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Neutrino phenomenology and stable dark matter with A_4

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ABSTRACT

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 θ_{13}. Only normal neutrino mass spectrum is possible and the effective mass entering the neutrinoless double beta decay rate is constrained to be $|m_{ee}| > 4 \times 10^{-4}$ eV.

1. 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) [1] observed, within one standard deviation, in the recent neutrino oscillation data [2,3], 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 [4] and S_4 [5] 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 [6], 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 [7] 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 [8], or from the breaking of Abelian $U(1)$ group [9–11] as, for example, in grand unified frameworks or in supersymmetry [12].

Recently a model where the stability of the DM arises from the spontaneous breakdown of a non-Abelian discrete flavor symmetry was proposed [13]. 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 [14–16]). In [13] 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 models where the stability of the dark matter is justified by non-Abelian discrete flavor symmetries. In this Letter 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 Letter is organized as follows: in Section 2 we give the field content of our model and derive the neutrino mass matrix, whose phenomenological implications are discussed in Section 3; in Section 4 we draw our conclusions.

2. The model

In Table 1 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 $1''$. The operators needed to generate neutrinos masses are the following:

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Table 1
Summary of relevant model quantum numbers.

	L_e	L_μ	L_τ	l_e^c	l_μ^c	l_τ^c	N_T	N_4	N_5	H	η
$SU(2)$	2	2	2	1	1	1	1	1	1	2	2
A_4	1	$1'$	$1''$	1	$1''$	$1'$	3	$1'$	$1''$	1	3

$$\mathcal{L} = y_1^\nu L_e (N_T \eta)_1 + y_2^\nu L_\mu (N_T \eta)_{1''} + y_3^\nu L_\tau (N_T \eta)_{1'} + y_4^\nu L_\tau N_4 H + y_5^\nu L_\mu N_5 H + M_1 N_T N_T + M_2 N_4 N_5 + \text{h.c.}$$

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 η do not take vev and we have:

$$\langle H^0 \rangle = v_h \neq 0, \qquad \langle \eta_1^0 \rangle = v_\eta \neq 0, \qquad \langle \eta_{2,3}^0 \rangle = 0. \tag{1}$$

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

$$m_D = \begin{pmatrix} y_1^\nu v_\eta & 0 & 0 & 0 & 0 \\ y_2^\nu v_\eta & 0 & 0 & 0 & y_5^\nu v_h \\ y_3^\nu v_\eta & 0 & 0 & y_4^\nu v_h & 0 \end{pmatrix},$$

$$m_M = \begin{pmatrix} M_1 & 0 & 0 & 0 & 0 \\ 0 & M_1 & 0 & 0 & 0 \\ 0 & 0 & M_1 & 0 & 0 \\ 0 & 0 & 0 & 0 & M_2 \\ 0 & 0 & 0 & M_2 & 0 \end{pmatrix}. \tag{2}$$

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

$$m_\nu = \begin{pmatrix} a^2 & ab & ac \\ ab & b^2 & bc+k \\ ac & bc+k & c^2 \end{pmatrix}, \tag{3}$$

where we defined:

$$a = \frac{y_1^\nu v_\eta}{\sqrt{M_1}}, \qquad b = \frac{y_2^\nu v_\eta}{\sqrt{M_1}},$$
$$c = \frac{y_3^\nu v_\eta}{\sqrt{M_1}}, \qquad k = \frac{y_4^\nu y_5^\nu v_h^2}{M_2}. \tag{4}$$

As discussed in more details in [13], the group A_4 is broken by the vev $\langle \eta_1^0 \rangle \sim (1,0,0)$ down to the subgroup Z_2, generated by the diagonal A_4 generator $S = \mathrm{Diag}\{1,-1,-1\}$; the Z_2 symmetry acts on the triples fields in the following way:

$$Z_2: \quad N_2 \to -N_2,\ h_2 \to -h_2,\ A_2 \to -A_2,$$
$$N_3 \to -N_3,\ h_3 \to -h_3,\ A_3 \to -A_3, \tag{5}$$

where h_i and A_i are respectively the CP-odd and CP-even components of the Higgs doublet η_i for $i = 2,3$ and $N_{2,3}$ are the components 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 [17,18].

