Neutron Majorana mass from exotic instantons in a Pati-Salam model

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ABSTRACT: We show how exotic stringy instantons can generate an effective interaction between color diquark sextets in a Pati-Salam model, inducing a Majorana mass term for the neutron. In particular, we discuss a simple quiver theory for a Pati-Salam like model, as an example in which the calculations of exotic instanton effects are simple and controllable. We discuss some different possibilities in order to generate $n - \bar{n}$ oscillations testable in the next generation of experiments, Majorana mass matrices for neutrini and a Post-Sphaleron Baryogenesis scenario. Connections with Dark Matter issues and the Higgs mass Hierarchy problem are discussed, in view of implications for LHC and rare processes physics. The model may be viewed as a completion of a Left-Right symmetric extension of the Standard Model, alternative to a GUT-inspired scenario. Combined measures in Neutron-Antineutron physics, FCNC, LHC, Dark Matter could rule out the proposed model or uncover aspects of physics at the Planck scale!

Contents

1 Introduction

How can Matter be generated in our Universe? And how are neutrino masses generated? Has the neutron a Majorana mass?

In principle, these three questions could appear unrelated. However, in Left-Right symmetric models with $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}_c(3) \times \mathrm{U}(1)_{B-L}$ gauge group, one can find intriguing and elegant connections between these three issues. A Left-Right model is naturally embedded in a Pati-Salam (P-S) model with $G_{224} = \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(4)_c$ [1], that in turn can be embedded in an SO(10) GUT.

As originally suggested in [2], new Higgses Δ_R in the $(\mathbf{1}, \mathbf{3}, \mathbf{10})$ (and Δ_L in the $(\mathbf{3}, \mathbf{1}, \mathbf{10}^*)$) of G_{224} can be introduced in PS models in order to spontaneously break Left-Right symmetry, through $\langle \Delta_R \rangle = v_R \neq 0$ and $\langle \Delta_L \rangle = 0$. This mechanism also produces Majorana masses for Right-Handed neutrinos, that can trigger a seesaw mechanism as suggested in [3], spontaneously breaking $\mathrm{U}(1)_{B-L}$ at the same time.[1] The new Higgs $\Delta_R(1,3,10) \equiv \Delta^c(1,3,10^*)$ decomposes with respect to $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(3)_C \times \mathrm{U}(1)_{B-L}$ as

$$\Delta_R(1,3,10) = \left\{ (1,3,1)_{-2} + (1,3,3)_{-2/3}, +(1,3,6)_{2/3} \right\}_R \tag{1.1}$$

with $\Delta^c_{l^c l^c}(1,3,1)_{-2}$ generating Right-Handed neutrini masses via $\langle \Delta^c_{\nu^c \nu^c} \rangle \nu^c \nu^c$.

In GUT SO(10), the $(\mathbf{1}, \mathbf{3}, \mathbf{10})$ of G_{224} and its conjugate are contained in the $\mathbf{126}$ representation.[2] But $\Delta^c(1,3,10)$ also contains *color sextet diquark fields* $\Delta^c_{q^c q^c}(1,3,6)_{2/3}$,

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[1]See also [4, 5] for more about seesaw mechanisms.

[2]The complete decomposition reads $\mathbf{126} \rightarrow (\mathbf{1}, \mathbf{3}, \mathbf{10}) + (\mathbf{3}, \mathbf{1}, \mathbf{10}^*) + (\mathbf{2}, \mathbf{2}, \mathbf{15}) + (\mathbf{1}, \mathbf{1}, \mathbf{6})$.

leading to possible new effects. In particular, these sextets can induce Baryon number violating processes beyond the Standard Model (BSM). Color sextets can also play an important role in some post-sphaleron baryogenesis mechanism [6–8]. In susy extensions, a quartic superpotential term

$$\mathcal{W}_4 = \Delta^c \Delta^c \Delta^c \Delta^c / \mathcal{M}_0$$

can appear that, among other terms, produces a term coupling three color sextets $\Delta^c_{q^c q^c}$ and one color singlet $\Delta^c_{\nu^c \nu^c}$, as $\Delta^c_{u^c u^c} \Delta^c_{d^c d^c} \Delta^c_{d^c d^c} \Delta_{\nu^c \nu^c}$. When the color singlet $\Delta^c_{\nu^c \nu^c}$ takes an expectation value, $U(1)_{B-L}$ is spontaneously broken and Right-handed neutrini get a mass [3]. Moreover a Majorana mass for the neutron is generated through the processes shown in figure 1-(a)-(b). This can be directly tested in Neutron-Antineutron transition experiments, as firstly proposed in [2]. As shown in [9], constraints from post-sphaleron baryogenesis and neutrino oscillations imply a precise prediction about neutron-antineutron transitions: an oscillation time $\tau_{n-\bar{n}} \approx 10^{10} s$ accessible to the next generation of experiments![3]

In principle, color sextet scalars could be as light as $1\,\mathrm{TeV}$ and they could be directly searched at the LHC, as proposed in [11]: dijet data put constraints on the couplings between colored scalars and quarks. In [12], bounds are shown in comparison with LHC data. On the other hand, processes mediated by FCNC could impose stronger constraints on the sextets with respect to LHC direct searches (see [13] for comparison with experimental limits). For example, the Δ^c_{dd} field couples to two down-type quarks dd, ss, bb and mediates $B^0_{d,s} - \bar{B}^0_{s,s}$, $K^0 - \bar{K}^0$ oscillations as well as B mesons decays. On the other hand Δ^c_{uu} mediates $D^0 - \bar{D}^0$ oscillations and D-decays like $D \to K\pi, \pi\pi$. These analyses show that for coupling strengths of order 10^{-2}, the mass of the color sextets has to exceed the TeV-scale.

