Neutrino phenomenology and stable dark matter with $A_4$

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We present a model based on the $A_4$ non-abelian discrete symmetry leading to a predictive five-parameter neutrino mass matrix and providing a stable dark matter candidate. We found an interesting correlation among the atmospheric and the reactor angles which predicts $\theta_{23} \sim \pi/4$ for very small reactor angle and deviation from maximal atmospheric mixing for large $\theta_{13}$. Only normal neutrino mass spectrum is possible and the effective mass entering the neutrinoless double beta decay rate is constrained to be $|m_{ee}| > 4 \cdot 10^{-4} \text{eV}$.

I. INTRODUCTION

Neutrino oscillation and dark matter (DM) are so far the most important evidences of physics beyond the standard model. Many models for neutrino masses and mixings have been studied in literature, some of them with the aim of reproducing the tri-bimaximal mixing pattern (TBM) observed, within one standard deviation, in the recent neutrino oscillation data, namely $\sin^2 \theta_{12} = 1/3$, $\sin^2 \theta_{23} = 1/2$ and $\sin^2 \theta_{13} = 0$. In this context, non-Abelian discrete flavor symmetries have been largely used since, especially in models based on $A_4$ and $S_4$ permutation groups, the TBM limit can be easily reproduced.

While we have robust evidence for dark matter from rotation curves of spiral galaxies, gravitational lensing, WMAP measurement, CMB anisotropy, structure formation, X-ray observations, bullet-clusters, we still do not have neither theoretical nor experimental indications about the nature of it; also the mechanism for DM stability is still not understood yet. Many models assume ad hoc extra Abelian symmetries in order to stabilize the dark matter. Such a symmetry can arise from the spontaneous breaking of a non-Abelian continuous symmetry, see for instance, or from the breaking of Abelian $U(1)$ group as, for example, in grand unified frameworks or in supersymmetry.

Recently a model where the stability of the DM arises from the spontaneous breakdown of a non-abelian discrete flavor symmetry was proposed. Such a model is based on the property that the $A_4$ group is spontaneously broken into its subgroup $Z_2$ which is responsible for the stability (for a paper with decaying dark matter with discrete non-Abelian symmetry see). In the same $Z_2$ is also acting in the neutrino sector, giving a vanishing reactor angle $\theta_{13} = 0$ and allowing only the inverted hierarchy for the neutrino mass spectrum.

However these properties are model-dependent features and cannot be considered as general results of

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models where the stability of the dark matter is justified by non-abelian discrete flavor symmetries. In this paper we want to provide an explicit example of a model, based on a simple extension of [13], with a richer neutrino phenomenology, predicting a normal mass ordering and $\theta_{13} \neq 0$.

The paper is organized as follows: in section II we give the field content of our model and derive the neutrino mass matrix, whose phenomenological implications are discussed in section III in section IV we draw our conclusions.

II. THE MODEL

In Tab. I we summarized the model quantum numbers. In contrast to [13], the right-handed neutrino $N_4$ is assigned to $1'$ instead of $1$ and we introduced one more right-handed neutrino $N_5$ assigned to $1''$.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|c|c|c|}
\hline
 & $L_e$ & $L_\mu$ & $L_\tau$ & $l'_e$ & $l'_\mu$ & $l'_\tau$ & $N_T$ & $N_4$ & $N_5$ & $H$ & $\eta$ \\
\hline
$SU(2)$ & 2 & 2 & 2 & 1 & 1 & 1 & 1 & 1 & 1 & 2 & 2 \\
$A_4$ & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 3 & 1 & 1' & 1'' & 1 & 3 \\
\hline
\end{tabular}
\caption{Summary of relevant model quantum numbers}
\end{table}

The operators needed to generate neutrinos masses are the following:

$$\mathcal{L} = y_1^\nu L_e (N_T \eta)_1 + y_2^\nu L_\mu (N_T \eta)_1 + y_3^\nu L_\tau (N_T \eta)_1 + y_4^\nu L_\tau N_4 H + y_5^\nu L_\mu N_5 H + M_1 N_T N_T + M_2 N_4 N_5 + \text{h.c.}$$

