Phys. Lett. B 879 (2026) 140586

Contents lists available at ScienceDirect

Physics Letters B

journal homepage: www.elsevier.com/locate/physletb

Letter

# Neutrino mass induced $n$–$\bar{n}$ oscillation

Ilja Doršner [a,b], Svjetlana Fajfer [b], Shaikh Saad [b,*]

[a] Department of Physics, Faculty of Science, University of Zagreb, Zagreb, HR-10000, Croatia
[b] Jožef Stefan Institute, Jamova 39, P. O. Box 3000, Ljubljana, SI-1001, Slovenia

ARTICLE INFO

Editor: Tuhin Roy

*Keywords:*
Neutrino mass
Neutron-antineutron oscillations
Unified theories

ABSTRACT

The Georgi-Glashow model is the simplest possible attempt at grand unification. However, due to its particle content, the model preserves a global $U(1)_{B-L}$ symmetry, where $B$ and $L$ are baryon and lepton numbers, respectively. It thus leaves neutrinos massless just as the Standard Model does. The extensions of the Georgi-Glashow model that break lepton number by two units, i.e., scenarios with $|\Delta L| = 2$ operator(s), naturally generate potentially viable Majorana neutrino masses. Since the dynamics that yields $|\Delta L| = 2$ interactions unavoidably induces $|\Delta B| = 2$ transitions via $B-L$ breaking, these extensions consequentially lead to intriguing processes such as neutron–antineutron ($n$–$\bar{n}$) oscillation. We investigate this intrinsic connection between neutrino mass generation and $n$–$\bar{n}$ oscillation within a number of representative extensions of the Georgi-Glashow model that can yield realistic neutrino masses and mixing parameters. These extensions include the tree-level seesaw mechanism realizations of the Type I, II, and III varieties, as well as the one-loop and two-loop radiative neutrino mass models.

## 1. Introduction

The Georgi–Glashow (GG) model [1] is the simplest realization of grand unification. It achieves quark–lepton unification as well as unification of the Standard Model (SM) interactions within $SU(5)$. However, since it predicts neutrinos to be massless, it leaves the origin of neutrino masses as one of the open questions that needs to be addressed.

We establish an intrinsic connection between Majorana neutrino mass generation in $SU(5)$ and $n$–$\bar{n}$ oscillation, both rooted in quark–lepton unification. Our work expands on a recent analysis [2] of two simple GG model extensions of the Type II and Type III varieties that simultaneously achieve observable $n$–$\bar{n}$ oscillation [3–6] and accommodate realistic SM fermion mass spectrum. More specifically, we apply approach advocated in Ref. [2] to other tree-level and radiative neutrino mass generation models to champion an unavoidable correlation between neutrino mass and $n$–$\bar{n}$ oscillation. Our work also dovetails a study [7] which enumerates, in scalar extensions of the GG model, all $d = 9$ operators contributing to $n$–$\bar{n}$ oscillation.

Our results are important because they imply that an observation of baryon number–violating processes such as $n$–$\bar{n}$ oscillation could provide an independent evidence for Majorana neutrinos, beside the neutrinoless double beta decay [8], if $SU(5)$ is realized in nature. Our focus on $n$–$\bar{n}$ oscillation is particularly timely given the current stringent bound from `Super-Kamiokande` [9] and upcoming searches at `DUNE` [10] and `NNBAR` [6] at `ESS` that are expected to significantly improve existing limits. Additional experimental handles include $|\Delta B| = 1$ two-body proton decays, $|\Delta L| = 2$ processes such as muon to positron conversion [11], and $|\Delta B| = 2$ transitions like di-nucleon decays [12–15]. A schematic summary of these implications is shown in Fig. 1.

In Section 2, we demonstrate that quark-lepton unification and Majorana neutrino mass generation mechanisms, in $SU(5)$, automatically yield baryon-number-violating interactions with $|\Delta B| = 2$. Section 3 presents the construction and discussion of neutrino mass mechanisms and their implications for $n$–$\bar{n}$ oscillation. We briefly conclude in Section 4.

## 2. Implications of Quark-Lepton unification and majorana neutrinos

The GG model consists of the following representations:

$$\bar{5}_{F\,i} = \left(d^c_i, L_i\right),\ \ 10_{F\,i} = \left(Q_i, u^c_i, e^c_i\right),\ \ 5_S = (T, H),\ \ 24_S, \tag{1}$$

where subscripts $F$ and $S$ denote fermions and scalars, respectively, and $i = 1, 2, 3$ is a family index. Here, $Q_i$, $u^c_i$, and $d^c_i$ are the SM quark multiplets, $L_i$ and $e^c_i$ represent leptonic multiplets, $T$ stands for the color triplet, while $H$ is the SM Higgs doublet. Given the above particle content, $U(1)_{B-L}$ remains an accidental global symmetry of the model [7,16–20]. To demonstrate this, let us look at the Yukawa part of the

* Corresponding author.
*E-mail addresses:* dorsner@phy.hr (I. Doršner), svjetlana.fajfer@ijs.si (S. Fajfer), shaikh.saad@ijs.si (S. Saad).

https://doi.org/10.1016/j.physletb.2026.140586
Received 17 March 2026; Received in revised form 11 May 2026; Accepted 21 May 2026
Available online 26 May 2026

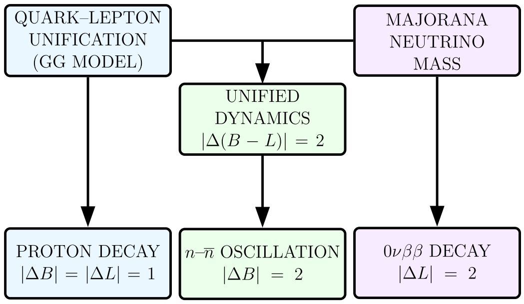



**Fig. 1.** Schematic illustrating the interplay of quark–lepton unification and neutrino mass generation with proton decay, $n$–$\bar{n}$ oscillation, and neutrinoless double beta decay ($0\nu\beta\beta$).