In this paper, we propose a (SUSY) PS model that is alternative to the SO(10) GUT inspired model mentioned above. We consider an (un)oriented open string model with intersecting D-branes, producing a susy PS like model. Models of this kind have been previously considered e.g. in [15], where an analysis of the mass spectrum and low-energy phenomenology has been carried out. In oriented string theory, a simple way to generate a U(N) gauge theory is to consider a stack of N D-branes, parallel to each other. In this way the open string excitations stretching between the N D-branes reproduce at low energy the fields in the adjoint of the U(N) gauge symmetry. In type IIA, compactified on a six-dimensional (CY) manifold, one can consider stacks of intersecting D6-branes, filling the 4D ordinary Minkowski spacetime, and wrapping internal 3-cycles. From strings connecting different stacks of branes, we can construct chiral fermions, localised at the four-dimensional intersections of two stacks of D6-branes a and b, in the bi-fundamental representation of U(N_a) × U(N_b) [16]. The net (positive-negative) number of intersections of two branes a and b is a topological invariant, representing the number of massless fermions. In the case in which D-branes are space-time filling, Ω-planes have to be introduced in order to cancel tadpoles and irreducible anomalies [17–22]. An Ω-plane implements a combination of world-sheet parity and a (non) geometric mirror-like involution in the target space. As a consequence Left- and Right-moving modes of the closed strings are identified; closed and open strings become un-oriented. More choices for the gauge groups and their representations are allowed [17–19]. In this way, one can produce stacks supporting U(N), SO(N) or

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[3]Another string-inspired mechanism for $n - \bar{n}$ was proposed in [10], in heterotic superstrings theories.
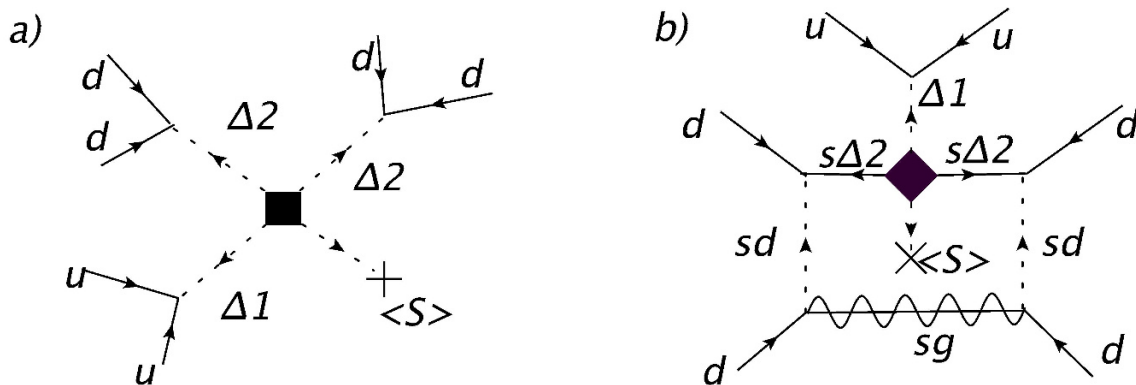
Figure 1. a) We show the main diagram for neutron-antineutron oscillation in a non-susy $SU(4)_c \times SU(2)_L \times SU(2)_R$ model [14]. The transition is induced by color sextets $\Delta_{u^c u^c}^c$ and $\Delta_{d^c d^c}^c$. b) We show the main diagram for neutron-antineutron oscillation in a supersymmetric $SU(4)_c \times SU(2)_L \times SU(2)_R$ model [14]. The transition is induced by color sextets $\Delta_{u^c u^c}^c$ and $\Delta_{d^c d^c}^c$ arising from the decomposition of $\Delta^c(1, 3, 10)$. The latter participate in a non-perturbative quartic superpotential term. The diagram involves also gaugini $\tilde{g}$ (gluini, zino or bino), squarks $\tilde{d}^c$, and susy partners of the color sextets $\tilde{\Delta}_{d^c d^c}^c$. In figures, we use the notation $\Delta_{1,2} = \Delta_{u^c u^c, d^c d^c}$ and s-particle for supersymmetric partners.

$Sp(2N)$ gauge groups. This is interesting in order to construct realistic gauge groups, with chiral matter in a globally consistent model [23, 24]. The closed strings propagate in the entire ten dimensional space-time: some mediate gravitational interactions, some behave as axions or scalar moduli fields.

In principle, one can construct a PS like gauge group $U(4) \times Sp_R(2) \times Sp_L(2)$ or $U(4) \times U_R(2) \times U_L(2)$ in terms of intersecting D-brane stacks and Ω-planes. In [15] the case $U(4) \times U_R(2) \times U_L(2)$ was analysed in some detail. In the present paper, we focus on the $U(4) \times Sp_R(2) \times Sp_L(2)$ case with an Ω^+-plane that requires a stack of four D-branes and its mirror image under Ω, producing $U(4)$, and two stacks of two D-branes each, identified with their own images under Ω, producing $Sp_L(2)$ and $Sp_R(2)$.

This model has extra anomalous $U(1)$'s that could seem dangerous from a gauge theory point of view. On the other hand, in string theory, Generalized Chern-Simon (GCS) terms appear that cancel anomalies [25, 26], in combination with a generalised Green-Schwarz mechanism [27, 28]. The extra Z' gauge bosons can get a mass through a Stückelberg mechanism [29–32]. We will return onto phenomenological implications of this in the next section. There is however a real problem in this scenario. It is not possible to represent $(1, 3, 10)$ in terms of open strings. Perturbative open strings have two ends and can at most carry fundamental charges with respect to two classical gauge groups. On the other hand, the triplet is the adjoint of $Sp(2)$, i.e. the symmetric product of two doublets, and the decaplet of $SU(4)$ is the symmetric product of two tetraplets. States in the $(1, 3, 10)$ (or its conjugate) would correspond to multi-pronged strings with two ends on the $U(4)$ stack and two ends on the $Sp(2)$ stack[4] that do not admit a perturbative description.

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[4] **3** is also the vector of SO(3) but this would prevent the existence of doublets, which are spinors.

