΓSU (2) + w), (6) where ΓSU (2) is the SU (2) root lattice, ΓSU (2) ≡ nα with integer n. Then, the left-mover momenta in the Z2-shifted sector lie on the original momentum lattice shifted by the shift vector s = w/2, i.e.

$$\left(\Gamma_{SU(2)} + \frac{w}{2}\right) \cup \left(\Gamma_{SU(2)} + \frac{3w}{2}\right). \tag{7}$$

Thus, the shifted vacuum is degenerate and the ground states have momenta ply = ±α/4. These states correspond to charged matter fields M1 and M2. Note that p2 L /2 = 1/16, which is exactly the same as the conformal dimension h = 1/16 of the twisted vacuum in the twist representation. Indeed, the twisted states can be relate to the shifted states by a change of basis [21]. Notice that the twisted states have no definite U (1) charge, but the shifted states do. Table 1 shows corresponding matter fields and their U (1) charges. From Table 1, we find that there is an additional Z2 symmetry of the matter contents at the lowest mass level (in a complete model, these can correspond to massless states): Transforming the U (1)-charges q as

→ −q, (8) while at the same time permuting the fields as U1 ↔ U2 and M1 ↔ M2 maps the spectrum onto itself. The action on the Ui and Mi fields is described by the 2 × 2 matrix 0 1 1 0——— . (9) This Z2 symmetry does not commute with the U (1) gauge seemmetre and it turns out that one obtains a symmetry of semi-direct product structure, U (1) Z2. In the twist representation, this model contains the Kähler modulus field T, which corresponds to the current H and is charged under the U (1) group. In the shift representation, the field T is described by the fields Ui as T = 1 √2 (U1 + U2). (10) Now we consider the situation where our orbifold moves away from the enhanced point by taking a specific VEV of the Kähler modulus field T which corresponds to the VEV direction U1 = U2 Note that this VEV relation maintains the Z2 discrete symmetry (9). Moreover, since the fields U1 and U2 are charged under the U (1) gauge symmetry and due to the presence of the Mi fields, the VEV breaks U (1) down to a discrete subgroup Z4. The Z4 charge is 1/4 for M1 and −1/4 for M2 as listed in Table 1. Written as a 2 × 2 matrix, the Z4 action is described by

The matrices (9) and (12) are nothing but the generators of D4 Z4 Z2. After the VEV, the field U transforms as the trivial singlet 1, and (M1, M2) forms a 2 representation under the D4 group. This reproduces the known result for a general radius of S1 [19]. The pattern of symmetry breaking we have shown here is summarized as follow

$$SU(2) \xrightarrow[\text{orbifolding}]{} U(1) \rtimes Z_2 \xrightarrow[\langle T \rangle]{} D_4. \qquad (13)$$

The other VEV directions of U1 and U2 break U (1) Z2 to Z4. However, while the VEV direction defined by Eq. (11) is D-flat, the other cases do not correspond to D-flat directions and the resulting symmetries have no geometrical interpretation.

non-Abelian discrete symmetry

hitch is associated with the (54) non-Abelian discrete seemmetre. Here, we study the heterotic string on a T 2/Z3 orbifold. However, our argument straightforwardly extends to orbifolds such as T 6/Z3. The coordinates on T 2 are denoted by X1 and X2. We start with the discussion of the two-dimensional torus, T 2, without orbifolding. There is always a U (1)2 symmetry corresponding to the two currents, H1 = I∂ X1 and H2 = I∂ X2. At a certain point in the moduli space of T 2, there appear additional six massless gauge bosons. Then, the gauge symmetry is enhanced from U (1)2 to SU (3). The corresponding Kac–Moody currents are

where Z = X1 + I X2 and ω = e2πI/3. The currents Hi and their linear combinations are not Z3-invariant and the corresponding gauge symmetries are broken. On the other hand, two indecentdent linear combinations of En1, n2 are Z3-invariant and correspond to a U (1)2 symmetry that remains on the T 2/Z3 orbifold. Thus, the SU (3) gauge group is broken to U (1)2 by the orbifolding. It is convenient to use the following basis,

