To my late father

Dihedral Flavor Symmetry
from Dimensional Deconstruction

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Abstract
Extra dimension deconstructed on a closed chain has naturally the symmetry of a regular polygon, the dihedral symmetry $D_N$. We assume that the fields are irreducible representations of the binary dihedral group $Q_{2N}$, which is the covering group of $D_N$. It is found that although the orbifold boundary conditions break the dihedral invariance explicitly, the $Q_{2N}$ symmetry appears as an intact internal, global flavor symmetry at low energies. A concrete predictive model based on $Q_{6N}$ with an odd $N$ is given.

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I. INTRODUCTION

The Yukawa sector of the standard model (SM) contains a large number of redundant parameters. The presence of the redundant parameters is not related to a symmetry in the SM. That is, they will appear in higher orders in perturbation theory even if they are set equal to zero at the level. These redundant parameters may become physical parameters when going beyond the SM, and, moreover, they can induce flavor changing neutral currents (FCNCs) and CP violating phenomena that are absent or strongly suppressed in the SM. One of the most well known examples is the case of the minimal supersymmetric model (MSSM). Since the SM can not control the redundant parameters, the size of the new FCNCs and CP violating phases may be unacceptably large unless there is some symmetry, or one fine tunes their values.

A natural guidance to constrain the Yukawa sector and to reduce the redundancy of this sector is a flavor symmetry. It has been recently realized that nonabelian discrete flavor symmetries, especially dihedral symmetries, can not only reduce the redundancy, but also partly explain the large mixing of neutrinos. When supersymmetrized, it has been found that the same flavor symmetries can suppress FCNCs that are caused by soft supersymmetry breaking terms [17, 18] (see also [19]-[23]).

In this letter we address the question of the origin of dihedral flavor symmetries. We will find that dimensional deconstruction [24, 25] is a possible origin of dihedral flavor symmetries.

II. DIHEDRAL INVARIANCE IN AN EXTRA DIMENSIONAL SPACE

Consider an extra dimension which is compactified on a closed one-dimensional lattice with $N$ sites. We assume that the lattice has the form of a regular polygon with $N$ edges as it is illustrated in Fig. 1.

FIG. 1: A regular polygon with $N = 12$ edges, which are located at $y = y_0, y_1, \ldots, y_{N-1}$.

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1 For recent reviews, see, for instance, [1] and references therein.
2 Models based on dihedral flavor symmetries, ranging from $D3(\simeq S_3)$ to $Q_6$ and $D_7$, have been recently discussed in [2]-[16].
The regular polygon is invariant under the symmetry operations of the dihedral group $D_N$. The $D_N$ operations are $2N$ discrete rotations, where $N$ of $2N$ rotations are combined with a parity transformation. Clearly, a discrete polygon rotation of $n \times \theta_N$, $n \in \{1, \ldots, N\}$ corresponds to a discrete translation of the lattice sites of $n \times a$, where $a$ is the lattice spacing and

$$\theta_N \equiv \frac{2\pi}{N}. \quad (1)$$

The coordinate of the extra dimension is denoted by $y$, and the $N$ sites are located at $y = y_0, y_1, \ldots, y_{N-1}$. ($y_{N+i}$ is identified with $y_i$.) Under a $D_N$ transformation, the set of coordinates $(y_0, y_1, \ldots, y_{N-1})$ changes to $(y'_0, y'_1, \ldots, y'_{N-1})$, which we express in terms of a $N \times N$ real matrix. The matrix for the fundamental rotation (i.e., a rotation of $\theta_N$) is given by

$$R_N = \begin{pmatrix} 0 & 0 & 0 & \cdots & 1 \\ 1 & 0 & 0 & \cdots & 0 \\ 0 & 1 & 0 & \cdots & 0 \\ \vdots \\ 0 & 0 & \cdots & 1 & 0 \end{pmatrix}, \quad (2)$$

and that for the parity transformation is

$$P_D = \begin{pmatrix} 1 & 0 & \cdots & 0 & 0 \\ 0 & \cdots & 0 & 0 & 1 \\ 0 & \cdots & 0 & 1 & 0 \\ \vdots \\ 0 & 1 & \cdots & 0 & 0 \end{pmatrix}. \quad (3)$$