On the other hand, we will show that a spontaneously breaking pattern to the SM, giving masses to the neutrini, can be recovered in this model. In fact, we will see that $\phi_{RR}(1,3,1)$ and $\phi_{LL}(1,1,3)$ appear as excitations of open strings with both ends attached to $\text{Sp}(2)_R$ or $\text{Sp}(2)_L$, while $\Delta(1,1,10)$ and its conjugate $\Delta^c(1,1,10^*)$ appear from open strings joining $\text{U}(4)$ and $\text{U}'(4)$ identified which one other under Ω. As a consequence the breaking $\text{U}(4) \times \text{Sp}_R(2) \times \text{Sp}_L(2) \to \text{SU}(3) \times \text{SU}_L(2) \times \text{U}(1)$ is not realized through $(1,3,10)$, but through $\phi_{RR}(1,3,1)$ and $\Delta(1,1,10)$:[5]

i) Left-Right symmetry breaking through the expectation values $\langle \phi_{RR} \rangle = v_R$ and $\langle \phi_{LL} \rangle = 0$;

ii) $\text{U}(1)_{B-L}$ symmetry breaking though the expectation of value $\langle S \rangle = v_{B-L}$, with S the color singlet contained in $(1,1,10)$. Alternatively, $\text{U}(1)_{B-L}$ can be broken dynamically by exotic instantons or spontaneously by the compactification.

Similarly to the case of $\Delta(1,3,10)$, color sextets are contained in $\Delta(1,1,10)$.

Our main suggestion is that the super-potential

$$\mathcal{W}_{\text{eff}} = S^c \Delta_6^c \Delta_6^c \Delta_6^c / \mathcal{M}_0 \tag{1.2}$$

be generated by non-perturbative quantum gravity effects peculiar to string theory, called "exotic instantons". These are associated to Euclidean branes ($E2$-branes in our case), wrapping internal 3-cycles, that could directly produce such interactions, in a calculable and controllable way in models like type IIA (un)oriented strings. We would like to stress that this class of instantons exists in string theory only, not in gauge theories. The resulting superpotential term is suppressed by the scale $\mathcal{M}_0 = M_S e^{+S_{E2}}$, where M_S is the string scale and $e^{+S_{E2}}$ depends on the 'size' of the 3-cycles wrapped by the relevant $E2$-brane. We would like to remark that the suppression scale is higher (in principle also much higher) than the string scale. This is a peculiarity of the non-renormalizable nature of such a non-perturbative term in the string effective action. As a consequence, the hierarchy depends on the particular model: $e^{+S_{E2}}$ can be approximately 1 for a 'small' 3-cycles or $e^{+S_{E2}} \gg 1$ for a 'large' 3-cycles. So, depending on the string scale M_s, assumed to be larger than some TeV's at least, and the size of the 3-cycle, it is possible to generate such an operator near the LHC scale or at a much higher scale. This leads to two very different branches for phenomenology. In particular, for $\mathcal{M}_0 \simeq 10^{13}\,\text{GeV}$, color sextets appear near the TeV scale, with potential implications in meson physics and at LHC, as mentioned above. On the other hand, for $\mathcal{M}_0 \simeq M_S \simeq 10\,\text{TeV}$, a post-sphaleron scenario is possible and testable at the next generation of experiments on neutron-antineutron oscillations, with heavy color sextets, at a scale $m_6 \gg \text{TeV}$ that can be generated by closed-string fluxes, as shown in [33] for quiver theories and reviewed below. In this case, there is no possibility to produce the sextets at LHC, and FCNC's in the meson sector are strongly suppressed. On the other hand, extra Z' at the TeV scale naturally appears in this scenario. Another relevant and peculiar possibility is a Left-Right breaking scale at TeV, compatible with

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[5]In [15] breaking triggered by Higgses in the $(1,2,4^*)$ was considered together with mass terms generated by exotic instantons.

neutron-antineutron physics and Post-Sphaleron scenario. This is not possible in a SO(10) scenario, as remarked in [9]. Our string-inspired scenario also naturally provides several candidates of WIMP dark matter as we will see.

We would like to mention that such an operator as (1.2) can emerge from stringy dynamics also in other kinds of models like F-theory [34–40], $E_8 \times E_8$ and SO(32) heterotic strings [41–49], generating an SO(10) GUT. For example in heterotic string theories *world-sheet instantons*, suppressed by $e^{-R^2/\alpha'}$ and thus perturbative in the string coupling g_s, can induce non-vanishing couplings of the desired kind from such amplitudes as $\langle V_{\mathbf{126}} V_{\mathbf{126}} V_{\mathbf{126}} V_{\mathbf{126}} \rangle$, for vertex operators $V_{\mathbf{126}}$ that can appear in twisted sectors. Unfortunately, in the F-theory case the calculations are more involved [50, 51].

The paper is organized as follows: in section 2, we briefly review what are stringy instantons and quivers. In section 3 we propose a simple and consistent quiver for a Pati-Salam model generating a Majorana mass for the neutron through exotic instantons. In section 4 we discuss some phenomenology resulting from this model. In section 5 we present our conclusions and final remarks.

2   Exotic instantons and quivers

In this section, we briefly review D-brane instantons and unoriented quiver theories.

2.1   Instantons

In 4-dimensional gauge theory, instantons are point-like configurations, that extremize the Euclidean action for a given topological charge. In string theory, instantons admit a simple geometric interpretation: they are special Euclidean branes wrapping some (internal) cycle. In theories with (unoriented) open strings, these are Euclidean D-branes (E-branes) that can intersect the 'physical' D-branes.[6] In (un-)oriented type IIA, gauge instantons can be classified as Euclidean D2 (E2) branes wrapping the same 3-cycle as a stack of "physical" D6-branes. In (un-)oriented IIB, instantons are E(-1) or E3 wrapping the same holomorphic divisor as a stack of "physical" D7-branes. In type I, instantons are E5 branes in the internal space, with the same magnetization as the D9, wrapping the entire CY_3.

