Neutrino Mass Matrix from $S_4$ Symmetry

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Abstract

The cubic symmetry $S_4$ contains $A_4$ and $S_3$, both of which have been used to study neutrino mass matrices. Using $S_4$ as the family symmetry of a complete supersymmetric theory of leptons, it is shown how the requirement of breaking $S_4$ at the seesaw scale without breaking supersymmetry enforces a special form of the neutrino mass matrix which exhibits maximal $\nu_\mu - \nu_\tau$ mixing as well as zero $U_{e3}$. In addition, $(\nu_e + \nu_\mu + \nu_\tau)/\sqrt{3}$ is naturally close to being a mass eigenstate, thus predicting $\tan^2 \theta_{12}$ to be near but not equal to 1/2.
In atmospheric neutrino oscillation data \[1\], the persistence of maximal $\nu_\mu - \nu_\tau$ mixing has raised the important theoretical question of whether it is due to an underlying symmetry. The naive response is that it is due to the exchange of $\nu_\mu$ with $\nu_\tau$, but since $\nu_\mu$ and $\nu_\tau$ are respective members of the $SU(2)_L$ doublets $(\nu_\mu, \mu)$ and $(\nu_\tau, \tau)$, this symmetry automatically implies the exchange of $\mu$ with $\tau$. As such, it cannot be sustained in the full Lagrangian, because the assumed diagonality of the charged-lepton mass matrix (with $m_\mu \neq m_\tau$) is then lost.

A related issue is whether or not $U_{e3} = 0$. If $\nu_\mu - \nu_\tau$ exchange were a legitimate symmetry, then this would also be “predicted”. On the other hand, $U_{e3} = 0$ by itself can be supported by a different symmetry of the full Lagrangian \[2\], although the latter is generally unable to shed any light on how maximal $\nu_\mu - \nu_\tau$ mixing could be achieved.

With the implementation of non-Abelian discrete symmetries such as $S_3$ \[3\], $D_4$ \[4\], and $A_4$ \[5\], as well as $Q_4$ \[6\] and $Q_6$ \[7\], it is indeed possible to have both maximal $\nu_\mu - \nu_\tau$ mixing and zero $U_{e3}$, as well as a prediction of the mixing angle in solar neutrino oscillations \[8\] in some cases. However, specific ad hoc assumptions regarding the symmetry breaking sector must be made, usually with the addition of arbitrary auxiliary Abelian discrete symmetries \[9\].

In this paper, the problem is solved by invoking a very simple requirement. The complete theory of leptons is assumed to be supersymmetric with $S_4$ as its family symmetry. Neutrino masses are assumed to come from the canonical seesaw mechanism \[10\] with heavy singlet neutral fermions $N$. The key is to allow $S_4$ to be broken by the Majorana mass matrix of $N$ at the seesaw scale, but not the supersymmetry. This requirement then fixes the pattern of symmetry breaking and thus the form of $\mathcal{M}_N$, and subsequently also $\mathcal{M}_\nu$, as shown below.

The group of the permutation of four objects is $S_4$ \[11\]. It is also the symmetry group of the hexahedron, i.e. the cube, one of five (and only five) perfect geometric solids, which
was identified by Plato with the element “earth” [12]. It has 24 elements divided into 5
equivalence classes, with $1, 1', 2, 3,$ and $3'$ as its 5 irreducible representations. Its character
table is given below.

Table 1: Character table of $S_4$.

| Class | $n$ | $h$ | $\chi_1$ | $\chi_1'$ | $\chi_2$ | $\chi_3$ | $\chi_3'$ |
|-------|-----|-----|-----------|------------|-----------|-----------|------------|
| $C_1$ | 1   | 1   | 1         | 1          | 2         | 3         | 3          |
| $C_2$ | 3   | 2   | 1         | 1          | 2         | -1        | -1         |
| $C_3$ | 8   | 3   | 1         | 1          | -1        | 0         | 0          |
| $C_4$ | 6   | 4   | 1         | -1         | 0         | -1        | 1          |
| $C_5$ | 6   | 2   | 1         | -1         | 0         | 1         | -1         |