The scalar and the lepton charged current sectors of our model are the same as in [13] and we refer to that paper for all the details. After electroweak symmetry breaking, two of the Higgs doublets contained in the $A_4$ triplet $\eta$ do not take vev and we have:

$$\langle H^0 \rangle = v_h \neq 0, \quad \langle \eta^0_1 \rangle = v_\eta \neq 0, \quad \langle \eta^0_{2,3} \rangle = 0.$$  \hspace{1cm} (1)

As a consequence, the Dirac and Majorana neutrino mass matrices are:

$$m_D = \begin{pmatrix}
    y_1^\nu v_\eta & 0 & 0 & 0 \\
    y_2^\nu v_\eta & 0 & 0 & y_4^\nu v_h \\
    y_3^\nu v_\eta & 0 & y_4^\nu v_h & 0
\end{pmatrix}, \quad m_M = \begin{pmatrix}
    M_1 & 0 & 0 & 0 & 0 \\
    0 & M_1 & 0 & 0 & 0 \\
    0 & 0 & M_1 & 0 & 0 \\
    0 & 0 & 0 & 0 & M_2 \\
    0 & 0 & 0 & M_2 & 0
\end{pmatrix}.$$  \hspace{1cm} (2)

Using the See-Saw type-I formula we get the following light neutrino mass matrix:

$$m_\nu = \begin{pmatrix}
    a^2 & ab & ac \\
    ab & b^2 & bc + k \\
    ac & bc + k & c^2
\end{pmatrix},$$  \hspace{1cm} (3)

where we defined:

$$a = \frac{y_1^\nu v_\eta}{\sqrt{M_1}} \quad b = \frac{y_2^\nu v_\eta}{\sqrt{M_1}} \quad c = \frac{y_3^\nu v_\eta}{\sqrt{M_1}} \quad k = \frac{y_4^\nu y_5^\nu v_h^2}{M_2}.$$  \hspace{1cm} (4)
As discussed in more details in \cite{13}, the group $A_4$ is broken by the vev $\langle \eta_i \rangle \sim (1, 0, 0)$ down to the subgroup $Z_2$, generated by the diagonal $A_4$ generator $S = \text{Diag}\{1, -1, -1\}$; the $Z_2$ symmetry acts on the triples fields in the following way:

$$
\begin{align*}
Z_2 : & \\
N_2 & \to -N_2, \quad h_2 \to -h_2, \quad A_2 \to -A_2, \\
N_3 & \to -N_3, \quad h_3 \to -h_3, \quad A_3 \to -A_3,
\end{align*}
$$

(5)

where $h_i$ and $A_i$ are respectively the CP-odd and CP-even component of the Higgs doublet $\eta_i$ for $i = 2, 3$ and $N_{2,3}$ are the component of the triplet $N_T$. This residual symmetry is responsible for the stability of the lightest combination of $h_2$ and $h_3$, which is the dark matter candidate. In fact it couples only to heavy right-handed neutrinos and not to quarks, supposed to be singlets under $A_4$. Such a scalar dark matter candidate is potentially detectable in nuclear recoil experiments \cite{17, 18}.

\section{Numerical Analysis of the Full Mass Matrix}

The most general complex and symmetric matrix has twelve independent real parameters which, after readsoeing the unphysical phases, reduce to nine. The matrix in eq. (3) has four complex parameters and then five real independent parameters: the moduli of $a, b$ and $c$ and the modulus and phase of the combination $bc + k = d e^{i \phi_d}$. In terms of them, the matrix in eq. (3) is rewritten in the following form:

$$
m_\nu = \begin{pmatrix} a^2 & ab & ac \\
ab & b^2 & de^{i \phi_d} \\
ac & de^{i \phi_d} & c^2 \end{pmatrix}.
$$

(6)

We can relate three of the five parameters in eq. (6) (i.e. $a, b$ and $c$) with the neutrino masses using the invariant quantities of the hermitian matrix $M_\nu^2 = m_\nu\, m_\nu^\dagger$:

$$
\begin{align*}
\text{Tr} (M_\nu^2) &= T = a^4 + 2a^2 (b^2 + c^2) + b^4 + c^4 + 2d^2 = \\
&= m_{\nu_1}^2 + m_{\nu_2}^2 + m_{\nu_3}^2 \\
\text{det} (M_\nu^2) &= a^4 (b^2 c^2 - 2b c d \cos \phi_d + d^2)^2 = \\
&= m_{\nu_1}^2 m_{\nu_2}^2 m_{\nu_3}^2 \\
\frac{1}{2} [T^2 - \text{Tr}(M_\nu^2 M_\nu^\dagger)] &= (b^2 c^2 - 2b c d \cos \phi_d + d^2) \cdot \\
&= 2a^4 + 2a^2 (b^2 + c^2) + b^4 c^2 + 2b c d \cos \phi_d + d^2 \\
&= m_{\nu_1}^2 (m_{\nu_2}^2 + m_{\nu_3}^2) + m_{\nu_2}^2 m_{\nu_3}^2,
\end{align*}
$$