2.2   Quivers

The effective low energy description of the dynamics of D-branes at Calabi-Yau singularities is captured by a quiver field theory. Usually, the (supersymmetric) quiver conventions are the following: the standard D-brane stacks are 'circle' nodes, the super-fields in the bifundamental representations of two D-brane stacks are oriented lines connecting the nodes, usually termed arrows whence the name 'quiver', 'triangle' nodes are Euclidean D-branes (instantons), grassmanian moduli or modulini are dashed lines connecting triangles and circles. Square nodes represent flavour branes, i.e. branes wrapping non-compact cycles so much so that $g_{YM,Dp}^2 = g_s(\alpha')^{p+1/2}/V_{p+1} \to 0$. These simple rules allow one to subsume the system of D-branes and open strings with a simple diagram. In this notation, perturbative

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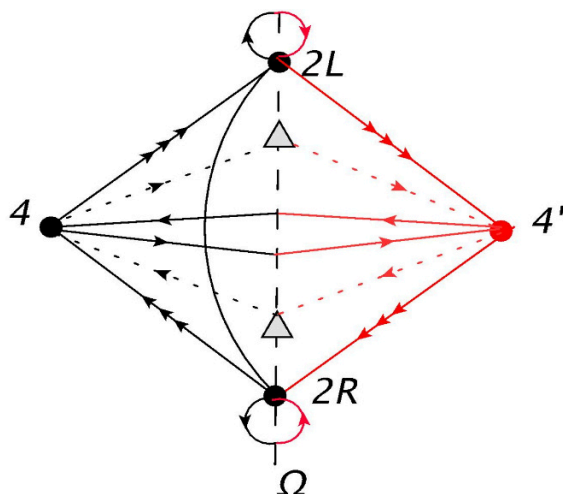[6]For a pedagogical review see e.g. [67–69].

Figure 2. Unoriented quiver for a Pati-Salam-like model $U(4) \times Sp(2)_L \times Sp(2)_R$ The circles labeled by $4, 4', 2L, 2R$ represent $U(4), U'(4), Sp(2)_L, Sp(2)_R$ stacks respectively. An Ω^+-plane identifies the $U(4)$ stack with its mirror image, $Sp(2)_{L,R}$ are stacks of two D6-branes laying exactly on top of the Ω-plane. The symmetric representations $\Delta(1, 1, 10)$ and $\Delta^c(1, 1, 10^*)$ appear in between the two stacks 4 and 4'. The triangles represent two possible $E2$-brane $O(1)$ instantons, laying on top of the Ω-plane.

interaction terms involving the matter super-fields correspond to closed oriented polygons, starting with triangles. On the other hand, interactions between standard super-fields and modulini also correspond to closed oriented polygons involving solid and dashed lines.[7]

3 An unoriented quiver for a Pati-Salam model

In this section, we construct a simple quiver for a Pati-Salam model inducing a Neutron Majorana mass for the neutron. We propose a simple quiver in figure 2, leading to a susy $\mathcal{N} = 1$ Pati-Salam like model, with all the necessary fields and Yukawa's for a spontaneous symmetry breaking pattern to the Standard Model and for the generation of a Neutron Majorana mass from Exotic Instantons.

The gauge group is $U(4) \times Sp(2) \times Sp(2)$: $U(4)$ is generated by a stacks of 4 D-branes and their images $U(4)'$ under an Ω^+-plane; $Sp(2)_{L,R}$ are generated by two stacks of two D-branes each on top of the Ω^+-plane. We also consider Exotic $O(1)$ Instantons corresponding to $E2$-branes on top of the Ω^+-plane.

In particular, the three generations of Left and Right fundamental representations $F_{L,R}$, containing quarks and leptons, are reproduced as excitations of open strings attached to the $U(4)$-stack and the Left or Right $Sp(2)_{L,R}$ stacks, respectively. We also get $\Delta = (1, 1, 10)$ and its conjugate from open strings attached to the $U(4)$-stack and its mirror image $U(4)'$-stack. $\phi_{RR} = (3, 1, 1)$ and $\phi_{LL} = (1, 3, 1)$ correspond to strings with both end-points attached on the $Sp(2)_{R,L}$ respectively. Higgs fields $h_{LR} = (2, 2, 1)$ are between

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[7]One should keep in mind that exotic instantons are brane wrapping empty nodes of the quiver. Their interactions are thus coded in the quiver or dimer.
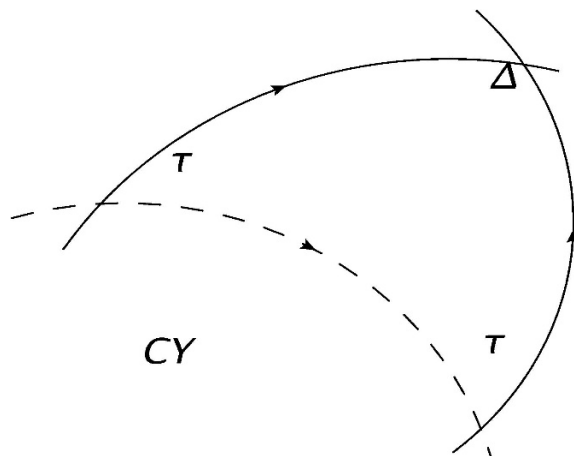
Figure 3. Amplitude in IIA (un)oriented string theory. Δ_{ab} sextets are excitations of strings attached between two intersecting $D6$-branes, represented in figure as black lines. The fermionic moduli (or modulini) result from strings localised at the intersection of one $D6$-brane and an exotic instanton Euclidean $D2$-brane (or $E2$ brane). They are represented in figure as dashed lines. This system is embedded in some Calabi-Yau compactification CY_3. In particular $D6$-branes are wrapping 3-cycles on CY_3 and $E2$ are wrapping a different 3-cycle.

$Sp(2)_L$ and $Sp(2)_R$. It is amusing to observe that a two-generation model of the same kind would result in unoriented Type IIB from 'fractional' D3-branes at a $\mathbf{C}^3/Z_4$ orbifold singularity.