The two three-dimensional representations differ only in the signs of their $C_4$ and $C_5$
matrices. Their group multiplication rules are similar to those of $A_4$ [13], namely

\[
\begin{align*}
3 \times 3 &= 1 + 2 + 3_S + 3'_A, \\
3' \times 3' &= 1 + 2 + 3_S + 3'_A, \\
3 \times 3' &= 1' + 2 + 3'_S + 3_A,
\end{align*}
\]

where the subscripts $S$ and $A$ refer to their symmetric and antisymmetric product combina-
tions respectively. The two-dimensional representation behaves exactly as its $S_3$ counterpart,

namely

\[
2 \times 2 = 1 + 1' + 2.
\]

The three $N$’s are assigned to the $3$ representation. To obtain a nontrivial $\mathcal{M}_N$, the
singlet Higgs superfields $\sigma_{1,2,3} \sim 3$ and $\zeta_{1,2} \sim 2$ are assumed. Consequently in the $N_{1,2,3}$
basis,

\[
\mathcal{M}_N = \begin{pmatrix}
A + f(\langle \zeta_2 \rangle + \langle \zeta_1 \rangle) & h\langle \sigma_3 \rangle & h\langle \sigma_2 \\
h\langle \sigma_3 \rangle & A + f(\omega \langle \zeta_2 \rangle + \omega^2 \langle \zeta_1 \rangle) & h\langle \sigma_1 \\
h\langle \sigma_2 \rangle & h\langle \sigma_1 \rangle & A + f(\omega^2 \langle \zeta_2 \rangle + \omega \langle \zeta_1 \rangle)
\end{pmatrix},
\]

3
where $\omega = \exp(2\pi i/3)$. The most general $S_4$-invariant superpotential of $\sigma$ and $\zeta$ is given by

$$W = \frac{1}{2}M(\sigma_1\sigma_1 + \sigma_2\sigma_2 + \sigma_3\sigma_3) + \lambda \sigma_1\sigma_2\sigma_3 + m\zeta_1\zeta_2 + \frac{1}{3}\rho(\zeta_1\zeta_1\zeta_1 + \zeta_2\zeta_2\zeta_2) + \frac{1}{2}\kappa(\sigma_1\sigma_1 + \omega\sigma_2\sigma_2 + \omega^2\sigma_3\sigma_3)\zeta_2 + \frac{1}{2}\kappa(\sigma_1\sigma_1 + \omega^2\sigma_2\sigma_2 + \omega\sigma_3\sigma_3)\zeta_1.$$

The resulting scalar potential is then

$$V = |M\sigma_1 + \lambda\sigma_2\sigma_3 + \kappa\sigma_1(\zeta_2 + \zeta_1)|^2$$
$$+ |M\sigma_2 + \lambda\sigma_3\sigma_1 + \kappa\sigma_2(\omega\zeta_2 + \omega^2\zeta_1)|^2$$
$$+ |M\sigma_3 + \lambda\sigma_1\sigma_2 + \kappa\sigma_3(\omega^2\zeta_2 + \omega\zeta_1)|^2$$
$$+ |m\zeta_1 + \rho\zeta_2\zeta_2 + \frac{1}{2}\kappa(\sigma_1\sigma_1 + \omega\sigma_2\sigma_2 + \omega^2\sigma_3\sigma_3)|^2$$
$$+ |m\zeta_2 + \rho\zeta_1\zeta_1 + \frac{1}{2}\kappa(\sigma_1\sigma_1 + \omega^2\sigma_2\sigma_2 + \omega\sigma_3\sigma_3)|^2.$$  

(6)

For supersymmetry to be unbroken, $V_{\text{min}} = 0$ is required. This is possible only if $\langle \zeta_1 \rangle = \langle \zeta_2 \rangle$ and $\langle \sigma_2 \rangle = \langle \sigma_3 \rangle$, for which

$$M\langle \sigma_1 \rangle + \lambda\langle \sigma_2 \rangle^2 + 2\kappa\langle \sigma_1 \rangle\langle \zeta_1 \rangle = 0,$$