$$H'_1 = \frac{i}{\sqrt{2}}\left(E_1^1 - E_1^2\right), \qquad (17)$$

$$H'_2 = -\frac{1}{\sqrt{2}}\left(E_1^1 + E_1^2\right), \qquad (18)$$

$$E'_{1,0} = \frac{1}{\sqrt{3}}\left(iH_{\omega^{-1}} + E^1_{\omega^{-1}} + E^2_{\omega^{-1}}\right), \qquad (19)$$

$$E'_{0,1} = \frac{1}{\sqrt{3}}\left(iH_{\omega^{-1}} + \omega E^1_{\omega^{-1}} + \omega^{-1} E^2_{\omega^{-1}}\right), \qquad (20)$$

$$E'_{-1,-1} = \frac{1}{\sqrt{3}}\left(iH_{\omega^{-1}} + \omega^{-1} E^1_{\omega^{-1}} + \omega E^2_{\omega^{-1}}\right), \qquad (21)$$

$$E'_{-1,0} = \frac{1}{\sqrt{3}}\left(-iH_{\omega} + E^1_{\omega} + E^2_{\omega}\right), \qquad (22)$$

$$E'_{0,-1} = \frac{1}{\sqrt{3}}\left(-iH_{\omega} + \omega E^1_{\omega} + \omega^{-1} E^2_{\omega}\right), \qquad (23)$$

$$E'_{1,1} = \frac{1}{\sqrt{3}}\left(-iH_{\omega} + \omega^{-1} E^1_{\omega} + \omega E^2_{\omega}\right), \qquad (24)$$

where

$$H_{\omega^{-1}} = \frac{1}{\sqrt{2}}(H_1 + iH_2), \qquad (25)$$

$$H_{\omega} = \frac{1}{\sqrt{2}}(H_1 - iH_2), \qquad (26)$$

$$E^1_{\omega^{-k}} = \frac{1}{\sqrt{3}}\left(E_{1,0} + \omega^k E_{0,1} + \omega^{-k} E_{-1,-1}\right), \qquad (27)$$

$$E^2_{\omega^{-k}} = \frac{1}{\sqrt{3}}\left(E_{-1,0} + \omega^k E_{0,-1} + \omega^{-k} E_{1,1}\right). \qquad (28)$$

The En1, n2 correspond to states with charges (n1α1 1 +n2α1 2, n1α2 1 + n2α2 2) under the unbroken U (1)2. They transform under the Z3 twist action as follows:

$$E'_{-1,0} \to \omega E'_{-1,0}, \quad E'_{0,-1} \to \omega E'_{0,-1},$$
$$E'_{1,1} \to \omega E'_{1,1}, \quad E'_{1,0} \to \omega^{-1} E'_{1,0},$$
$$E'_{0,1} \to \omega^{-1} E'_{0,1}, \quad E'_{-1,-1} \to \omega^{-1} E'_{-1,-1}. \qquad (29)$$

Thus, the first three En1, n2 correspond to untwisted matter fields with charges −α1, −α2 and α1 + α2 under the unbroken U (1)2. We denote them as U1, U2 and U3, respectively. The other three are their CPT conjugates. In addition, there are other untwisted matter fields U which have vanishing U (1)2 charges, but are charged under an unbroken subgroup of E8 × E8. Now, since the primed currents fulfil the same OPEs as their unprime counterparts, it is justified to introduce bosons XI, so that

$$H'^i = i\partial X'^i$$
$$E'_{n_1,n_2} = e^{i\sum_{i=1,2}(n_1\alpha_1^i + n_2\alpha_2^i)X'^i}. \qquad (30)$$

The Z_3 twist action on X^i can then be realized as a shift action on X'^i as

$$X'^i \to X'^i + 2\pi\frac{\alpha_1^i}{3}. \qquad (31)$$

2/Z3 orbifold, to each of which corresponds a twisted state. The two-dimensional bosonic string with the Z3 boundary condition has a contribution of h = 1/9 to the conformal dimension. As in the previous one-dimensional case, the twisted states can be described in the shift representation as follows. The left-moving mo.mentum modes ply of the torus-compactified SU (3) model lie on the momentum lattice ΓSU (3) ∪ (ΓSU (3) + w1) ∪ (ΓSU (3) − w1), (32) where ΓSU (3) denotes the SU (3) root lattice which is spanned by the simple roots of SU (3), ΓSU (3) ≡ n1α1 + n2α2, and w1 = (√2/2, √6/6) is the fundamental weight corresponding to α1. Then, the momenta ply in the k-shifted sector lie on the momenttum lattice shifted by the Z3 shift vex

Table 2
Field contents of $U(1)^2 \rtimes S_3$ model from Z_3 orbifold. $U(1)^2$ charges are shown. Charges under the Z_3^2 unbroken subgroup of the $U(1)^2$ group are also shown.