The perturbative super-potential that we obtain from the quiver reads

$$\mathcal{W}_{\text{eff}} \sim y_1 h_{LR} F_L F_R + \frac{1}{M_1} F_L \phi_{LL} F_L \Delta + \frac{1}{M_2} F_R \phi_{RR} F_R \Delta^c + \frac{1}{M_3} h_{LR} \phi_{RR} h_{RL} \phi_{LL} + \mu h_{LR} h_{RL}$$

$$+ m_\Delta \Delta \Delta^c + \frac{1}{4M_4}(\Delta \Delta^c)^2 + \frac{1}{2} m_L \phi_{LL}^2 + \frac{1}{2} m_R \phi_{RR}^2 + \frac{1}{3!} a_L \phi_{LL}^3 + \frac{1}{3!} a_R \phi_{RR}^3 \qquad (3.1)$$

where $\Delta \equiv \Delta_{qq}$ and the mass scales $M_{1,2,3,4}$ depend on the particular global completion of the model: they could be near M_S or at lower scales. In a T-dual Type IIB context, the mass terms m_Δ and $m_{L,R}$ can be generated by RR-RR or NS-NS three-forms fluxes in the bulk:

$$m_{\Delta_{qq}} \sim \Gamma^{ijk} \langle \tau H_{ijk}^{(qq)} + i F_{ijk}^{(qq)} \rangle$$

$$m_{L,R} \sim \Gamma^{ijk} \langle \tau H_{ijk}^{(L,R)} + i F_{ijk}^{(L,R)} \rangle$$

with H RR-RR and F NS-NS three-forms and in general $H^{dd}, F^{dd} \neq H^{uu}, F^{uu}$ depend on the choice of fluxes through the relevant cycles wrapped by the D-branes.[8]

On the other hand, the non-perturbative superpotential term:

$$\mathcal{W}_{E2} = \frac{1}{\mathcal{M}_0} \epsilon^{ijkl} \epsilon^{i'j'k'l'} \Delta_{ii'}^c \Delta_{jj'}^c \Delta_{kk'}^c \Delta_{ll'}^c \qquad (3.2)$$

can be generated by an $E2$-brane instanton assuming the relevant $E2$ intersect twice the U(4) stack of $D6$-branes, necessary to produce a four-Δ^c (as well as a four-Δ) interaction.

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[8]Mass deformed quivers and dimers have been recently investigated in [33].

The fermionic modulini τ_α^i, with $i = 1, \ldots, 4$ and $\alpha, \beta = 1, 2$ interact with the super-fields Δ's via

$$\mathcal{L}_{E2-D6-D6} \sim \tau_\alpha^i \Delta_{(ij)}^c \tau_\beta^j \epsilon^{\alpha\beta} + \tau_{i\alpha'}^c \Delta^{(ij)} \tau_{j\beta'}^c \epsilon^{\alpha'\beta'} . \tag{3.3}$$

These interactions are induced by mixed disk amplitudes, that emerge at the intersections between two $D6$-brane stacks and one $E2$ instanton as the one in figure 3.[9] In our case the two $D6$-branes are actually the $D6$-branes of the U(4) stack and their images. Integrating out the fermionic modulini produces two ϵ^{ijkl} so that schematically

$$\mathcal{W}_{E2} = \frac{1}{\mathcal{M}_0} \int d^8\tau e^{\mathcal{L}_{E2-D6-D6}} = \frac{1}{\mathcal{M}_0} \epsilon^{ijkl} \epsilon^{i'j'k'l'} \Delta_{ii'}^c \Delta_{jj'}^c \Delta_{kk'}^c \Delta_{ll'}^c . \tag{3.4}$$

The suppression scale $\mathcal{M}_0$ is related to the string scale M_S as $\mathcal{M}_0 \sim M_S e^{+S_{E2}}$, where S_{E2} depends on the closed string moduli that parametrize the (complex) size of the 3-cycles, wrapped by the $E2$-instantons.[10]

After the decomposition $\mathrm{SU}(4) \to \mathrm{SU}(3) \times \mathrm{U}(1)_{B-L}$, $\mathbf{10}_+ \to (\mathbf{6}_{+2/3}, \mathbf{3}_{-2/3}, \mathbf{1}_{-2})$, we denote by Δ_6 the 'diquark' super-field in the $\mathbf{6}$, T_3 in the triplet $\mathbf{3}$ and S the singlet $\mathbf{1}$ and find

$$\frac{1}{M_2} \phi_{\dot\alpha\beta}^c F_i^{c\dot\alpha} F_j^{c\dot\beta} \Delta^{c(ij)} \to \frac{1}{M_2} \phi_{\dot\alpha\beta} \left[Q_i^{c\dot\alpha} Q_j^{c\dot\beta} \Delta_6^{cij} + 2 Q_i^{c\dot\alpha} L^{c\dot\beta} T_3^{ci} + L^{c\dot\alpha} L^{c\dot\beta} S^c \right] \tag{3.5}$$

$$m_\Delta \Delta\Delta^c + \frac{1}{4M_4}(\Delta\Delta^c)^2 \to m_\Delta(\Delta_6\Delta_6^c + TT^c + SS^c) + \frac{1}{4M_4}(\Delta_6\Delta_6^c + TT^c + SS^c)^2 \tag{3.6}$$

$$\frac{1}{\mathcal{M}_0} \epsilon_{ijkl} \epsilon_{i'j'k'l'} \Delta^{cii'} \Delta^{cjj'} \Delta^{ckk'} \Delta^{cll'} \tag{3.7}$$

$$\to \frac{1}{\mathcal{M}_0} \left[4 \, \epsilon_{ijk}^{\mathrm{SU}(3)} \epsilon_{i'j'k'}^{\mathrm{SU}(3)} \Delta_6^{cii'} \Delta_6^{cjj'} \Delta_6^{ckk'} S + 6 \, \epsilon_{ijk}^{\mathrm{SU}(3)} \epsilon_{i'j'k'}^{\mathrm{SU}(3)} \Delta_6^{cii'} \Delta_6^{cjj'} T_3^{ck} T_3^{ck'} \right].$$