(8)

$$M + \lambda\langle \sigma_1 \rangle - \kappa\langle \zeta_1 \rangle = 0,$$

(9)

$$m\langle \zeta_1 \rangle + \rho\langle \zeta_1 \rangle^2 + \frac{1}{2}\kappa(\langle \sigma_1 \rangle^2 - \langle \sigma_2 \rangle^2) = 0,$$

(10)

where $\omega + \omega^2 = -1$ has been used. Thus $\mathcal{M}_N$ is fixed to be of the form

$$\mathcal{M}_N = \begin{pmatrix} A + 2B & C & C \\ C & A - B & D \\ C & D & A - B \end{pmatrix},$$

(11)

where $B = f\langle \zeta_1 \rangle$, $C = h\langle \sigma_2 \rangle$, and $D = h\langle \sigma_1 \rangle$.

The residual symmetry of the theory is then $Z_2$, under which $(N_2 - N_3)$, $(\sigma_2 - \sigma_3)$, $(\zeta_1 - \zeta_2)$ are odd, and $N_1$, $(N_2 + N_3)$, $\sigma_1$, $(\sigma_2 + \sigma_3)$, $(\zeta_1 + \zeta_2)$ are even. If $\langle \zeta_{1,2} \rangle = 0$, then $B = 0$, $C = D$, and the residual symmetry is $S_3$, under which $(N_1 + N_2 + N_3)$, $(\sigma_1 + \sigma_2 + \sigma_3) \sim \mathbb{1}$;
\[ (N_1 + \omega N_2 + \omega^2 N_3, N_1 + \omega^2 N_2 + \omega N_3, (\sigma_1 + \omega \sigma_2 + \omega^2 \sigma_3, \sigma_1 + \omega^2 \sigma_2 + \omega \sigma_3, (\zeta_1, \zeta_2) \sim 2. \] The \( Z_2 \) (and the approximate \( S_3 \)) symmetry of \( M_N \) is preserved in \( M_N^{-1} \).

Consider now the leptons \((\nu_i, l_i)\) and \(l_j^c\) of the Standard Model. They are also assigned to the 3 representation of \( S_4 \). As for the Higgs doublet superfields, they are assumed to be \( 1 + 2 \), with one set coupling \((\nu_i, l_i)\) to \( l_j^c\) and the other coupling \((\nu_i, l_i)\) to \( N_j\), as in any supersymmetric extension of the Standard Model. Nonzero vacuum expectation values of the scalar components of these Higgs superfields break the electroweak symmetry, but because of their \( S_4 \) structure, the \( l_i l_j^c\) mass matrix is diagonal, as well as that of \( \nu_i N_j\). In particular,

\[
\begin{pmatrix}
m_e \\
m_\mu \\
m_\tau
\end{pmatrix} = \begin{pmatrix}
1 & 1 & 1 \\
1 & \omega & \omega^2 \\
1 & \omega^2 & \omega
\end{pmatrix}
\begin{pmatrix}
y_1 v_1 \\
y_2 v_2 \\
y_2 v_3
\end{pmatrix},
\]

where \( v_1 \) is the vacuum expectation value of the 1 representation, \( v_{2,3} \) are those of the 2, and \( y_{1,2} \) are their respective Yukawa couplings. In the \( \nu N \) sector, the vacuum expectation values of the corresponding Higgs 2 representation are assumed to be zero, thus \( m_1 = m_2 = m_3(= m_D) \) for all the Dirac neutrino masses. Using the seesaw mechanism, the observed Majorana neutrino mass matrix in the \( \nu_{e,\mu,\tau} \) basis is then given by

\[
\mathcal{M}_\nu = -\mathcal{M}_D \mathcal{M}_N^{-1} \mathcal{M}_D^T = \begin{pmatrix}
a + 2b & c & c \\
c & a - b & d \\
c & d & a - b
\end{pmatrix},
\]