3. Numerical analysis of the full mass matrix

The most general complex and symmetric matrix has twelve independent real parameters which, after readsorbing 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 $bc + k = de^{i\phi_d}$. In terms of them, the matrix in Eq. (3) is rewritten in the following form:

$$m_\nu = \begin{pmatrix} a^2 & ab & ac \\ ab & b^2 & de^{i\phi_d} \\ ac & de^{i\phi_d} & c^2 \end{pmatrix}. \tag{6}$$

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 $M_\nu^2 = m_\nu m_\nu^\dagger$:

$$\begin{aligned}
\mathrm{Tr}(M_\nu^2) &= T = a^4 + 2a^2(b^2+c^2) + b^4 + c^4 + 2d^2 \\
&= m_{\nu_1}^2 + m_{\nu_2}^2 + m_{\nu_3}^2, \\
\det(M_\nu^2) &= a^4 (b^2c^2 - 2bcd\cos\phi_d + d^2)^2 \\
&= m_{\nu_1}^2 m_{\nu_2}^2 m_{\nu_3}^2, \\
\frac{1}{2}\big[T^2 &- \mathrm{Tr}(M_\nu^2 M_\nu^2)\big] \\
&= (b^2c^2 - 2bcd\cos\phi_d + d^2) \\
&\quad \times \big[2a^4 + 2a^2(b^2+c^2) + b^2c^2 + 2bcd\cos\phi_d + d^2\big] \\
&= m_{\nu_1}^2(m_{\nu_2}^2 + m_{\nu_3}^2) + m_{\nu_2}^2 m_{\nu_3}^2,
\end{aligned} \tag{7}$$

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_{ν_1}, the two square mass differences Δm^2_{sol} and Δm^2_{atm} and the other two parameters of the neutrino mass matrix, i.e. d and ϕ_d; to study the correlations among the neutrino observables predicted by the model, we perform a parameter space scan extracting randomly the previous parameters in the following intervals:

$$d \in [-1,1], \qquad \phi_d \in [-\pi,\pi), \qquad m_{\nu_1} \in \big[10^{-5}, 1\big]\ \mathrm{eV},$$
$$7.1 \times 10^{-5}\ \mathrm{eV}^2 < \Delta m^2_{sol} < 8.3 \times 10^{-5}\ \mathrm{eV}^2,$$
$$0.027 < r = \frac{\Delta m^2_{sol}}{|\Delta m^2_{atm}|} < 0.040.$$

Then, we require that the obtained mixing angles are within their current 3σ 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 analitically discussed this point in the appendices: Appendix A.1 for the real case, Appendix A.2 for the μ–τ invariant case and Appendix A.3 for the exact TBM case). Among all possible correlations, we found that the (θ_{13}–θ_{23}) one is quite interesting, as shown in Fig. 1. As

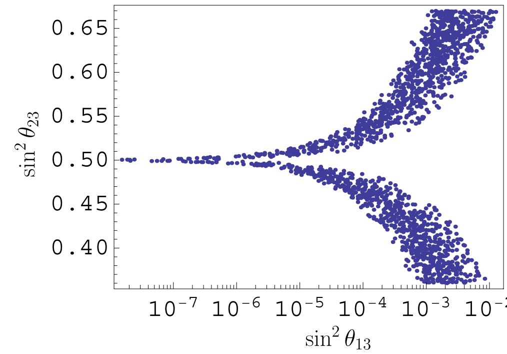

Fig. 1. Correlation among the θ_{13} and θ_{23} angles as predicted by the model.