The complete super-potential after the decomposition $\mathrm{SU}(4) \to \mathrm{SU}(3) \times \mathrm{U}(1)_{B-L}$ reads

$$\mathcal{W} \sim y_1 h_{\alpha\dot\alpha} Q^{i\alpha} Q^{ci\dot\alpha} + y_1 h_{\alpha\dot\alpha} L^\alpha L^{c\dot\alpha} + \frac{1}{M_1} \phi_{\alpha\beta} \left(Q^{i\alpha} Q^{j\beta} \Delta_{ij} + Q^{i\alpha} L^\beta T_{3i} + L^\alpha L^\beta S \right) \tag{3.8}$$

$$+ \frac{1}{M_2} \phi_{\alpha\beta}^c (Q^{ci\alpha} Q^{cj\beta} \Delta_{ij}^c + Q^{ci\alpha} L^{c\beta} T_{3i}^c + L^{c\alpha} L^{c\beta} S^c) + \mu h_{\alpha\dot\alpha} h^{\alpha\dot\alpha}$$

$$+ \frac{1}{M_3} h^c \phi^c h\phi + m_\Delta(\Delta_6\Delta_6^c + TT^c + SS^c) + \frac{1}{4M_4}(\Delta_6\Delta_6^c + TT^c + SS^c)^2 + V(\phi_{L,R})$$

$$+ \frac{1}{\mathcal{M}_0} \left[4 \, \epsilon_{ijk}^{\mathrm{SU}(3)} \epsilon_{i'j'k'}^{\mathrm{SU}(3)} \Delta_6^{cii'} \Delta_6^{cjj'} \Delta_6^{ckk'} S + 6 \, \epsilon_{ijk}^{\mathrm{SU}(3)} \epsilon_{i'j'k'}^{\mathrm{SU}(3)} \Delta_6^{cii'} \Delta_6^{cjj'} T_3^{ck} T_3^{ck'} + c.c \right]$$

where $V(\phi_{L,R}) = m_{L,R}\phi_{L,R}^2 + a_{L,R}\phi_{L,R}^3/3$.

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[9]For similar calculations in related contexts, see [70–77].

[10]In general, the calculations could be much more complicated, in the presence of bulk fluxes, that can also induce soft susy breaking mass terms for the susy partners. For example gaugino mass terms with $M_\lambda \sim \Omega_{0,3}^{ijk} \langle \tau H_{ijk} + i F_{ijk} \rangle$ can be generated in Type IIB contexts by internal 3-form fluxes. In the presence of fluxes, one has to verify that physical branes and instantons are not lifted, i.e. the cycles they wrap and their intersections are not eliminated. With the introduction of bulk fluxes, one also has to consider the back-reactions on the exotic instantons, that could change then number of zero modes. This could modify our present analysis.

When $\langle\phi_{RR}\rangle = v_R$ and $\langle\phi_{LL}\rangle = 0$, Left-Right symmetry is spontaneously broken. When $\langle S\rangle = v_{B-L}$, SU(4) and its subgroup U(1)$_{B-L}$ are spontaneously broken. MB: contrary to U(1)$_4$ that extends SU(4) to U(4) and is anomalous, both U(1)$_{B-L}$ and U(1)$_Y$ are non-anomalous being associated to traceless generators of the non-abelian groups SU(4) and Sp(2)$_R$. A Majorana mass term for the Neutrino is generated[11] as $m_N \sim v_R v_{B-L}/M_2$. For example, $m_N \simeq 10^{12}$ GeV can be obtained if $v_R \simeq M_2$ and $v_{B-L} \simeq 10^{12}$ GeV. In this model the generation of a neutrino Majorana mass is connected to the generation of a Neutron Majorana mass.

In fact, when S takes an expectation value, a cubic interaction term

$$(v_{B-L}/\mathcal{M}_0)\epsilon_{ijk}^{\mathrm{SU(3)}}\epsilon_{i'j'k'}^{\mathrm{SU(3)}}\Delta_6^{cii'}\Delta_6^{cjj'}\Delta_6^{ckk'}$$

is generated. In section 4, we will discuss the consequences, for Neutron-Antineutron physics and for LHC phenomenology, in more in details.

3.1 No-proton decays

The quiver in figure 2 is constructed in such a way that, starting from a R-parity preserving susy Pati-Salam model, $E2, E2'$-instantons non-perturbatively generate only two R-parity violating operators: (3.3) and its complex conjugate $\mathcal{W}_{E2'} = \Delta\Delta\Delta\Delta/\mathcal{M}_0'$, respectively. For example, superpotentials like $\mathcal{W}_{ex1} = F_L F_L F_L F_L/\Lambda_1$ and $\mathcal{W}_{ex2} = F_R F_R F_R F_R/\Lambda_2$ are not generated, simply for construction: exotic instantons that could generate such superpotentials are not introduced in figure 2. As a consequence, there is no proton decay in our model. This is a nice feature of a dynamical R-parity breaking from Exotic Instantons, with respect to an explicit R-violation: we can generate specific superpotentials without other dangerous ones.

4 Neutron-Antineutron oscillation through color diquark sextets

In a susy PS-like model SU(4) $\times$ SU(2)$_R$ $\times$ SU(2)$_L$, we can construct a diagram like the one in figure 1 for Neutron-Antineutron oscillation, through the 'exotic' interaction $\Delta_{10}^c \Delta_{10}^c \Delta_{10}^c \Delta_{10}^c/\mathcal{M}_0$, containing

$$\mathcal{W}_{\Delta B=2} = \frac{1}{\mathcal{M}_0}\epsilon^{u^c d^c d^c \nu^c}\epsilon^{u'^c d'^c d'^c \nu'^c}\Delta_{u^c u'^c}^c \Delta_{d^c d'^c}^c \Delta_{d^c d'^c}^c S_{\nu^c \nu'^c}^c \tag{4.1}$$

(with $S_{\nu^c \nu^c}^c \equiv \Delta_{\nu^c \nu^c}^c$). The operator (4.1) induces a neutron-antineutron transition depicted in figure 2, as a result of the super-potential term $\tilde{f}_{11}v_R Q^c Q^c \Delta^c/M_2$, whose components include $f_{11}\Delta_{u^c u^c}^c u^c u^c$ and $f_{11}\Delta_{d^c d^c}^c d^c d^c$, with $f_{11} = \tilde{f}_{11}v_R/M_2$, and $\tilde{f}_{11}$ Yukawa couplings.