where

\[
a = [-A^2 + B^2 + (2C^2 + D^2)/3]m_D^2/det \mathcal{M}_N, \\
b = [B(A - B) - (C^2 - D^2)/3]m_D^2/det \mathcal{M}_N, \\
c = C(A - B - D)m_D^2/det \mathcal{M}_N, \\
d = [D(A + 2B) - C^2]m_D^2/det \mathcal{M}_N, \\
det \mathcal{M}_N = (A - B - D)[(A + 2B)(A - B + D) - 2C^2].
\]
This form of $M_\nu$ is precisely that advocated in Ref. [14] on purely phenomenological grounds. As expected, the $S_3$ limit is obtained if $B = 0$ and $C = D$, for which $b = 0$ and $c = d$, resulting in $(\nu_e + \nu_\mu + \nu_\tau)/\sqrt{3}$ as a mass eigenstate. Note that $S_3$ as well as the residual $Z_2$ symmetry are assumed broken softly at the scale of supersymmetry breaking. This allows the electroweak vacuum expectation values to be chosen as they are.

In the basis $[\nu_e, (\nu_\mu + \nu_\tau)/\sqrt{2}, (-\nu_\mu + \nu_\tau)/\sqrt{2}]$, the neutrino mass matrix of Eq. (13) becomes

$$M_\nu = \begin{pmatrix} a + 2b & \sqrt{2}c & 0 \\ \sqrt{2}c & a - b + d & 0 \\ 0 & 0 & a - b - d \end{pmatrix},$$

which exhibits maximal $\nu_\mu - \nu_\tau$ mixing and zero $U_{e3}$ as advertised. The mixing angle of the $2 \times 2$ submatrix (for all parameters real) can be simply read off as

$$\tan 2\theta_{12} = \frac{2\sqrt{2}c}{d - 3b} = \frac{2\sqrt{2}c}{c - 3b - (c - d)},$$

which reduces to $2\sqrt{2}$ in the $S_3$ limit of $b = 0$ and $c = d$. This would imply $\tan^2 \theta_{12} = 1/2$, resulting in a mixing pattern proposed some time ago [15]. However in this limit, $\Delta m_{atm}^2$ vanishes as well [16]. Therefore, $\tan^2 \theta_{12}$ should not be equal to $1/2$, but since it is natural for $b$ and $c - d$ to be small compared to $c$, its deviation from $1/2$ is expected to be small. For example,

$$(3b + c - d)/c = -0.15 \implies \tan^2 \theta_{12} = 0.45,$$

in excellent agreement with data [8].

The limit $\Delta m_{sol}^2 = 0$ implies $2a + b + d = 0$, and

$$\Delta m_{atm}^2 \equiv m_3^2 - (m_2^2 + m_1^2)/2 = (a - b - d)^2 - (d - 3b)^2/4 - 2c^2 = 6bd - 2(c^2 - d^2) \simeq [6b - 4(c - d)]c,$$

which can be either positive or negative, corresponding to a normal or inverted ordering of neutrino masses respectively. If $c - d = 0$, then $\tan^2 \theta_{12} < 1/2$ implies an inverted ordering
as in the models of Ref. [16]. It is also possible to set \( B = 0 \) alone by choosing \( f = 0 \) in Eq. (5), so that Eq. (20) becomes

\[
\tan 2\theta_{12} = \frac{2\sqrt{2}C}{D},
\]

and \( \Delta m^2_{sol} = 0 \) implies

\[
\Delta m^2_{atm} \simeq 9(C - D)C^3(m_D^2/det\mathcal{M}_N)^2.
\]

Thus \( \tan^2 \theta_{12} < 1/2 \) would also imply an inverted ordering of neutrino masses in this case.

The effective neutrino mass \( m_{ee} \) measured in neutrinoless double beta decay is simply given by the magnitude of the \( \nu_e\nu_e \) entry of \( \mathcal{M}_\nu \), i.e. \( |a + 2b| \). Using Eqs. (21), (22), and \( |\Delta m^2_{atm}| = 2.5 \times 10^{-3} \text{ eV}^2 \), it is in the range 0.10 eV for \( c - d = 0 \) (inverted ordering) and 0.05 eV (normal ordering) for \( b = 0 \).