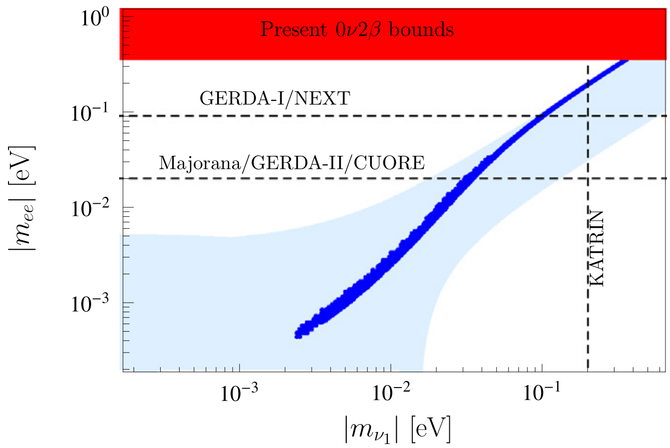

Fig. 2. Values of the effective mass $|m_{ee}|$ (in eV) allowed in our model. The two dashed horizontal lines represent the experimental sensitivities of some of the forthcoming experiments while the dashed vertical line is the upper limit from tritium β-decay experiment. For references to experiments see [20–24].

we can see, for very small θ_{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 θ_{13} two different branches above and below $\theta_{23} = \pi/4$ develop and for $\theta_{13} \gtrsim 1°$ maximal 2–3 mixing is strongly excluded. It is interesting to observe that this result can be easily verified at the forthcoming neutrino experiments [19] even with a reduced sensitivity to θ_{13} since the largest deviation from $\theta_{23} = \pi/4$ is obtained for θ_{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 $\mathcal{O}(10^{-2})$, which ensures that no fine-tuning is required among the matrix elements of Eq. (3), and with the phase ϕ 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_{ν_1}. The result of our simulation can be found in Fig. 2, where we can see that a lower bound $(|m_{\nu_1}|, |m_{ee}|) \gtrsim (2 \times 10^{-3}, 4 \times 10^{-4})$ 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.

4. Conclusion

In this Letter 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 $\theta_{13} \gtrsim 1°$ maximal atmospheric angle is strongly disfavored whereas for small reactor angle the θ_{23} is close to maximal. The model gives only normal neutrino mass hierarchy and predicts a lower bound for the neutrinoless double beta decay $|m_{ee}| \gtrsim 4 \times 10^{-4}$ eV.

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Appendix A. Mass hierarchy

A.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:

$$\det(m_\nu) = -a^2(-bc + d)^2. \tag{A.1}$$

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:

$$m_\nu \simeq \begin{pmatrix} 1 & 0 & 0 \\ 0 & x & z \\ 0 & y & f \end{pmatrix}, \tag{A.2}$$

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_{ν_3} 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_{ν_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.

A.2. Neutrino mass hierarchies for the vanishing reactor angle

We work here in the limit of vanishing θ_{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

$$m_{\nu 3} = \left(c^2 - d e^{i\phi_d}\right). \tag{A.3}$$

The absolute value of this mass squared is

$$|m_{\nu 3}|^2 = c^4 - 2c^2 d \cos\phi_d + d^2. \tag{A.4}$$

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

$$\begin{aligned}\det\left(M_\nu^2\right) &= a^4\left(c^4 - 2c^2 d\cos\phi_d + d^2\right)^2 \\ &= |m_{\nu_1}|^2 |m_{\nu_2}|^2 |m_{\nu_3}|^2.\end{aligned} \tag{A.5}$$