(7)

where the symbols Tr and Det refer to the trace and determinant of a matrix, respectively. The expressions of $a, b$ and $c$ in terms of neutrino masses, derived from the inversion of the previous relations, produce lengthly formulae that we do not present here but included in our numerical simulations. As independent parameters, we decide to take $m_{\nu_1}$, the two square mass differences $\Delta m_{\text{sol}}^2$ and $\Delta m_{\text{atm}}^2$, and the other two parameters of the neutrino mass matrix, i.e. $d$ and $\phi_d$; to study the correlations among the neutrino observables predicted by the model, we perfom a parameter space scan extracting randomly the previous parameters in the following.
intervals:

\[ d \in [-1, 1], \quad \phi_d \in [-\pi, \pi), \quad m_{\nu_1} \in [10^{-5}, 1] \text{ eV}, \]

\[ 7.1 \times 10^{-5} \text{ eV}^2 < \Delta m^2_{\text{sol}} < 8.3 \times 10^{-5} \text{ eV}^2, \]

\[ 0.027 < r = \frac{\Delta m^2_{\text{sol}}}{|\Delta m^2_{\text{atm}}|} < 0.040. \]

Then, we require that the obtained mixing angles are within their current 3\(\sigma\) confidence level \[2\]. We carefully checked that \(|d| > 1\) is not compatible with data. We made the calculation for both inverted and normal neutrino mass hierarchies finding that only the latter is allowed in our framework (we have analytically discussed this point in the appendices: Appendix A1 for the real case, Appendix A2 for the the \(\mu - \tau\) invariant case and Appendix A3 for the exact TBM case). Among all possible correlations, we found that the \((\theta_{13} - \theta_{23})\) one is quite interesting, as shown in fig.(1). As we can see, for very small \(\theta_{13}\) the atmospheric angle is fixed to its maximal value, as a consequence of the fact that the model has the TBM limit built-in (see next section). Increasing \(\theta_{13}\) two different branches above and below \(\theta_{23} = \pi/4\) develop and for \(\theta_{13} \gtrsim 1^\circ\) maximal 2 – 3 mixing is strongly excluded. It is interesting to observe that this result can be easily verified at the forthcoming neutrino experiments \[21\] even with a reduced sensitivity to \(\theta_{13}\) since the largest deviation from \(\theta_{23} = \pi/4\) is obtained for \(\theta_{13}\) close to its current upper bound. We checked that the previous correlation is obtained with \(a, b, c\) and \(d\) all at the same order of magnitude \(O(10^{-2})\), which ensures that no fine-tuning is required among the matrix elements of eq.(3), and with the phase \(\phi\) uniformly distributed in the \([-\pi, \pi)\) interval. We also verified that the model does not produce any other relevant correlations among the mixing parameters.

The other interesting neutrino observable that we want to discuss is the effective mass \(|m_{ee}|\) entering the neutrinoless double beta decay rate. Since our model is only compatible with the normal hierarchy, we expect it to be quite small for small \(m_{\nu_1}\). The result of our simulation can be found in fig.(2), where we can see that a lower bound \((|m_{\nu_1}|, |m_{ee}|) \gtrsim (2 \cdot 10^{-3}, 4 \cdot 10^{-4}) \text{ eV}\) can be set. The existence of such a lower
bound on $|m_{ee}|$ can be justified in the following way: since we are working in the basis where the charged leptons are diagonal, a vanishing $|m_{ee}|$ would imply a vanishing parameter $a$; in that case, the mass matrix in eq.(6) would have a zero mass eigenvalue associated to the eigenvector $(1, 0, 0)^T$, which is not compatible with the neutrino experimental data.