Then the $2N$ group elements of $D_N$ are:

$$\mathcal{G}_{D_N} = \{R_N, (R_N)^2, \ldots, (R_N)^N = 1, R_N P_D, (R_N)^2 P_D, \ldots, (R_N)^N P_D = P_D\}. \quad (4)$$

Using the properties, $P_D^2 = 1$ and $P_D R_N P_D = (R_N)^{-1}$, one can convince oneself that $\mathcal{G}_{D_N}$ is indeed a group.

There exist two-dimensional representations for $\tilde{R}_N$ and $\tilde{P}_D$ [12]:

$$\tilde{R}_N = \begin{pmatrix} \cos \theta_N & \sin \theta_N \\ -\sin \theta_N & \cos \theta_N \end{pmatrix}, \quad \tilde{P}_D = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad (5)$$

which are useful representations in finding irreducible representations (irreps) of $D_N$ ($\theta_N$ is given in (1)). It follows that $D_N$ is a subgroup of $SO(3)$, which one sees if one embeds $\tilde{R}_N$ and $\tilde{P}_D$ into $3 \times 3$ matrices

$$\tilde{R}_N \rightarrow \begin{pmatrix} \cos \theta_N & \sin \theta_N & 0 \\ -\sin \theta_N & \cos \theta_N & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad \tilde{P}_D \rightarrow \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & -1 \end{pmatrix}. \quad (6)$$
Therefore, $D_N$ has only real representations.

$SU(2)$ is the universal covering group of $SO(3)$, and has pseudo real and real irreps. $Q_{2N}$ is a finite subgroup of $SU(2)$. It can be interpreted as the covering group of $D_N$ in the sense that the defining matrices $\tilde{R}_{2N}$ and $\tilde{P}_Q$ for $Q_{2N}$ satisfy

\[(\tilde{R}_{2N})^2 = \tilde{R}_N, \quad (\tilde{P}_Q)^4 = (\tilde{P}_D)^2 = 1,\]

where

\[
\tilde{R}_{2N} = \left( \begin{array}{cc}
\cos \frac{\theta_N}{2} & \sin \frac{\theta_N}{2} \\
-\sin \frac{\theta_N}{2} & \cos \frac{\theta_N}{2}
\end{array} \right), \quad \tilde{P}_Q = \left( \begin{array}{cc}
i & 0 \\
0 & -i
\end{array} \right).
\]

The set of $4N$ elements of $Q_{2N}$ is given by

\[G_{Q_{2N}} = \{ \tilde{R}_{2N}, (\tilde{R}_{2N})^2, \ldots, (\tilde{R}_{2N})^{2N} = 1, \tilde{R}_{2N}\tilde{P}_Q, (\tilde{R}_{2N})^2\tilde{P}_Q, \ldots, (\tilde{R}_{2N})^{2N}\tilde{P}_Q = \tilde{P}_Q \}. \]

There exist only one- and two-dimensional irreps for $D_N$ and $Q_{2N}$. For $Q_{2N}$, there are $N-1$ different two-dimensional irreps, which we denote by $2_\ell$, $\ell = 1, \ldots, N-1$.

$2_\ell$ with odd $\ell$ is a pseudo real representation, while $2_\ell$ with even $\ell$ is a real representation, where $2_\ell$ with even $\ell$ is exactly $2_{\ell/2}$ of $D_N$. Under the fundamental rotation (i.e., a rotation of $\theta_N$ which is defined in (11)), $2_\ell$ transforms with the matrix

\[
\tilde{R}_{2N}(2_\ell) = (\tilde{R}_{2N})^\ell = \left( \begin{array}{cc}
\cos(\ell \theta_N / 2) & \sin(\ell \theta_N / 2) \\
-\sin(\ell \theta_N / 2) & \cos(\ell \theta_N / 2)
\end{array} \right).
\]