The process in figure 1-(b) produces an effective operator $G_{n-\bar{n}}(udd)^2$ with

$$G_{n-\bar{n}} \simeq \frac{g_3^2}{16\pi}\frac{f_{11}^2 v_{BL}}{M_{\Delta_{u^c u^c}^c}^2 M_{\Delta_{d^c d^c}^c}^2 M_{\mathrm{SUSY}}\mathcal{M}_0}. \tag{4.2}$$

We can now discuss different choices of the parameters leading to very different branches for phenomenology. The motivation of such a variety of possibilities is related

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[11]Dirac masses are generated via Yukawa couplings when h_{LR} gets a VEV.

to the fact that in (4.2) one can produce a scale of 300–1000 TeV, testable in the next generation of experiments, with different choices of the other parameters.

The cases with $\mathcal{M}_0 \simeq 10^{19}\,\text{GeV}$ and $\mathcal{M}_0 \simeq 10^{13}\,\text{GeV}$ are equivalent to the GUT SO(10) inspired scenario, discussed in [14] in figure 1-(b). In these cases $M_{\Delta_{u^c u^c}^c} \sim 1\,\text{TeV}$. Both cases are well compatible with the mechanism proposed in the previous section. In fact a scenario in which $\mathcal{M}_0 \simeq M_S$ can be envisaged, if $e^{S_{E_2}} \sim 1$ i.e. small 3-cycles wrapped by E_2 in CY. A priori, the string scale can be considered as a free parameter, it can be as high as $10^{19}\,\text{TeV}$ as low as a few TeV's. For instance, if $M_S = 1 \div 10\,\text{TeV}$, the hierarchy problem of the Higgs mass is automatically solved, and $\mathcal{M}_0$ can be as high as $10^{13}\,\text{GeV}$ (or more) if $e^{S_{E_2}} = 10^{10}$ i.e. for an E_2 wrapping a large 3-cycle in the CY.

In TeV-scale gravity scenari, one can also consider an alternative scenario in which $\mathcal{M}_0 \simeq M_S \simeq 10\,\text{TeV}$, with $e^{S_{E2}} \sim 1$. Compatibly with $n - \bar{n}$ limits, the four-sextets' interaction (3.2) is suppressed at much lower scales, with $\mathcal{M}_0 \sim 10\,\text{TeV}$. In this case, the masses of the sextets have to be much higher than $1\,\text{TeV}$, eluding completely a direct observation at LHC and in FCNCs. On the other hand, a post sphaleron baryogenesis scenario as the one proposed in [14] remains possible.

So, combined observations from the next generation of experiments on neutron-antineutron oscillations, LHC, FCNC physics, neutrino physics can provide precious informations not only on a possible PS extension on the SM and its region of the parameters, but also on the dynamical scale generated by the Exotic instantons, and as a consequence on the geometric structure of the Calabi-Yau compactifications, in particular the 3-cycle wrapped by the Exotic Instanton! This is a fundamental information for realistic model building in string phenomenology. Exotic instantons could be *portals* from low energy physics to the quantum gravity scale! In the next section, we will briefly discuss connections with dark matter and the hierarchy problem. Note that in our scenario, $\langle S \rangle \neq \langle \phi_{RR} \rangle$ in general. This is an important difference with respect to Babu-Mohapatra model cited above: $n - \bar{n}$ oscillation time of order $10^{10}\,\text{s}$ is compatible with a Left-Right symmetry restoration at TeV scale, with intriguing implications for LHC. A recent anomaly with significance near 3σ, compatible with Left-Right symmetry, in $pp \to l_1 l_2 jj$, was seen by CMS [80]. In a Left-Right model, this is interpreted as sequential W_R and N_R production as [81–84]

$$pp \to W_R \to l_1 N_R \to l_1 l_2 W_R^* \to l_1 l_2 jj\,.$$

However, this interpretation requires $g_R(M_{W_R}) \simeq 0.6 g_L(M_{W_R})$, with $1.8\,\text{TeV} < \text{M}_{\text{W}_\text{R}} < 2.4\,\text{TeV}$. Curiously, this situation is not compatible with a D-parity preserving SO(10) GUT scenario. For P-S models emerging from SO(10) $\to$ SU(4) $\times$ SU(2)$_L$ $\times$ SU(2)$_R$ D-parity is a symmetry. It is the external automorphisms that exchanges the two SU(2) groups, SU(2)$_L$ $\leftrightarrow$ SU(2)$_R$, and at the same time acts by conjugation on SU(4) representations. More explicitly $D_P = \Gamma_7 \gamma_5$ is a symmetry when $g_L = g_R$. At the unification scale $g_L = g_R$, the two coupling constants have the same running if the field content is LR symmetric. For a LR interpretation of CMS anomaly one needs $g_L \neq g_R$, compatibly with a PS model not emerging from SO(10) without D-parity altogether or at least an SO(10) model where D-parity is broken at a high scale. On the other hand, D-parity is not a symmetry in string-

inspired models. The gauge couplings depend on the size of the internal cycle wrapped by the D-branes. In orbifolds or Calabi-Yau singularities, one can tune the blow-up modes so that different cycle have the same size e.g. vanishing, but generically this is not the case and one can start with $g_L \neq g_R$ already at the string scale.