![Graph showing the effective mass $|m_{ee}|$ in eV allowed in the model.](image)

**FIG. 2:** Values of the effective mass $|m_{ee}|$ (in eV) allowed in our model. The two dashed horizontal lines represent the experimental sensitivities of some of the forthcoming experiments while the dashed vertical line is the upper limit from tritium $\beta$-decay experiment. For references to experiments see [22–26].

### IV. CONCLUSION

In this paper we have studied a model explaining the stability of dark and giving an interesting neutrino phenomenology based on the $A_4$ flavor symmetry. The model is an extension of the standard model and contains four Higgs doublets and five heavy right-handed neutrinos. Light neutrino masses are generated with the type-I seesaw and the resulting Majorana mass matrix has only five free parameters. We made a numerical scan of the allowed parameter space and found a correlation between the atmospheric and the reactor angles. In particular, for $\theta_{13} \lesssim 1^\circ$ maximal atmospheric angle is strongly disfavoured whereas for small reactor angle the $\theta_{23}$ is close to maximal. The model gives only normal neutrino mass hierarchy and predicts a lower bound for the neutrinoless double beta decay $|m_{ee}| \gtrsim 4 \cdot 10^{-4}$ eV.

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Appendix A: Mass hierarchy

1. Neutrino mass hierarchies for the real mass matrix

In this section we show that our model allows only the normal hierarchy in the simplified case the mass matrix in eq. (3) is real. Since the absolute neutrino mass scale is unknown, we can restrict ourselves to study the situation where the lightest neutrino mass eigenvalue is (almost) vanishing. For our discussion it is then enough to consider the determinant of such a matrix, which results:

$$\det(m_\nu) = -a^2(-bc + d)^2.$$  \hspace{1cm} (A1)

For an almost vanishing determinant, we can have two options, namely $a \sim 0$ and $d \sim bc$. In the first case, the neutrino mass matrix allows a diagonalizing matrix of the form:

$$m_\nu \simeq \begin{pmatrix} 1 & 0 & 0 \\ 0 & x & z \\ 0 & y & f \end{pmatrix},$$  \hspace{1cm} (A2)

whose first line implies a vanishing solar mixing angle. The phenomenological interesting case corresponds then to $a \neq 0$ but, even in this case, the null eigenvalue will always be associated with an eigenvector different from $(0, 1, -1)^T$, so it cannot be $m_{\nu_1}$ anyway. This excludes the inverted hierarchy as a viable neutrino mass spectrum. The second choice $d = bc$ always produces two vanishing neutrino masses; if one of them would be $m_{\nu_3}$ then the atmospheric mass difference would be vanishing; on the other hand, if $m_{\nu_1} = 0$, then one can have $\Delta m^2_{21} = 0$ but $\Delta m^2_{31} \neq 0$. Then the normal mass scheme provides the only framework where to account for such results.

2. Neutrino mass hierarchies for the vanishing reactor angle

We work here in the limit of vanishing $\theta_{13}$, which is a good approximation as given by the experimental data. In our case this limit implies $b = c$ in eq. (6). This matrix has an eigenvector $(0, -1, 1)$ with eigenvalue

$$m_{\nu_3} = (c^2 - de^{i\phi_d}).$$  \hspace{1cm} (A3)

The absolute value of this mass squared is

$$|m_{\nu_3}|^2 = c^4 - 2c^2d\cos \phi_d + d^2.$$  \hspace{1cm} (A4)

From the invariant equations (7), the determinant of the squared mass matrix is

$$\det(M^2_\nu) = a^4 (c^4 - 2c^2d\cos \phi_d + d^2)^2 = |m_{\nu_1}|^2 |m_{\nu_2}|^2 |m_{\nu_3}|^2.$$  \hspace{1cm} (A5)

From eqs. (A1) and (A3) we have the relation

$$a^4 = \frac{|m_{\nu_1}|^2 |m_{\nu_2}|^2}{|m_{\nu_3}|^2};$$  \hspace{1cm} (A6)
we also know that \(|m_{bb}| = a^2\) and therefore
\[
|m_{bb}|^2 = |m_{\nu 1}|^2 \frac{|m_{\nu 2}|^2}{|m_{\nu 3}|^2} = |m_{\nu 1}|^2 \left( c_\odot^4 + s_\odot^4 + s_\odot^4 \frac{\Delta m_{12}^2}{|m_{\nu 1}|^2} + c_\odot^2 s_\odot^2 \frac{\sqrt{\Delta m_{12}^2 + |m_{\nu 1}|^2}}{|m_{\nu 1}|} \cos \alpha \right)
\]
(A7)