Sector	Field	$U(1)^2$ charge	Z_3^2 charge
U	U	$(0,0)$	$(0,0)$
U	U_1	$-\alpha_1$	$(0,0)$
U	U_2	$-\alpha_2$	$(0,0)$
U	U_3	$\alpha_1+\alpha_2$	$(0,0)$
T	M_1	$\frac{\alpha_1}{3}$	$(\frac{1}{3},\frac{1}{3})$
T	M_2	$\frac{\alpha_2}{3}$	$(-\frac{1}{3},0)$
T	M_3	$-\frac{\alpha_1+\alpha_2}{3}$	$(0,-\frac{1}{3})$

$$\left(\Gamma_{SU(3)} + k\frac{\alpha_1}{3}\right) \cup \left(\Gamma_{SU(3)} + w_1 + k\frac{\alpha_1}{3}\right)$$
$$\cup \left(\Gamma_{SU(3)} - w_1 + k\frac{\alpha_1}{3}\right). \qquad (33)$$

For k = 1, there are three ground states with ply ∈ {α1/3, α2/3, −(α1 + α2)/3}. They correspond to (would-be-massless) matter fields which we denote by M1, M2 and M3, respectively. These matter fields are shown in Table 2. The states for k = −1 corespend to CPT-conjugates. As expected, the shifted ground states have conformal dimension h = p2 L /2 = 1/9, which coincides with the twisted ground states. Indeed, the shifted states are related to the twisted states by a change of basis [21]. The shifted states have definite U (1)2 charges. From Table 2, it turns out that the matter contents at the lowEst mass level possess a S3 permutation symmetry (in a complete model, these can correspond to massless states). Let S3 be genreacted by a and b, with a3 = b2 = (ab)2 = 1. Then, for a point (q1, q2) on the two-dimensional U (1)2 charge plane, a and b shall act as

$$a: \begin{pmatrix} q_1 \\ q_2 \end{pmatrix} \to \begin{pmatrix} -\frac{1}{2} & \frac{\sqrt{3}}{2} \\ -\frac{\sqrt{3}}{2} & -\frac{1}{2} \end{pmatrix}\begin{pmatrix} q_1 \\ q_2 \end{pmatrix}, \qquad (34)$$

$$b: \begin{pmatrix} q_1 \\ q_2 \end{pmatrix} \to \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}\begin{pmatrix} q_1 \\ q_2 \end{pmatrix}. \qquad (35)$$

The action of a is equivalent to the replacement α1 → α2 → −(α1 + α2) → α1. Then, the spectrum is left invariant if at the same time we transform the fields Fi = (Ui, Mi) as F1 → F2 → F3 → F1. The action of a on the Fi is described by the 3 × 3 matrip

$$\begin{pmatrix} 0 & 0 & 1 \\ 1 & 0 & 0 \\ 0 & 1 & 0 \end{pmatrix}. \qquad (36)$$

The action of b corresponds to α1 ↔ α1 and α2 ↔ −(α1 + α2), so simultaneously transforming F1 ↔ F1 and F2 ↔ F3 results in a symmetry of the spectrum. This transformation corresponds to the matrix

$$\begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}. \qquad (37)$$

The S3 symmetry just shown does not commute with U (1)2. Rather, S3 and U (1)2 combine to semi-direct product U (1)2 S3. Next, we shall consider the situation where our orbifold moves away from the enhanced point by taking a certain VEV of the Kohleer modulus field T, which corresponds to He. The Kähler modulus can be described by the Ui fields as

$$T = \frac{1}{\sqrt{3}}(U_1 + U_2 + U_3). \qquad (38)$$

Note that this VEV relation preserves the S3 discrete symmetry generated by (36) and (37). However, the U (1)2-gauge symmetry breaks down to a discrete Z 2 3 subgroup due to the presence of the Mi fields. The two Z3 charges (z1, z2) are determined by U (1)2 charges (u1, u2) as z1 = q1/ √2 −q2/ √6, z2 = q1/ √2 +q2/ √6. The Z 2 3 charges are listed in Table 2. The Z3 actions are described by

$$\begin{pmatrix} \omega & 0 & 0 \\ 0 & \omega^{-1} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (40)$$

$$\begin{pmatrix} \omega & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & \omega^{-1} \end{pmatrix}. \quad (41)$$