Lagrangian

$$\mathscr{L}_Y = 10_F\bar{5}_F 5^*_S + 10_F 10_F 5_S \to (QL + u^c d^c)T^* + (QQ + u^c e^c)T + (e^c L + Qd^c)H^* + (Qu^c)H. \quad (2)$$

One can trivially check that interactions of Eq. (2) preserve a global $U(1)_{B-L}$ symmetry with

$$Q_{+1/3},\ u^c_{-1/3},\ d^c_{-1/3},\ L_{-1},\ e^c_{+1},\ T_{-2/3},\ H_0, \quad (3)$$

where the subscripts represent the $U(1)_{B-L}$ charges. To correct the charged fermion mass relations in the GG model, one can minimally extend it by introducing a single $45_S$ [21]. But, even with the additional Yukawa interactions, $U(1)_{B-L}$ remains unbroken [7]. Note that $U(1)_{B-L}$ conservation allows for two-body proton decays such as $p \to e^+\pi^0$, where both initial and final states have $B-L=+1$.

Introduction of Majorana neutrino masses, which requires $|\Delta L|=2$, automatically breaks $U(1)_{B-L}$ by two units, thus inducing $|\Delta B|=2$ processes. We can illustrate this via an effective field theory approach [22]. Specifically, an analogue of the $d=5$ Weinberg operator within the GG framework yields

$$\mathscr{L}_5 = \frac{C_{ij}}{\Lambda}\bar{5}_{F\,i}\bar{5}_{F\,j}5_S 5_S \to \underbrace{\left(\frac{C^{\nu\nu}_{ij}}{\Lambda_1}\underbrace{L_i L_j H H}_{\Delta L=+2} + \frac{C^{dd}_{ij}}{\Lambda_2}\underbrace{d^c_i d^c_j T T}_{\Delta B=-2}\right)}_{\Delta(B-L)=-2}, \quad (4)$$

where $\Lambda$ denotes a new physics scale and $C_{ij}$'s are the associated Wilson coefficients. We accommodate model dependence of relevant scales and couplings in the $SU(5)$ broken phase through introduction of $\Lambda_{1,2}$ and $C^{\nu\nu,dd}_{ij}$. The dynamics generating Majorana neutrino masses clearly induces $|\Delta B|=2$ processes. Namely, the last term in Eq. (4) generates $n$–$\bar{n}$ oscillation through the following $d=9$ operator:

$$\frac{C^{dd}_{11}}{\Lambda_2}d^c_1 d^c_1 TT \to \frac{C^{dd}_{11}C^{ud}_{11}}{\Lambda_2\Lambda_3^4}d^c_1 d^c_1\left(u^c_1 d^c_1\right)\left(u^c_1 d^c_1\right). \quad (5)$$

To better understand advocated connection between quark-lepton unification and Majorana neutrino mass generation within the framework of renormalizable $SU(5)$ models, let us consider the following two examples:

- **Type I seesaw.** If one introduces a Majorana fermion $1_F = \nu^c(1,1,0)$, the Yukawa sector

  $$\mathscr{L}_Y \supset y_i\,\bar{5}_{F\,i}1_F 5_S \to y^\nu_i\, L_i \nu^c H + y^d_i\, d^c_i \nu^c T, \quad (6)$$

  requires $\nu^c$ to carry $B-L$ charge of $+1$. Therefore, its Majorana mass term

  $$\mathscr{L}_Y \supset \underbrace{M^\nu\,\nu^c\nu^c}_{\Delta(B-L)=-2}, \quad (7)$$

  breaks $B-L$ by two units.

- **Type II seesaw.** This case requires $15_S = \Delta_1(1,3,1) + \Delta_3(3,2,1/6) + \Delta_6(\bar{6},1,-2/3)$. The Yukawa interaction reads

  $$\mathscr{L}_Y \supset y_{ij}\,\bar{5}_{F\,i}\bar{5}_{F\,j}15_S \to y^\nu_{ij}\,L_i L_j \Delta_1 + y^{Ld}_{ij}\,L_i d^c_j \Delta_3 + y^{dd}_{ij}\,d^c_i d^c_j \Delta_6, \quad (8)$$

  which assigns $B-L$ charges of $-2$, $+4/3$, and $+2/3$ to $\Delta_1$, $\Delta_3$, and $\Delta_6$, respectively. The scalar potential includes

  $$V \supset \mu\, 5_S 5_S 15^*_S \to \underbrace{\overline{\mu}\,HH\Delta^*_1 + \hat{\mu}\,HT\Delta^*_3 + \widetilde{\mu}\,TT\Delta^*_6}_{\Delta(B-L)=-2}, \quad (9)$$

  where each term breaks $B-L$ by two units.

This analysis yields the same conclusion in any other renormalizable $SU(5)$ model of Majorana neutrino mass.

## 3. Models

We focus on minimal renormalizable extensions of the GG model that generate phenomenologically viable Majorana neutrino masses and, consequentially, yield $n$–$\bar{n}$ oscillation signatures. Here, minimality refers both to the number of employed representations beyond those of the GG model as well as the dimensionality of these additional representations. More precisely, we do not consider extensions with more than two additional representations, where their dimensionality should not exceed fifty. The GG model extensions relevant for our work are listed in Table 1. The associated neutrino mass diagrams are presented in the upper panels of Fig. 2 through Fig. 6, whereas the corresponding $n$–$\bar{n}$ topologies are given in the lower panels of Fig. 2 through Fig. 6.

The second column of Table 1 lists the particle content of each scenario, the third column gives the neutrino mass generation order, and the fourth and fifth columns specify the class and topology of $n$–$\bar{n}$ oscillation. The last column lists the mediators of the $d=9$ operators responsible for $n$–$\bar{n}$. The **CLASS** indicates the $U(1)_{B-L}$ breaking mechanism. **CLASS I** requires at least one Majorana fermion, **CLASS II** requires scalar(s), while **CLASS III** requires both a Dirac fermion and a scalar. If the model generates multiple contributions to $n$–$\bar{n}$ oscillation, we lists in Table 1 particle content yielding the dominant transition rate.

**Type I seesaw model**: One needs at least two right-handed neutrinos to generate viable neutrino mass parameters in the Type I seesaw. The interplay between a scalar $(\bar{3},1,1/3) \in 5^*_S$ and a fermion $1_F$ produces the $n$–$\bar{n}$ diagram of **TOPOLOGY B**, as shown in the lower panel of Fig. 2. The fermion featured in **TOPOLOGY B** is always of Majorana nature, and the $|\Delta(B-L)|=2$ breaking arises from its Majorana mass. Scalar $(\bar{3},1,1/3) \in 5^*_S$ can mediate proton decay and its contribution towards $n$–$\bar{n}$ signal is expected to be suppressed.