where \(\alpha\) is the Majorana phase and \(c_\odot \equiv \cos \theta_{12}\) and \(s_\odot \equiv \sin \theta_{12}\). From the previous relation we get:
\[
\frac{|m_{\nu 2}|^2}{|m_{\nu 3}|^2} = \left( c_\odot^4 + s_\odot^4 + s_\odot^4 \frac{\Delta m_{12}^2}{|m_{\nu 1}|^2} + c_\odot^2 s_\odot^2 \frac{\sqrt{\Delta m_{12}^2 + |m_{\nu 1}|^2}}{|m_{\nu 1}|} \cos \alpha \right)
\]
(A8)

which, for \(|m_{\nu 1}| >> \sqrt{\Delta m_{12}^2}\), can be approximated by:
\[
\frac{|m_{\nu 2}|^2}{|m_{\nu 3}|^2} \approx \left( c_\odot^4 + s_\odot^4 + c_\odot^2 s_\odot^2 \cos \alpha \right) = 1 - c_\odot^2 s_\odot^2 (2 - \cos \alpha) < 1
\]
(A9)

That means that \(|m_{\nu 2}|/|m_{\nu 3}| < 1\) and implies a normal hierarchy. In the limit \(m_{\nu 1} \sim \Delta m_{12}^2\) we cannot have the inverse hierarchy because the minimal value for \(m_{\nu 1}\) is \(m_{\nu 1}^{min} = \sqrt{\Delta m_{13}^2} > \sqrt{\Delta m_{12}^2}\).

Notice that the relation (A8) can be expressed as an implicit function of \(m_{\nu 1}\), \(\Delta m_{12}^2\), \(\Delta m_{13}^2\), \(\theta_{12}\) and the Majorana phase \(\alpha\). Using the experimental information on the mass differences and the solar angle we can numerically evaluate the minimum allowed value for \(m_{\nu 1}\).

3. TBM limit

In this appendix we briefly discuss the TBM limit of the mass matrix in eq. (3). This can be achieved imposing the following relations among the parameters:
\[
\begin{align*}
b &= c \\
k &= a^2 + ab - 2b^2
\end{align*}
\]
(A10)

(the first relation is enough to get a \(\mu - \tau\) invariant mass matrix). Then the masses depend on two complex parameters that can be ordered according to:
\[
\begin{align*}
m_1 &= a(a - b) \\
m_2 &= a(a + 2b) \\
m_3 &= -a^2 - ab + 2b^2.
\end{align*}
\]
(A11)

It is easier to study the phenomenology redifining
\[
\begin{align*}
a &= |a| e^{i \phi_a} \\
b &= |b| e^{i \phi_b} \\
\frac{|b|}{|a|} = t & \quad (\phi_a - \phi_b) = \Delta \phi
\end{align*}
\]

so that the solar and atmospheric mass differences are:
\[
\begin{align*}
\Delta m_{21}^2 &= 3|a|^4 t (2 \cos \Delta \phi + t) \\
\Delta m_{31}^2 &= 4|a|^4 t (\cos \Delta \phi + t - 2t \cos^2 \Delta \phi - t^2 \cos \Delta \phi + t^3).
\end{align*}
\]
(A12)
The model is only compatible with a normal hierarchy spectrum because the simultaneous requirements
\[ \Delta m^2_{21} > 0 \quad \text{and} \quad \Delta m^2_{31} < 0 \]
gives
\[ -\frac{t}{2} < \cos \Delta \phi < 1 \]  
(A13)
\[ \cos \Delta \phi < -t \]
which are obviously incompatible. Moreover, it easy to check that the conditions \(|m_1| > 0\) and \(|m_3| \lesssim 0.5\) eV imply:
\[ 0.07 \lesssim t \lesssim 5 \]  
(A14)
For \(|m_{ee}|\) we found a lower bound \(|m_{ee}| > 6 \cdot 10^{-4} \text{ eV}\); other neutrino mass matrix models with two complex parameters predicts different lower limits in the TBM limit, for instance \(|m_{ee}| > 7 \cdot 10^{-3} \text{ eV}\) in \cite{27} and \(|m_{ee}| = 0\) in \cite{28}.

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