It is straightforward to calculate the Clebsch-Gordan coefficients for tensor products of irreps [12]. There exist four different one-dimensional irreps of $Q_{2N}$. Because of the relation (7), each of them has a definite $Z_4$ charge. Further, under the fundamental rotation, they either remain unchanged or change their sign. Therefore, one-dimensional irreps can be characterized according to $Z_2 \times Z_4$ charge:

\[
1_{+,0}, \quad 1_{-,0}, \quad 1_{+,2}, \quad 1_{-,2} \quad \text{for} \quad N = 2, 4, 6, \ldots, \quad (12)
\]

\[
1_{+,0}, \quad 1_{-,1}, \quad 1_{+,2}, \quad 1_{-,3} \quad \text{for} \quad N = 3, 5, 7, \ldots, \quad (13)
\]

where the $1_{+,0}$ is the true singlet of $Q_{2N}$, and only $1_{-,1}$ and $1_{-,3}$ are complex irreps. Note that all the real representations of $Q_{2N}$ are exactly those of $D_N$, which is one of the reasons why we would like to call $Q_{2N}$ as the covering group of $D_N$.

### III. FIELD THEORY WITH THE DIHEDRAL INVARINACE

Let us now discuss how to construct field theory models with a dihedral invariance. We denote the five-dimensional coordinate by

\[z^M = (x^\mu, y) \quad \text{with} \quad \mu = 0, \ldots, 3. \]
The coordinates $y_i$ of the lattice sites transform to $y'_i$ with $N \times N$ matrices of $D_N$, which are given in [2] and [3]. Then it is natural to assume that the fields defined on the lattice are irreps of $Q_{2N}$ which is the covering group of $D_N$. That is

$$\phi(x, y) \rightarrow \phi'(x, y) = \tilde{Q}_{2N} \phi(x, \tilde{D}_N^{-1} y), \quad \tilde{Q}_{2N} \in Q_{2N} \text{ and } \tilde{D}_N \in D_N.$$  

(15)

In Table 1 explicit expressions of the matrices corresponding to the fundamental rotation and the parity transformation are given, where we assume that the gauge fields belong to the true singlet $1_{+, 0}$.

| Irreps | $1_{+, 0}$ | $1_{+, 2}$ | $1_{-, 0}$ | $1_{-, 1}$ | $1_{-, 2}$ | $2_{2\ell - 1}$ | $2_{2\ell}$ |
|--------|------------|------------|------------|------------|------------|---------------|------------|
| rotation | 1          | 1          | −1         | −1         | −1         | $(R_{2N})^{2\ell - 1}$ | $(R_{2N})^{2\ell}$ |
| parity     | 1          | −1         | 1          | $i$        | −1         | $P_Q$         | $P_D$      |
| reality     | r          | r          | r          | c          | c          | pr            | r          |

Table 1. Explicit expressions of the matrices corresponding to the fundamental rotation (i.e., a rotation of $\theta_N$ given in [1]) and the parity transformation. $\tilde{R}_{2N}$, $\tilde{P}_Q$ and $\tilde{P}_D$ are given in [8] and [9], respectively, where $\ell \in \mathbb{N}$ and $\leq (N - 1)/2$, $r$=real, $c$=complex, $pr$=pseudo real. All the real irreps of $Q_{2N}$ are those of $D_N$. Complex one-dimensional irreps exist only for $N = 3, 5, 7, \ldots$, while the real one-dimensional irreps $1_{-, 0}$ and $1_{-, 2}$ exist only for $N = 2, 4, 6, \ldots$.

Given the details of the $Q_{2N}$ irreps, it is then straightforward to construct an invariant action [24, 25, 26]. Supersymmetrization can also be straightforwardly done [26].