**Type II seesaw model**: The Type II seesaw corresponds to **CLASS II**, with the $n$–$\bar{n}$ transition via **TOPOLOGY A** that is mediated by $(\bar{3},1,1/3) \in 5^*_S$ and $(6,1,-2/3) \in 15_S$. The neutrino mass and $n$–$\bar{n}$ diagrams are featured in Fig. 3a and b, respectively.

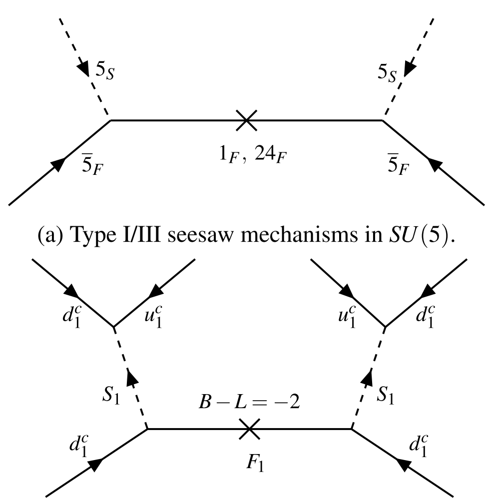
**Type II seesaw variant**: This interesting variant employs two color sextets, $S_1 = (6,1,1/3) \in 45^*_S$ and $(6,1,-2/3) \in 15_S$, to mediate $n$–$\bar{n}$ oscillation. Since these sextets do not induce proton decay, the oscillation time can be within the observable range [2]. The neutrino mass and $n$–$\bar{n}$ diagrams appear in the upper and lower panels of Fig. 3, respectively.

**Type III seesaw model**: One requires two copies of $24_F$ to generate viable neutrino mass parameters even though a single $24_F$ already contains two states acting as right-handed neutrinos. The $B-L$ symmetry is broken by the Majorana mass of $24_F$. The neutrino mass and $n$–$\bar{n}$ diagrams are shown in the upper and lower panels of Fig. 2, respectively.

**Type III seesaw variant**: One can replace one $24_F$ with $45_S$ to achieve consistency with neutrino oscillation data. The $n$–$\bar{n}$ mediators, $(6,1,1/3) \in 45^*_S$ and $(8,1,0) \in 24_F$, do not induce proton decay and can enhance the oscillation rate [2]. Mediators of $n$–$\bar{n}$ oscillation can also be $(\bar{3},3,1/3) \in 45^*_S$ and $(1,3,0) \in 24_F$, but the relevant couplings of the scalar multiplet are antisymmetric in flavor space and cannot generate

**Table 1**
Summary of the viable neutrino mass models and the associated particle content within the GG framework of quark-lepton unification. The third column indicates the order of the neutrino mass generation. The fields listed in the sixth column correspond to those providing the leading $n$–$\bar{n}$ contribution, with their $SU(3)\times SU(2)\times U(1)$ transformation properties explicitly shown in the second line. Superscript $q$ in $F^q$ represents the type of quark $q$ that fermion $F$ interacts with, while $S^{q\hat{q}}$ represents a scalar with an interaction of the $S^{q\hat{q}}q\hat{q}$ type. The $n$–$\bar{n}$ diagram topologies are denoted with **A**, **AI**, **AII**, **B**, and **C**, where all these topologies can be found in lower panels of Figs. 2–6.

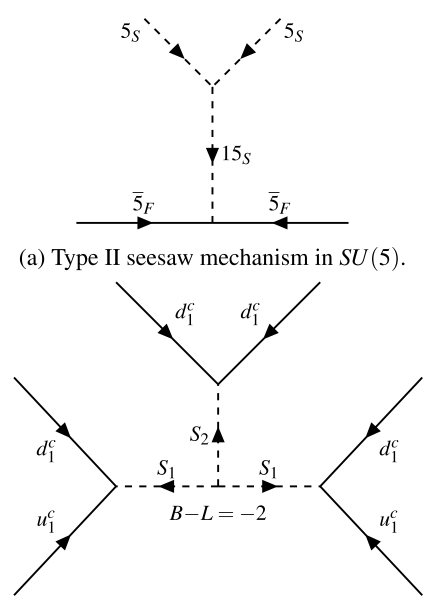
| Model | Content | $\nu$ Mass | Class | Topology | $n$–$\bar{n}$ Mediators |
|---|---|---|---|---|---|
| Georgi–Glashow (GG) [1] | $\overline{5}_{F\,i}+10_{F\,i}+5_S+24_S$ | – | – | – | – |
| Type I seesaw [23–27] | $GG+1_F+1_F$ | tree level | **I** | **B** | $S_1^{u^c d^c}=5_S^*$, $F_1^{d^c}=1_F$<br>$S_1=(\overline{3},1,1/3)$, $F_1=(1,1,0)$ |
| Type II seesaw [28–31] | $GG+15_S$ [2,32–41] | tree level | **II** | **A** | $S_1^{u^c d^c}=5_S^*$, $S_2^{d^c d^c}=15_S$<br>$S_1=(\overline{3},1,1/3)$, $S_2=(6,1,-2/3)$ |
| Type II seesaw variant | $GG+15_S+45_S$ [2] | tree level | **II** | **A** | $S_1^{u^c d^c}=45_S^*$, $S_2^{d^c d^c}=15_S$<br>$S_1=(6,1,1/3)$, $S_2=(6,1,-2/3)$ |
| Type III seesaw [42] | $GG+24_F+24_F$ [32,43–46] | tree level | **I** | **B** | $S_1^{u^c d^c}=5_S^*$, $F_1^{d^c}=24_F$ $S_1=(\overline{3},1,1/3)$, $F_1=(1,1,0),(8,1,0)$ |
| Type III seesaw variant | $GG+24_F+45_S$ [2] | tree level | **I** | **B** | $S_1^{u^c d^c}=45_S^*$, $F_1^{d^c}=24_F$<br>$S_1=(6,1,1/3)$, $F_1=(8,1,0)$ |
| BNT [47] | $GG+15_F+35_S$ [48–51] | tree & one-loop | **III** | **C** | $S_1^{u^c d^c}=5_S$, $S_2=35_S$, $F_1^{d^c}=15_F$<br>$S_1=(3,1,-1/3)$, $S_2=(\overline{10},1,1)$,<br>$F_1=(6,1,-2/3)$ |
| Zee [52] | $GG+10_S+45_S$ [53–55] | one-loop | **II** | **AI** | $S_1^{u^c d^c}=45_S^*$, $S_2^{d^c d^c}=10_S$<br>$S_1=(6,1,1/3)$, $S_2=(\overline{3},1,-2/3)$ |
| Zee-Babu [56,57] | $GG+10_S+50_S$ [58,59] | two-loop | **II** | **AII** | $S_1^{d^c d^c}=10_S$, $S_2^{u^c u^c}=50_S$<br>$S_1=(\overline{3},1,-2/3)$, $S_2=(6,1,4/3)$ |
| Zee-Babu variant | $GG+15_S+50_S$ | tree & two-loop | **II** | **A** | $S_1^{d^c d^c}=15_S$, $S_2^{u^c u^c}=50_S$<br>$S_1=(6,1,-2/3)$, $S_2=(6,1,4/3)$<br>$S_1^{\overline{QQ}}=50_S^*$, $S_2^{d^c d^c}=15_S$<br>$S_1=(6,3,1/3)$, $S_2=(6,1,-2/3)$ |