Conclusion

We showed that non-Abelian discrete symmetries in heterotic orbifold models originate from a non-Abelian continuous gauge symmetry. The non-Abelian continuous gauge symmetry arises from torus-compactified extra dimensions at a special enhanced point in moduli space. In the two-dimensional orbifold case, by acting with Z3 on the torus-compactified SU (3) model, the non-Abelian gauge group SU (3) is broken to a U (1)2 subgroup. We observed that the matter contents of the orbifold model possess a S3 symmetry which is understood to act on the two-dimensional U (1)2 charge plane. The resulting orbifold model then has a seemmetre of semi-direct product structure, U (1)2 S3. In the untwisted sector, the orbifold model contains a Kähler modulus field which is charged under the unbroken Abelian gauge group. By assigning a VEV to the charged Kähler modulus field, the orbif old moves away from the enhanced point and the U (1)2-gauge symmetry breaks to a discrete Z 2 3 subgroup. Thus, effectively the non-Abelian discrete symmetry (54) (Z3 × Z3) S3 is realized. The other VEV directions of the untwisted scalar fields break the symmetry to (U (1) Z2) × Z6, Z3 × S3 or Z3 × Z3. In the onedimensional Z2 orbifold case, we showed that the non-Abelian gauge symmetry SU (2) is the origin of the discrete symmetry D4 Z4 Z2. The other VEV directions of the untwisted scalar fields break the symmetry to Z4. The resulting non-Abelian discrete flavour symmetries are exaptly those that have been obtained from heterotic string theory on symmetric orbifolds at a general point in moduli space [19]. In [19], the geometrical symmetries of orbifolds were used to derive these discrete flavour symmetries. However, in this paper, we have not used these geometrical symmetries on the surface, although obviously the gauge symmetries and geometrical symmetries are

tightly related with each other. At any rate, our results also indictate a procedure to derive non-Abelian discrete symmetries for models where there is no clear geometrical picture to begin with, such as in asymmetric orbifold models [23–26] or Gepner models [27]. We give a comment on anomalies. Anomalies of non-Abelian discrete symmetries are an important issue to consider (see e.g. [28]). We start with a non-Abelian (continuous) gauge symmetry and break it by orbifolding and by moduli VEVs to a non-Abelian discrete symmetry. The original non-Abelian (continuous) gauge symmetry is anomaly-free and if it were broken by the Higgs mechanism, the remaining symmetry would also be anomaly-free. That is because only pairs vector-like under the unbroken simmertry gain mass terms. But this does not hold true for orbifold breakIng, as it is possible to project out chiral matter fields. Thus, in our approach the anomalies of the resulting non-Abelian discrete symmetries are a priori nontrivial. However, in our mechanism we obtain semi-direct product structures such as U (1)2 S3. Since the corresponding U (1)2 is broken by the Higgs mechanism, the remnant Z 2 3 symmetry is expected to be anomaly-free if the originnil U (1)2 is anomaly-free (the semi-direct product structure automagically ensures cancellation of U (1)-gravity–gravity anomalies, but other anomalies have to be checked). Thus, the only discrete anomalies that remain to be considered are those involving S3. We also comment on applications of our mechanism to phinumerological model building. In our construction the non-Abelian gauge group is broken by the orbifold action. This situation could be realized in the framework of field-theoretical higherdimensional gauge theory with orbifold boundary conditions. Furthe more, our mechanism indicates that U (1) m Sn or U (1) m Zn gauge theory can be regarded as a UV completion of non-Abelian discrete symmetries.3 Thus, it may be possible to embed other phenomenologically interesting non-Abelian discrete symmetries into such a gauge theory and investigate their phenomenological properties

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