IV. ORBITFOLD BOUNDARY CONDITIONS AND $Q_{2N}$ FLAVOR SYMMETRY

In the case of a continuous extra dimension, orbifold boundary conditions are used to suppress unnecessary light fields and also to obtain four-dimensional chiral fields. We shall discuss next how an internal $Q_{2N}$ flavor symmetry can appear even if orbifold boundary conditions break the dihedral invariance [15]. Let $\phi(x, y)$ be a generic field which satisfies the periodic boundary condition, $\phi(x, y) = \phi(x, y + Na)$. Then the field $\phi(x, y)$ can be decomposed into the cosine and sine modes:

$$\phi(x, y) = \frac{\phi(x)}{\sqrt{N}} + \sum_{i=1}^{i_{\text{max}}} \phi_{+, i}(x) \cos (k_i y) + \sum_{i=1}^{i'_{\text{max}}} \phi_{-, i}(x) \sin (k_i y),$$  

(16)

where

$$\phi(x) = \frac{1}{\sqrt{N}} \sum_{n=0}^{N-1} \phi(x, y_n),$$  

(17)

$$k_i = \frac{2\pi i}{aN}, i \in \mathbb{N}, \quad i_{\text{max}} = \left\{ \begin{array}{ll} \frac{i'_{\text{max}} + 1}{2} = N/2 - 1 & \text{for } \text{even } N \\ \frac{i'_{\text{max}}}{2} = (N - 1)/2 & \text{for } \text{odd } N \end{array} \right.,$$  

(18)

3 $D_N$ may be understood as a twisted product of $Z_N$ and $Z_2$. Witten [27] has considered this $Z_N$ (the symmetry of the boundary of a deconstructed disc) to solve the triplet-doublet splitting problem in GUTs. 4 Nonabelian discrete family symmetries appearing in extra dimension models of [2, 28], for instance, are not directly related to a symmetry of the extra dimension.
\( \phi(x) \) is the zero mode. As in the continuous case, we can drop the cosine or sine modes by imposing an appropriate boundary condition: Under the parity transformation (3), i.e.,

\[
y_0 \rightarrow y_0' = y_0, \quad y_1 \rightarrow y_1' = y_{N-1}, \quad \ldots, \quad y_i \rightarrow y_i' = y_{N-i}, \quad \ldots,
\]

the zero mode \( \phi(x) \) and the cosine modes are even, while the sine modes are odd.

Since the \( D_N \) transformation mixes the cosine and sine modes, the orbifold boundary conditions break the dihedral invariance explicitly. However, the \( Q_{2N} \) invariant construction of an action discussed in the previous section ensures that the \( Q_{2N} \) invariance remains intact as a global, internal symmetry. This is because there is no derivative with respect to \( y \) is used in the construction. So, the theory with orbifold boundary conditions is invariant under the internal transformation

\[
\phi(x, y) \rightarrow \phi'(x, y) = \tilde{Q}_{2N} \phi(x, y), \quad \tilde{Q}_{2N} \in Q_{2N},
\]

which should be compared with (15). The internal symmetry is nothing but a global flavor symmetry based on \( Q_{2N} \).

V. AN EXAMPLE

In what follows, we would like to discuss a concrete model. One of the successful Ansätze for the quark mass matrices is of a nearest neighbor interaction (NNI) type \[29, 30, 31\]

\[
M = \begin{pmatrix}
0 & C & 0 \\
\pm C & 0 & B \\
0 & B' & A
\end{pmatrix},
\]  

(21)

In [12] it has been proposed to derive the mass matrix (21) solely from a dihedral symmetry, and concluded that two conditions should be met: (i) There should be real as well as pseudo real nonsinglet representations, and (ii) there should be the up and down type Higgs \( SU(2)_L \) doublets (type II Higgs). The smallest finite group that allows both real and pseudo real nonsinglet representations is \( Q_6 \) as we have seen. Further, the Higgs sector of the MSSM fits the desired Higgs structure. Therefore, we assume supersymmetry in four dimensions. The \( D_3(S_3) \) model of [6] with a \( Z_2 \) symmetry in the leptonic sector is one of the most predictive models for the leptonic sector. However, the \( Z_2 \) symmetry in the quark sector is broken, so that the \( Z_2 \) symmetry should be seen as an approximate symmetry in that model. It was found, however, that this leptonic sector can be reproduced in a supersymmetric \( Q_6 \) model without introducing an additional discrete symmetry into the leptonic sector [12]. In Table 2 we write the \( Q_6 \) assignment of the quark and lepton chiral supermultiplets \[^5\]:

| \( Q \) | \( U^c, D^c, L, E^c, N^c, H^u, H^d \) | \( Q_3 \) | \( U^c_3, D^c_3, H^u_3, H^d_3 \) | \( L_3, E^c_3 \) | \( N^c_3 \) |
|---|---|---|---|---|---|
| \( Q_6 \) | \( 2, 2 \) | \( 1_{+2} \) | \( 1_{-1} \) | \( 1_{+0} \) | \( 1_{-3} \) |

\[^5\] The same model exists for \( Q_{2N} \) if \( N \) is odd and a multiple of 3.
Table 2. $Q_6$ assignment of the matter supermultiplets. $Q, Q_3, L, L_3$ and $H^u, H^d, H^c, H^3_d$ stand for $SU(2)_L$ doublets supermultiplets for quarks, leptons and Higgs bosons, respectively. Similarly, $SU(2)_L$ singlet supermultiplets for quarks, charged leptons and neutrinos are denoted by $U^c, U^c_3, D^c, D^c_3, E^c, E^c_3$ and $N^c, N^c_3$. This is an alternative assignment to the one given in the footnote of \cite{12}. The present assignment can more suppress the proton decay \cite{32}. The assignment for the mirror supermultiplets can be simply read off from Table 2.

We impose the following orbifold boundary conditions: All the mirror chiral supermultiplets are odd under the parity transformation \cite{19}. Similarly, the $N = 1$ chiral supermultiplets, which are the $N = 2$ superpartners of the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge supermultiplets, are also odd. It is then clear that the zero modes of the gauge, matter and Higgs supermultiplets coincide with those of the supersymmetric $Q_6$ model of \cite{12}, and hence it is the low energy effective theory. The low energy Yukawa superpotential $W_Y$ is given by

$$W_Y = W_Q + W_L,$$

where

$$W_Q = Y_u^Q Q_3^u U_3^q H_3^u + Y^u U_3^q Q_3^u H_3^u - Y^u U_3^q Q_3^u H_3^u + Y^u U_3^q Q_3^u H_3^u + Y^u U_3^q Q_3^u H_3^u + Y^u U_3^q Q_3^u H_3^u + Y^u U_3^q Q_3^u H_3^u,$$

$$W_L = Y^c f^{IJK} L_1^c E_3^c H_3^c + Y^c L_3^c (H_3^c E_1^c + H_3^c E_2^c) + Y^c (L_1 H_1^c + L_2 H_2^c) E_3^c + Y^c (L_1 H_1^c + L_2 H_2^c) E_3^c + Y^c (L_1 H_1^c + L_2 H_2^c) E_3^c,$$

and $f^{122} = f^{221} = f^{222} = -f^{111} = 1$. In \cite{12} it has been found that by introducing a certain set of gauge singlet Higgs supermultiplets it is possible to construct a Higgs sector in such a way that CP phases can be spontaneously induced. Therefore, all the parameters appearing in the Lagrangian including the soft supersymmetry breaking (SSB) sector are real. Consequently, no $CP$ violating processes induced by SSB terms are possible in this model, satisfying the most stringent experimental constraint coming from the EDM of the neutron and the electron \cite{33,34}. Since the Higgs sector is also $Q_6$ invariant, it is straightforward to derive it from dimensional deconstruction. Consequently, the quark sector contains only 8 real parameters with one independent phase to describe the quark masses and their mixing, and the lepton sector contains only 6 real parameters with one independent phase to describe 12 independent physical parameters. Predictions in the $|V_{ub}| - \sin 2\phi_1$ planes are shown in Fig. 2, while Fig. 3 shows the predictions in the $\sin 2\phi_1 - \phi_3$ planes. As we can see from Fig. 2 and 3, with accurate measurements of the Cabibbo-Kobayashi-Maskawa matrix elements, the predictions could be tested.