(a) Type I/III seesaw mechanisms in $SU(5)$.

(b) **TOPOLOGY B** for $n$–$\bar{n}$ oscillation.

**Fig. 2.** Type I/III seesaws (upper panel) and the associated topology, i.e., **TOPOLOGY B**, behind $n$–$\bar{n}$ oscillation (lower panel).




(a) Type II seesaw mechanism in $SU(5)$.

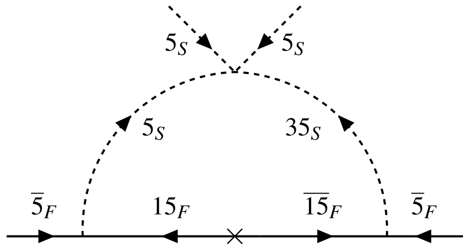
(b) **TOPOLOGY A** for $n$–$\bar{n}$ oscillation.

**Fig. 3.** Type II seesaw mechanism (upper panel) and the associated topology, i.e., **TOPOLOGY A**, behind $n$–$\bar{n}$ oscillation (lower panel).

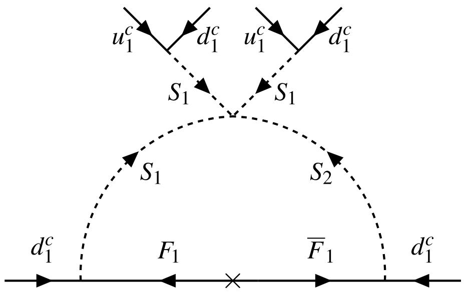
the leading-order contribution. We accordingly omit this particular scenario from Table 1.

**BNT model**: The Babu–Nandi–Tavartkiladze (BNT) model is the only model in Table 1 employing a Dirac fermion to generate both neutrino masses and $n$–$\bar{n}$ oscillation. The upper panel of Fig. 4 shows the dominant contribution to the neutrino mass. A tree-level term also exists but it is negligible [48]. The $n$–$\bar{n}$ mediators are $(3,1,-1/3)\in 5_S$, $(\overline{10},1,1)\in 35_S$, and $(6,1,-2/3)\in 15_F$, corresponding to **TOPOLOGY C.**

**Zee model**: The $GG+10_S$ model can generate neutrino masses and $n$–$\bar{n}$ process but cannot accommodate realistic neutrino oscillation data.

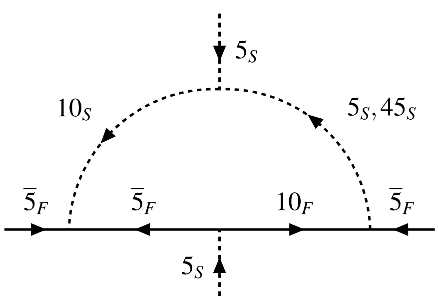



(a) BNT mechanism in $SU(5)$.

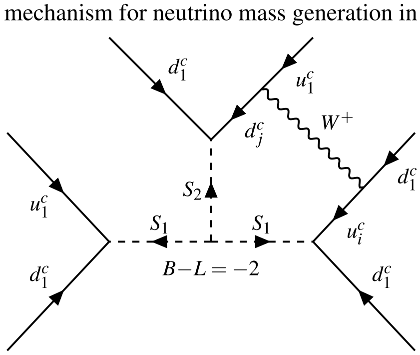



(b) **TOPOLOGY C** for $n$–$\bar{n}$ oscillation in the BNT embedding.

**Fig. 4.** Babu–Nandi–Tavartkiladze (BNT) mechanism (upper panel) and associated $n$–$\bar{n}$ oscillation diagram (lower panel).




(a) Zee mechanism for neutrino mass generation in $SU(5)$.

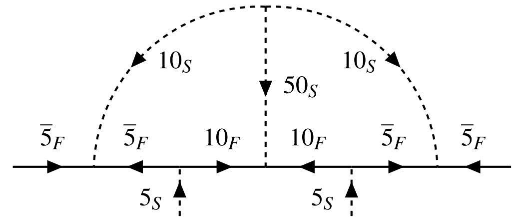



(b) **TOPOLOGY AI** for $n$–$\bar{n}$ oscillation with $i = 1,2,3$ and $j = 2,3$.

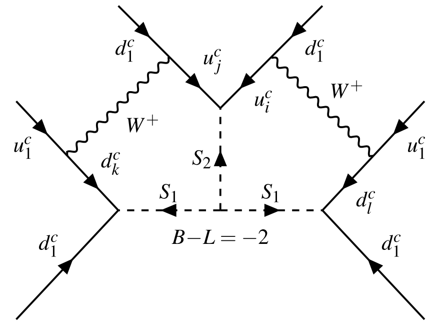
**Fig. 5.** Zee mechanism (upper panel) and the associated topology, i.e., **TOPOLOGY AI**, behind $n$–$\bar{n}$ oscillation (lower panel).