From Eqs. (A.4) and (A.5) we have the relation

$$a^4 = \frac{|m_{\nu_1}|^2 |m_{\nu_2}|^2}{|m_{\nu_3}|^2}; \tag{A.6}$$

we also know that $|m_{bb}| = a^2$ and therefore

$$\begin{aligned} |m_{bb}|^2 &= |m_{\nu 1}|^2 \frac{|m_{\nu_2}|^2}{|m_{\nu_3}|^2} \\ &= |m_{\nu 1}|^2 \Bigg(c_\odot^4 + s_\odot^4 + s_\odot^4 \frac{\Delta m_{12}^2}{|m_{\nu 1}|^2} \\ &\quad + c_\odot^2 s_\odot^2 \frac{\sqrt{\Delta m_{12}^2 + |m_{\nu 1}|^2}}{|m_{\nu 1}|} \cos\alpha \Bigg) \end{aligned} \tag{A.7}$$

where α is the Majorana phase and $c_\odot \equiv \cos\theta_{12}$ and $s_\odot \equiv \sin\theta_{12}$. From the previous relation we get:

$$\begin{aligned} \frac{|m_{\nu_2}|^2}{|m_{\nu_3}|^2} &= \Bigg(c_\odot^4 + s_\odot^4 + s_\odot^4 \frac{\Delta m_{12}^2}{|m_{\nu 1}|^2} \\ &\quad + c_\odot^2 s_\odot^2 \frac{\sqrt{\Delta m_{12}^2 + |m_{\nu 1}|^2}}{|m_{\nu 1}|} \cos\alpha \Bigg) \end{aligned} \tag{A.8}$$

which, for $|m_{\nu 1}| \gg \sqrt{\Delta m_{12}^2}$, can be approximated by:

$$\begin{aligned} \frac{|m_{\nu_2}|^2}{|m_{\nu_3}|^2} &\sim \left(c_\odot^4 + s_\odot^4 + c_\odot^2 s_\odot^2 \cos\alpha \right) \\ &= 1 - c_\odot^2 s_\odot^2 (2 - \cos\alpha) < 1. \end{aligned} \tag{A.9}$$

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 (A.8) can be expressed as an implicit function of $m_{\nu 1}$, Δm_{12}^2, Δm_{13}^2, θ_{12} and the Majorana phase α. Using the experimental information on the mass differences and the solar angle we can numerically evaluate the minimum allowed value for $m_{\nu 1}$.

A.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:

$$\begin{aligned} &b = c, \\ &k = a^2 + ab - 2b^2 \end{aligned} \tag{A.10}$$

(the first relation is enough to get a μ–τ invariant mass matrix). Then the masses depend on two complex parameters that can be ordered according to:

$$\begin{aligned} &m_1 = a(a-b), \\ &m_2 = a(a+2b), \\ &m_3 = -a^2 - ab + 2b^2. \end{aligned} \tag{A.11}$$

It is easier to study the phenomenology redefining

$$\begin{aligned} &a = |a| e^{i\phi_a}, \\ &b = |b| e^{i\phi_b}, \\ &\frac{|b|}{|a|} = t, \qquad (\phi_a - \phi_b) = \Delta\phi \end{aligned}$$

so that the solar and atmospheric mass differences are:

$$\begin{aligned} &\Delta m_{21}^2 = 3|a|^4 t (2\cos\Delta\phi + t), \\ &\Delta m_{31}^2 = 4|a|^4 t \big(\cos\Delta\phi + t - 2t\cos^2\Delta\phi - t^2\cos\Delta\phi + t^3\big). \end{aligned} \tag{A.12}$$

The model is only compatible with a normal hierarchy spectrum because the simultaneous requirements $\Delta m_{21}^2 > 0$ and $\Delta m_{31}^2 < 0$ give

$$\begin{aligned} &-\frac{t}{2} < \cos\Delta\phi < 1, \\ &\cos\Delta\phi < -t, \end{aligned} \tag{A.13}$$

which are obviously incompatible. Moreover, it easy to check that the conditions $|m_1| > 0$ and $|m_3| \lesssim 0.5$ eV imply:

$$0.07 \lesssim t \lesssim 5. \tag{A.14}$$

For $|m_{ee}|$ we found a lower bound $|m_{ee}| > 6 \times 10^{-4}$ eV; other neutrino mass matrix models with two complex parameters predict different lower limits in the TBM limit, for instance $|m_{ee}| > 7 \times 10^{-3}$ eV in [25] and $|m_{ee}| = 0$ in [26].

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