Returning to the condition \( 2a + b + d = 0 \) for \( \Delta m^2_{sol} = 0 \), under which \( \mathcal{M}_\nu \) of Eq. (13) becomes

\[
\mathcal{M}_\nu = \begin{pmatrix} a + 2b & c & c \\ c & a - b & -2a - b \\ c & -2a - b & a - b \end{pmatrix},
\]

it should be noted that this form is invariant under the transformation [17]

\[
UM\nu U^T = \mathcal{M}_\nu,
\]

where

\[
U = \begin{pmatrix} i2\sqrt{2}/3 & i/3\sqrt{2} \\ i/3\sqrt{2} & 1/2 - i\sqrt{2}/3 \\ i/3\sqrt{2} & -1/2 - i\sqrt{2}/3 \end{pmatrix}, \quad U^2 = \begin{pmatrix} -1 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & -1 & 0 \end{pmatrix},
\]

if \( c = -2a - 4b \), for which \( \tan^2 \theta_{12} = 1/2 \). This limit of \( \mathcal{M}_\nu \) corresponds thus to a \( Z_4 \) symmetry.

In conclusion, the family symmetry \( S_4 \) has been advocated as the origin of the observed pattern of neutrino mixing in the context of a complete supersymmetric theory. The key
is the requirement that supersymmetry is unbroken at the seesaw scale where $S_4$ is broken. This fixes the pattern of $S_4$ breaking to retain a residual $Z_2$ (as well as an approximate $S_3$) symmetry in the Majorana mass matrix of the heavy singlet neutrinos. At the much lower scale of supersymmetry breaking, $S_4$ as well as $Z_2$ are allowed to be broken by soft terms. However, the $S_4$-invariant Yukawa terms ensure that the charged-lepton mass matrix is diagonal, and the neutrino Dirac masses are equal. This results in a Majorana neutrino mass matrix which exhibits maximal $\nu_\mu - \nu_\tau$ mixing and zero $U_{e3}$. In the $S_3$ limit, $(\nu_e + \nu_\mu + \nu_\tau)/\sqrt{3}$ is a mass eigenstate, thus predicting $\tan^2 \theta_{12} = 1/2$, but $\Delta m^2_{atm} = 0$ as well. Allowing a small deviation from this limit can result in $\tan^2 \theta_{12} = 0.45$ and a nonzero $\Delta m^2_{atm}$, in excellent agreement with data.

I thank Michele Frigerio for an important comment. This work was supported in part by the U. S. Department of Energy under Grant No. DE-FG03-94ER40837.
Appendix The matrices of the 3 representation of $S_4$ are given by

$$C_1 : \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (28)$$

$$C_2 : \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & -1 \end{pmatrix}, \begin{pmatrix} -1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -1 \end{pmatrix}, \begin{pmatrix} -1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (29)$$

$$C_3 : \begin{pmatrix} 0 & 0 & 1 \\ 1 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \begin{pmatrix} 0 & 1 & 0 \\ -1 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 1 \\ -1 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 0 & 1 \\ 1 & 0 & 0 \\ 0 & 1 & 0 \end{pmatrix}, \quad (30)$$

$$C_4 : \begin{pmatrix} -1 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & 0 & 1 \end{pmatrix}, \begin{pmatrix} 0 & -1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}, \begin{pmatrix} -1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & -1 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 0 & 1 \\ 0 & -1 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 1 & 0 \\ -1 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}, \quad (31)$$

$$C_5 : \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \begin{pmatrix} 0 & 0 & 1 \\ 0 & 1 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & 1 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 0 & -1 \\ 0 & 1 & 0 \\ -1 & 0 & 0 \end{pmatrix}, \begin{pmatrix} 0 & -1 & 0 \\ -1 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (32)$$

The matrices of the 3' representation are the same as those of 3 for $C_{1,2,3}$ and opposite in sign for $C_{4,5}$. Those of the 2 representation are the same as in $S_3$, i.e.

$$C_{1,2} : \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}, \quad (33)$$

$$C_3 : \begin{pmatrix} \omega & 0 \\ 0 & \omega^2 \end{pmatrix}, \begin{pmatrix} \omega^2 & 0 \\ 0 & \omega \end{pmatrix}, \quad (34)$$

$$C_{4,5} : \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \begin{pmatrix} 0 & \omega \\ \omega^2 & 0 \end{pmatrix}, \begin{pmatrix} 0 & \omega^2 \\ \omega & 0 \end{pmatrix}, \quad (35)$$

each appearing four times.
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