The GG + $10_S + 45_S$ extension resolves this. Here, the $n$–$\bar{n}$ mediators are $S_1 = (6,1,1/3) \in 45_S^*$ and $(\overline{3},1,-2/3) \in 10_S$. Due to antisymmetric nature of Yukawa couplings of $10_S$, a $W$-boson exchange is required for the transition. The neutrino mass and $n$–$\bar{n}$ diagrams are shown in the upper and lower panels of Fig. 5, respectively, with the latter corresponding to **TOPOLOGY AI**.

**Zee-Babu model**: In the Zee–Babu model, both neutrino masses and $n$–$\bar{n}$ oscillation arise at two loops, as shown in Fig. 6. The $n$–$\bar{n}$ medi-




(a) Zee-Babu mechanism for neutrino mass generation.




(b) **TOPOLOGY AII** for $n$–$\bar{n}$ oscillation with $i,j = 1,2,3$ and $k,l = 2,3$.

**Fig. 6.** Zee-Babu mechanism in $SU(5)$ (upper panel) and the associated topology behind $n$–$\bar{n}$ oscillation (lower panel).

ators are $S_1 = (\overline{3},1,-2/3) \in 10_S$ and $(6,1,4/3) \in 50_S$, corresponding to **TOPOLOGY AII**. Due to antisymmetric nature of Yukawa coupling of $10_S$ and the fact that scalar cubic interaction involves two $10_S$ propagators, two $W$-boson exchanges are required to generate $n$–$\bar{n}$ transition.

**Zee-Babu variant**: We note that $10_S$ of the minimal Zee–Babu model can be replaced by $15_S$, which would enhance the $n$–$\bar{n}$ rate due to introduction of symmetric Yukawa couplings of $15_S$ to the SM fermions. There would thus be no need for two $W$-boson exchanges to generate $n$–$\bar{n}$ oscillation. In this particular variant a singly charged component within $(1,3,1) \in 15_S$ replaces $(1,1,1) \in 10_S$ in the Zee–Babu two-loop neutrino mass diagram of Fig. 6a.

This variant also adds a tree-level contribution to neutrino masses, which might be subdominant in parts of parameter space. For example, this could be trivially achieved by assigning a tiny vacuum expectation value to scalar $(1,3,1) \in 15_S$. If that is the case, the leading $n$–$\bar{n}$ contributions would be via **TOPOLOGY A** of Fig. 3b and could arise entirely via color sextets, with two mediator options: $(6,1,-2/3) \in 15_S$ and $(6,1,4/3) \in 50_S$, or $(6,3,1/3) \in 50_S^*$ and $(6,1,-2/3) \in 15_S$. The latter is the only case with four external left-handed quarks. While both contributions can yield unsuppressed $n$–$\bar{n}$ transition rates we assume, for simplicity, that the former contribution dominates as it involves only right-handed SM fermions in the external legs. This particular choice of dominance does not affect the overall conclusions presented below.

Note that for all minimal, phenomenologically viable neutrino mass models, the neutrino mass and corresponding $n$–$\bar{n}$ diagrams both occur at the same perturbative order. Our discussion sets aside issues such as charged fermion mass generation and the viability of gauge coupling unification. A study addressing these aspects and their experimental implications will be presented elsewhere. We will, however, estimate whether a viable neutrino mass scale is compatible with the experimentally accessible $n$–$\bar{n}$ signal.

The fields responsible for the neutrino mass generation are specified in Table 2. To generate experimentally viable neutrino mass scale [64] one requires $\Lambda_1/|C^{\nu\nu}_{\tau\tau}| \simeq 6\times10^{14}$ GeV, whereas current bounds from `KamLAND-Zen` [65] and the projected reach of `nEXO` [66] on $0\nu\beta\beta$ imply $6\times10^{15}$ GeV $\gtrsim \Lambda_1/|C^{\nu\nu}_{ee}| \gtrsim 2.5\times10^{14}$ GeV. We also list in Table 2 the scales and associated couplings relevant for $n$–$\bar{n}$ oscillation, as intro-

**Table 2**
The mass scales and couplings relevant for neutrino masses and $n$–$\bar{n}$ process, where all the mass parameters are given in TeV units. $V_{ij}$ are the quark mixing parameters and $g_2$ is the $SU(2)$ gauge coupling of the SM. PD denotes viability with respect to proton decay limits.

| Model | $C^{\nu\nu}_{\tau\tau}$ | $\Lambda_1$ | $C^{dd}_{11} C^{ud}_{11}$ | $\Lambda_2 \Lambda_3^4$ | (NNBAR) $\geq \Lambda_2\Lambda_3^4 \geq$ (SK) | $F_k/S_k$ | LHC | PD |
|---|---|---|---|---|---|---|---|---|
| Type I | $(y^\nu)^2$ | $M_{F_1} \simeq 6\times10^{11}$ | $(y^{dd})^2(y^{ud})^2$ | $M_{F_1} M_{S_1}^4$ | $3 \geq M_{S_1} \geq 1.8$ | - | ✗ | ✗ |
| Type II<br>Type II variant | $y^\nu$ | $\frac{M_{S_3}^2}{\tilde{\mu}} \simeq 6\times10^{11}$ | $y^{dd}(y^{ud})^2$ | $\frac{M_{S_1}^4 M_{S_2}^2}{\tilde{\mu}}$ | $4.2\times10^3 \geq M_{S_2}^{1/4} M_{S_1} \geq 2.6\times10^3$ | $S_3 : (1,3,1)$ | ✓<br>✓ | ✗<br>✓ |
| Type III<br>Type III variant | $(y^\nu)^2$ | $M_{F_1} \simeq 6\times10^{11}$ | $(y^{dd})^2(y^{ud})^2$ | $M_{F_2} M_{S_1}^4$ | $2.6\times10^3 \geq M_{F_2}^{1/4} M_{S_1} \geq 1.6\times10^3$ | $F_2 : (8,1,0)$ | ✓<br>✓ | ✗<br>✓ |
| BNT | $\frac{y^\nu \hat{y}^\nu}{16\pi^2}$ | $\frac{M_{S_3}^2}{M_{F_2}} \simeq 6\times10^{11}$ | $\frac{y^d \hat{y}^d (y^{ud})^2}{16\pi^2}$ | $\frac{M_{S_1}^4 M_{S_2}^2}{M_{F_1}}$ | $1.3\times10^3 \geq \frac{M_{S_1} M_{S_2}^{1/2}}{M_{F_1}^{1/4}} \geq 8.3\times10^2$ | $F_2 : (1,3,1)$<br>$S_3 : (1,4,-3/2)$ | ✓ | ✗ |
| Zee | $\frac{y^\nu \hat{y}^\nu y_\tau}{16\pi^2}$ | $\frac{M_{S_3}^2}{\tilde{\mu}} \simeq 4\times10^7$ | $\frac{g_2^2 V_{ub} V_{td}}{16\pi^2} \frac{m_b m_t}{m_W^2} y^{ds}\left(y^{ud}\right)^2$ | $\frac{M_{S_1}^4 M_{S_2}^2}{\tilde{\mu}}$ | $125 \geq M_{S_1} M_{S_2}^{1/4} \geq 77$ | $S_3 : (1,1,1)$ | ✓ | ✓ |
| Zee-Babu<br>Zee-Babu varian | $\frac{y^\nu \hat{y}^\nu y_\mu^2}{256\pi^4}$ | $\frac{M_{S_3}^2}{\tilde{\mu}} \simeq 140$ | $\frac{g_2^4 V_{ub}^2 V_{td}^2}{256\pi^4} \frac{m_b^2 m_t^2}{m_W^4} \left(y^{ds}\right)^2 \; y^{uu}$<br>$y^{uu}(y^{dd})^2$ | $\frac{M_{S_1}^4 M_{S_2}^2}{\tilde{\mu}}$ | $3.7 \geq M_{S_2}^{1/4} M_{S_1} \geq 2.3$<br>$4.2\times10^3 \geq M_{S_2}^{1/4} M_{S_1} \geq 2.6\times10^3$ | $S_3 : (1,1,1)$<br>$S_3 : (1,3,1)$ | ✗<br>✓ | ✓<br>✓ |