The predictions in the leptonic sector are summarized as follows \cite{7}:

1. Inverted neutrino mass spectrum, i.e., $m_{\nu_3} < m_{\nu_1}, m_{\nu_2}$.

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\footnote{The Higgs sector of the model of \cite{12} possesses a permutation symmetry $H_1^{u(d)} \leftrightarrow H_2^{u(d)}$, which ensures the stability of the VEV $H_1^{u(d)} = H_2^{u(d)}$. The resulting mass quark matrices are equivalent to \cite{21}. The lepton sector given in \cite{12} can be obtained by the interchange $1 \leftrightarrow 2$.}

\footnote{Large mixing of neutrinos may be obtained in dimensional deconstruction models in a different mechanism. See, for instance, \cite{33,34}.}
2. \( m_{ee}^2 / \Delta m_{23}^2 = \frac{(1+2r^2_12 + t^4_12 - r^4_12)^2}{4r^2_12 (1+r^2_12)(1+r^2_12 - r^4_12) \cos^2 \phi_\nu} - \tan^2 \phi_\nu \) (\( r = \Delta m_{21}^2/\Delta m_{23}^2, t_{12} = \tan \theta_{12} \)), where \( \phi_\nu \) is an independent phase.

3. \( \sin \theta_{13} \simeq m_e / \sqrt{2} m_\mu \simeq 3.4 \times 10^{-3} \) and \( \tan \theta_{23} \simeq 1 - (m_e / \sqrt{2} m_\mu)^2 = 1 - O(10^{-5}) \).

4. The prediction of \( < m_{ee} > \) is shown in Fig. 4.

FIG. 2: Predictions in the \( |V_{ub}| - \sin 2\phi_1 \) plane. The uncertainties result from those in the quark masses and in \( |V_{us}| \) and \( |V_{cb}| \), where we have used \( |V_{us}| = 0.2240 \pm 0.0036 \) and \( |V_{cb}| = (41.5 \pm 0.8) \times 10^{-3} \) [36]. The vertical and horizontal lines correspond to the experimental values, \( \sin 2\beta(\phi_1) = 0.726 \pm 0.037 \) and \( |V_{ub}| = (36.7 \pm 4.7) \times 10^{-4} \) [37, 38].

FIG. 3: Predictions in the \( \sin 2\phi_1 - \phi_3 \) plane. The vertical and horizontal lines correspond to the experimental values, \( \sin 2\beta(\phi_1) = 0.726 \pm 0.037 \) and \( \phi_3 = (60^\circ \pm 14^\circ) \) [37, 38].

FIG. 4: The effective Majorana mass \( < m_{ee} > \) as a function of \( \sin \phi_\nu \) with \( \sin^2 \theta_{12} = 0.3 \) and \( \Delta m_{21}^2 = 6.9 \times 10^{-5} \text{ eV}^2 \) [39]. The dashed, solid and dot-dashed lines stand for \( \Delta m_{23}^2 = 1.4, 2.3 \) and \( 3.0 \times 10^{-3} \text{ eV}^2 \), respectively.
We emphasize that the smallness of $\sin \theta_{13}$ and the almost maximal mixing of the atmospheric neutrinos are consequences of the $Q_6$ flavor symmetry. $\sin \theta_{13}$ in the present model may be too small to be measured in a laboratory experiment [41], but the tiny deviation from zero ($\sin^2 \theta_{13} \approx m_e^2/2m_\mu^2 \approx 10^{-5}$) are important in supernova neutrino oscillations [40].

VI. CONCLUSION

In this letter we have looked for a possible origin of dihedral symmetries. It has been recently realized that a flavor symmetry based on a dihedral group can be used to soften the flavor problem of the SM and the MSSM. We have considered an extra dimension compactified on a closed chain, which is assumed to have the form of a regular polygon. Since the symmetry group of the regular polygon is the dihedral group $D_N$, we assumed that the fields are irreps of the covering group of $D_N$, which is the binary dihedral group $Q_{2N}$. The construction of an action with the dihedral invariance is straightforward, and moreover we found that the $Q_{2N}$ symmetry remains as an intact, internal flavor symmetry even if the original dihedral invariance is broken by orbifold boundary conditions. We hope that with our finding we can come closer to a deep understanding of the origin of a flavor symmetry based on a nonabelian finite group.

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