$pp \to l_1 l_2 jj$ is not the only peculiar channel suggested by our model for LHC, and we will return to other ones in the next section. We conclude this section with another observation. A scenario in which S doesn't take an expectation value at all can be envisaged in our model. In this case, U(1)$_{B-L}$ is not spontaneously broken by S. However, as mentioned above, exotic intantons can dynamically break U(1)$_{B-L}$. For instance, a Majorana mass matrix for RH neutrini can be generated by exotic instantons rather than by S, as cited above. In this case, S could also be a light particle, if a residual discrete symmetry of U(1)$_{B-L}$ stabilizes it. In other words, S can behave as a Majoron, but it is not exactly a Majoron [85]. We can call it an *exoticon*. We suppose that the exoticon interacts with the three color sextets with a coupling μ_S. So, in this case, we have to replace v_{B-L} with μ_S in (4.2). The exoticon carries $B - L = -2$, so $n \to \bar{n}S$ does not violate $B - L$. We also note another important difference with respect to Majorons: Majoron mass $m_\phi = y_L v_L$, with v_L vev of a global U(1)$_L$ (and y_L coupling), is related to its interaction with neutrini, as $g_{\phi\nu\nu} = m_\nu/v_L$; such a relation, in general, is not satisfied by exoticons. A massive exoticon cannot be emitted in a $n-\bar{n}$ transition, in the vacuum: CPT symmetry protects neutron by transitions $n \to \bar{n} + S$, i.e. $m_n = m_{\bar{n}}$. However, in a nuclear environment, such a transition is allowed! Such a transition is followed by annihilation of the antineutron with another neutron in the nuclear environment, as $(Z, A) \to (Z, A - 2) + 3\pi$. We can roughly estimate the corresponding decay width as $\Gamma \simeq (\delta m/\mu_S)^2 \langle \Delta \mathcal{E} \rangle$, where $\langle \Delta \mathcal{E} \rangle \simeq 10 \div 100\,\mathrm{MeV}$ is the average energy in the nuclear environment. Limits on $n - \bar{n}$ oscillation in the nuclei are $\Gamma_{n\bar{n}}^{-1} \sim 10^{-32}\,\mathrm{yr}$ [86], corresponding to $\mu_S > 10^{30}\delta m \simeq \mathrm{keV}$. Another spectacular signature of an exoticon could be a nucleon-nucleon disappearances as $nn \to S$, $\Delta B = 2$. This could be detected as a nuclear transition $(Z, A) \to (Z, A - 2) + \mathrm{missing\,energy}$. We can easily estimate the rate of such a transition as $\Gamma \sim \kappa_{np}(\mathcal{M}_{n\bar{n}})^{-3}m_N^{14}G_{n-\bar{n}}^2\,\mathrm{GeV}$, where $\kappa_{np} \sim 10^{-6}$ approximately accounts for the hadronic non-perturbative correction. Such an estimate leads us to conclude that such a process is very suppressed, roughly as $10^{40}\,\mathrm{yr}$. Finally, an exoticon can be also detected in a neutrinoless double-beta-decay, as a Majoron. However, there are several important differences with respect to the Majoron: a $0\nu2\beta + S$ does not violate lepton number, it is an *apparent violation*. For a $0\nu2\beta + S$ process, limits on the exoticon production imply $(m_\nu/\mu_S) < 10^{-5}$ [87, 88] that corresponds to a bound $\mu_S > 10\,\mathrm{keV}$. Limits from supernovae cooling processes $\nu \to \nu^c S$ or $\nu\nu \to SS$ are competitive (for electronic neutrini $m_\nu/\mu_S < 10^{-5}$) [89].

5 Other comments on phenomenology, dark matter and hierarchy problem

In Minimal non-susy Left-Right models, identifying a candidate for cold dark matter is difficult. On the other hand, our model automatically suggests several candidates for cold dark matter. In fact, our model predicts the presence of neutralini and stuckelberg axini

(or stuckelini), mixing with each other, as in [90, 91]. They are good candidate for WIMP dark matter. On the other hand, extra Z' for each anomalous U(1) can get a mass from Stuckelberg mechanism, that is not necessary of the order of string scale M_S. In particular, a scenario in which $M_S \simeq 10 \div 10^3$ TeV can be envisaged, alleviating hierarchy problem of Higgs mass. In this case, $m_{Z'} \simeq 10^{-4} M_S \div M_S \simeq 1 \div 10$ TeV, with Generalized Chern Simon terms inducing peculiar decays as $Z' \to Z\gamma$ [25–32, 90–92].[12] Curiously, the presence of apparent flavor violations in $B \to llK$, detected in LHCb [100–103] could be a hint in favor of a new Z' [104–106]. In our model, we also predict extra decays $B \to l^+l^-l'^+l'^- K$ suppressed by the GCS couplings with respect to $B \to l^+l^- K$. Extra Z' from anomalous symmetries are different with respect to Z'_R (Z-boson of the SU(2)$_R$), and kinetic mixings $Z - Z'$ or $Z' - Z'_R$ can be envisaged. Z'_R can also interact Z, γ, Z', γ through G.C.S, and a complete study, of resulting cascade processes, is beyond the purpose of this paper. Concerning the hierarchy problem of the Higgs mass, our model is compatible with TeV-scale supersymmetry, but this model has more undetermined parameters than the MSSM, i.e. it could be more elusive and more difficult to constrain.

6 Conclusions and remarks

In this paper, we have shown how to generate a Majorana mass for the neutron, inducing neutron-antineutron transitions with $|\Delta B| = 2$ in the context of a Pati-Salam Left-Right model. Indeed exotic instantons can produce an effective interaction involving color diquark sextets, leading to a Majorana mass term for the neutron. We would like to stress that in the present context no processes with $|\Delta B| = 1$ are allowed that could lead to fast proton decay. We have discussed some possible phenomenological implications, in the main branches of the parameters space, for LHC, FCNC, Dark Matter, $0\nu 2\beta$-decay. A unifying picture of Dark Matter, Hierarchy problem of the Higgs mass, Baryogenesis and Neutrino mass emerges in a very simple unoriented quiver! In this sense, the model elaborated here could represent a serious alternative to GUT inspired models.

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[12]We would like to stress that GCS terms generate UV divergent triangles that are cured by considering UV completions with KK states or string excitations. For issues in scattering amplitudes and collider physics see [96], for recent discussion about string theory and causality see [94, 95]. For non-local (string inspired) quantum field theories see [97–99].

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