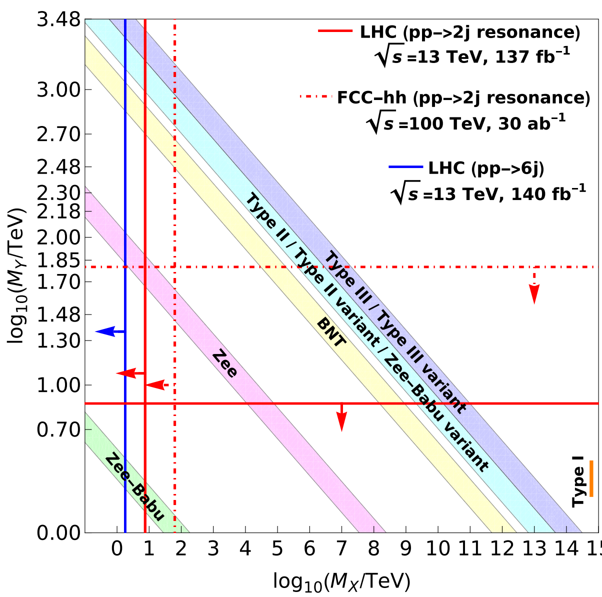



**Fig. 7.** $Y = S_1$ in all models. For BNT model, we fix the mass of $F_1$ to be equal to the mass of $S_2$. $X$ is to be inferred from Table 2. LHC/FCC-hh $pp \to 2j$ [60–62] ($pp \to 6j$ [63]) bound is applicable to di-quark scalars (color octet fermion). Each band corresponds to observable range of $n$–$\bar{n}$ process with parameter space to the left ruled out by SK, whereas the region to the right is beyond experimental sensitivity.

duced in Eq. (5), for all models under consideration. The corresponding experimental accessibility of $n$–$\bar{n}$ oscillation signal is shown in Fig. 7, where we implement the current bound from Super-Kamiokande [9] and future reach from NNBAR [6]. We set all Yukawa couplings to one and the relevant cubic scalar coupling to $\tilde{\mu} = M_{S_2}$.

Fig. 7 demonstrates that if $n$–$\bar{n}$ oscillation signal is observed, the Zee-Babu and the Type I seesaw models cannot be its source due to LHC bounds on dijet resonances. Also, the Type I, II, III, and BNT models may be in tension with proton decay limits unless the leptoquark couplings of the color triplet are extremely suppressed. For example, if one is to have a TeV scale color triplet with order-one diquark couplings, its lepton–quark Yukawa couplings should be $\lesssim 10^{-24}$ to sufficiently suppress associated proton decay signatures [67]. One possible approach is to abandon renormalization requirement altogether and consider higher-dimensional operators that allow for arbitrary suppression of the relevant color triplet Yukawa interactions, as discussed in Refs. [55,68]. The exact suppression implementation would still depend on the details of the model and is beyond the scope of this work. This is something one should keep in mind when interpreting the last column of Table 2. On the other hand, the variants of the Type II, III, and Zee-Babu models, as well as the Zee model itself are all viable with regard to proton decay constraints and may potentially yield collider-accessible signatures.

While Fig. 7 displays the most optimistic scenario with regard to observation of $n$–$\bar{n}$ oscillation signal, we clarify how the picture would change if the relevant Yukawa or cubic couplings are varied.

Let us first address the Type II scenario. The oscillation lifetime exhibits the following parameter dependence [2]

$$\tau_{\mathbf{A}} = \frac{20}{3\mathscr{M}} \frac{1}{y^{dd}(y^{ud})^2} \frac{M_{S_1}^4 M_{S_2}^2}{\tilde{\mu}}, \tag{10}$$

where **A** denotes the topology type and $\mathscr{M}$ is the relevant hadronic matrix element [69–73]. Consequently, achieving oscillation times within the observable range with smaller Yukawa coupling(s) would require lighter mediator state(s), implying that the observation bands shown in Fig. 7 would shift diagonally downward. A similar effect would occur if the cubic coupling $\tilde{\mu}$ is reduced. Note that this dimensionful coupling is constrained from above, i.e., $\tilde{\mu} \lesssim 3\times\max(M_{S_1}, M_{S_2})$, if one is to avoid charge-breaking minima [74–76]. We, in particular, use $\tilde{\mu} = M_{S_2}$ to generate Fig. 7.

For the Type III seesaw, the parameter dependence is [2]

$$\tau_{\mathbf{B}} = \frac{1}{\mathscr{M}} \frac{1}{(y^{dd})^2(y^{ud})^2} M_{F_1} M_{S_1}^4, \tag{11}$$

where **B** denotes the topology type. If one would reduce the relevant Yukawa couplings, the viable new physics mass scale would again need to go to lower values in order to have observable $n$–$\bar{n}$ oscillation signal. For instance, assuming unity for all Yukawa couplings and degenerate masses $M_{F_1} \approx M_{S_1} \approx 450$ TeV, the $n$–$\bar{n}$ transition rate is potentially measurable. If these couplings are reduced by an order of magnitude, the mass scales must decrease to $M_{F_1} \approx M_{S_1} \approx 70$ TeV to maintain observability. The experimentally accessible range is bounded from below by the Super-Kamiokande limit $\tau^{\text{SK}}_{n-\bar{n}} \geq 4.7\times10^8$ s [9], while the upper bound is set by the NNBAR experimental sensitivity $\tau^{\text{NNBAR}}_{n-\bar{n}} \leq 3\times10^9$ s [6].

The case of the Type I seesaw is slightly different since the mass $M_{F_1}$ not only enters $n$–$\bar{n}$ prediction but it is also directly connected to the neutrino mass scale. For instance, if all Yukawa couplings are reduced by a factor of 10, $M_{F_1}$ should need to be simultaneously lowered by a factor of $10^2$ in order to preserve the correct neutrino mass scale. With these changes, an observable $n$–$\bar{n}$ oscillation could still be achieved but only if $M_{S_1}$ is in the sub-TeV range. This, on the other hand, would be in conflict with the LHC data. The same issue arises if one is to reduce diquark couplings entering prediction for the $n$–$\bar{n}$ oscillation signal within the Type I seesaw scenario. In short, Fig. 7 depicts expectations for all possible scenarios considered if the relevant Yukawa couplings are set to one.

In the present work, all the couplings and mass parameters are evaluated at the $n$–$\bar{n}$ mediator scale set by max($M_{S_i}$) and/or max($M_{F_j}$). We expect renormalization-group (RG) effects on the tree level and radiative neutrino-mass generation sectors to induce at most moderate quantitative corrections without qualitative impact on the viable parameter regions shown in Fig. 7. For example, since the mediator masses in the Zee-Babu model lie in the multi-TeV range, the logarithmic running between the mediator and electroweak scales should be relatively modest. Also, our analysis is intended as a leading-order phenomenological study rather than a precision calculation. A complete effective field theory matching and RG evolution analysis for the scenarios we address here would become a necessity if and when $n$–$\bar{n}$ oscillation signal is observed.

Moreover, in plotting Fig. 7, we adopt the relevant hadronic matrix elements from Table 1 of Ref. [72]. Although for the $d = 9$ operators we do not explicitly include the RG running of the corresponding transition matrix elements, these effects are known to induce only moderate numerical corrections. In particular, Ref. [72] evolves the matrix elements from $\mu_0 = 2\,\text{GeV}$ to the relevant new-physics scale (even up to the GUT scale) to find that they remain relatively stable under RG evolution, preserving the same order of magnitude across the phenomenologically relevant mass range. Even the largest RG-induced variation over the full range from the GUT scale down to low energies corresponds only to an enhancement of the matrix element by a factor of a few (in particular, for $\mathcal{M}_5$, it is a factor of three enhancement) [72]. Therefore, for the mediator mass scales considered in this work, we do not expect the omission of these running effects to qualitatively alter the viable parameter regions shown in Fig. 7.

## 4. Conclusions

We have demonstrated that Majorana neutrino mass generation and $n$–$\bar{n}$ oscillation are intrinsically linked if $SU(5)$ is realized in nature. Our study explores the simplest extensions of the GG model that yield realistic neutrino sector and, in doing so, automatically induce $n$–$\bar{n}$ oscillation. This implies the existence of additional baryon-number-violating processes beyond the ever-present proton decay mediated by gauge bosons and scalars. Majorana masses can be probed via $|\Delta L| = 2$ processes, such as neutrinoless double beta decay and lepton flavor–violating transitions like muon to positron conversion. Similarly, $|\Delta B| = 2$ processes, including $n$–$\bar{n}$ oscillation and di-nucleon decays, offer complementary experimental avenues, motivating dedicated searches for nucleon decay and neutron–antineutron oscillations.

We summarize our crucial findings here:

- Within the economical models considered in this work, the mediators responsible for neutron–antineutron oscillations can be either scalar or fermionic states. The scalar mediators include color-triplet and color-sextet scalars, while the beyond SM fermions can be singlets, color sextets, or color octets.
- Models that contain at least one scalar color triplet $(3, 1, -1/3)$ mediator cannot lead to observable $n$–$\bar{n}$ oscillations because of the associated proton decay constraints. Consequently, an observation of $n$–$\bar{n}$ oscillations in future experiments would strongly disfavor the minimal versions of Type I, Type II, and Type III seesaw models, as well as the BNT model.
- The Zee-Babu model cannot generate observable $n$–$\bar{n}$ oscillations since it requires scalar diquarks — a color triplet $(\bar{3}, 1, -2/3)$ and a color sextet — in a mass range already excluded by LHC searches. This exclusion originates from the antisymmetric Yukawa couplings of the color-triplet scalar, where two such mediators are needed to generate $n$–$\bar{n}$ process.
- Models compatible with observable $n$–$\bar{n}$ oscillations include the minimally extended Type II and Type III seesaw models, the extended Zee-Babu model, and the minimal Zee model. In the Type II and Zee-Babu scenarios, the relevant mediators are scalar color sextets that do not induce proton decay. In the Zee model, one mediator is the color-triplet scalar $(\bar{3}, 1, -2/3)$, which is likewise free from proton decay constraints. In the extended Type III scenario, the oscillation is mediated by a color-sextet scalar together with a color-octet fermion, again avoiding proton decay bounds.
- If future experiments observe $n$–$\bar{n}$ oscillations, the Zee model discussed here will be fully testable at the `FCC-hh` through searches for low-mass diquarks. In contrast, even in the absence of such a discovery at the `FCC-hh`, viable parameter space for observable $n$–$\bar{n}$ oscillation would still remain for the Type II and Type III seesaw model variants, as well as the Zee-Babu model variant.

Our dedicated study, therefore, provides clear guidance for model building in the exciting event of an observed $n$–$\bar{n}$ oscillation signal.

## Data availability

No data was used for the research described in the article.

## Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

## Acknowledgments

I.D. acknowledges the financial support from the Slovenian Research Agency (research core funding No. P1-0035) and hospitality of the CERN Theory Department, where part of this work was completed. S.F. and S.S. acknowledge the financial support from the Slovenian Research Agency (research core funding No. P1-0035 and N1-0321).

## References

[1] H. Georgi, S.L. Glashow, Unity of all elementary particle forces, Phys. Rev. Lett. 32 (1974) 438–441. https://doi.org/10.1103/PhysRevLett.32.438
[2] I. Dorsner, S. Fajfer, S. Saad, Beyond neutrino mass: observable $n$-$\bar{n}$ oscillations in UV complete seesaw models (2025). arXiv:2509.00145 [hep-ph]
[3] V.A. Kuzmin, Cp violation and baryon asymmetry of the universe, Pisma Zh. Eksp. Teor. Fiz. 12 (1970) 335–337.
[4] R.N. Mohapatra, R.E. Marshak, Local B-L symmetry of electroweak interactions, majorana neutrinos and neutron oscillations, Phys. Rev. Lett. 44 (1980) 1316–1319. [Erratum: Phys. Rev. Lett. 44 (1980) 1643]. https://doi.org/10.1103/PhysRevLett.44.1316
[5] I.I. Phillips, D. G., et al., Neutron-antineutron oscillations: theoretical status and experimental prospects, Phys. Rept. 612 (2016) 1–45. arXiv:1410.1100 [hep-ex], https://doi.org/10.1016/j.physrep.2015.11.001
[6] A. Addazi, et al., New high-sensitivity searches for neutrons converting into antineutrons and/or sterile neutrons at the HIBEAM/NNBAR experiment at the European spallation source, J. Phys. G 48 (7) (2021) 070501. arXiv:2006.04907 [physics.ins-det], https://doi.org/10.1088/1361-6471/abf429
[7] S. Rao, R.E. Shrock, Six fermion $(B-L)$ violating operators of arbitrary generational structure, Nucl. Phys. B 232 (1984) 143–179. https://doi.org/10.1016/0550-3213(84)90365-1
[8] M.J. Dolinski, A.W.P. Poon, W. Rodejohann, Neutrinoless double-beta decay: status and prospects, Ann. Rev. Nucl. Part. Sci. 69 (2019) 219–251. arXiv:1902.04097 [nucl-ex] https://doi.org/10.1146/annurev-nucl-101918-023407
[9] Super-Kamiokande Collaboration, K. Abe, et al., Neutron-antineutron oscillation search using a 0.37 megaton-years exposure of Super-Kamiokande, Phys. Rev. D 103 (1) (2021) 012008. arXiv:2012.02607 [hep-ex], https://doi.org/10.1103/PhysRevD.103.012008
[10] DUNE Collaboration, R. Acciarri, et al., Long-Baseline Neutrino Facility (LBNF) and Deep Underground Neutrino Experiment (DUNE): Conceptual Design Report, Volume 2: The Physics Program for DUNE at LBNF (2015). arXiv:1512.06148 [physics.ins-det]
[11] M. Lee, M. MacKenzie, Muon to positron conversion, Universe 8 (2022) 227. arXiv:2110.07093 [hep-ex], https://doi.org/10.3390/universe8040227
[12] K. Aitken, D. McKeen, T. Neder, A.E. Nelson, Baryogenesis from oscillations of charmed or beautiful baryons, Phys. Rev. D 96 (7) (2017) 075009. arXiv:1708.01259 [hep-ph], https://doi.org/10.1103/PhysRevD.96.075009
[13] S. Girmohanta, R. Shrock, Improved upper limits on baryon-Number violating dinucleon decays to dileptons, Phys. Lett. B 803 (2020) 135296. arXiv:1910.08356 [hep-ph], https://doi.org/10.1016/j.physletb.2020.135296

[14] P.S.B. Dev, et al., Searches for baryon number violation in neutrino experiments: a white paper, J. Phys. G 51 (3) (2024) 033001. arXiv:2203.08771 [hep-ex], https://doi.org/10.1088/1361-6471/ad1658
[15] A.B. Beneito, I, S. Fajfer, A.A. Petrov, New Avenues for $|\Delta B| = 2$ Processes Beyond Neutron-Antineutron Oscillations, (2025). arXiv:2511.05657 [hep-ph]
[16] S. Weinberg, Baryon and lepton nonconserving processes, Phys. Rev. Lett. 43 (1979) 1566–1570. https://doi.org/10.1103/PhysRevLett.43.1566
[17] S. Weinberg, Varieties of baryon and lepton nonconservation, Phys. Rev. D 22 (1980) 1694. https://doi.org/10.1103/PhysRevD.22.1694
[18] F. Wilczek, A. Zee, Conservation or violation of $B^{-}$l in proton decay, Phys. Lett. B 88 (1979) 311–314. https://doi.org/10.1016/0370-2693(79)90475-1
[19] F. Wilczek, A. Zee, Operator analysis of nucleon decay, Phys. Rev. Lett. 43 (1979) 1571–1573. https://doi.org/10.1103/PhysRevLett.43.1571
[20] H.A. Weldon, A. Zee, Operator analysis of new physics, Nucl. Phys. B 173 (1980) 269–290. https://doi.org/10.1016/0550-3213(80)90218-7
[21] H. Georgi, C. Jarlskog, A new lepton - Quark mass relation in a unified theory, Phys. Lett. B 86 (1979) 297–300. https://doi.org/10.1016/0370-2693(79)90842-6
[22] A. de Gouvea, J. Herrero-Garcia, A. Kobach, Neutrino masses, grand unification, and baryon number violation, Phys. Rev. D 90 (1) (2014) 016011. arXiv:1404.4057 [hep-ph], https://doi.org/10.1103/PhysRevD.90.016011
[23] P. Minkowski, $\mu \to e\gamma$ At a rate of one out of $10^9$ muon decays?, Phys. Lett. 67B (1977) 421–428. https://doi.org/10.1016/0370-2693(77)90435-X
[24] T. Yanagida, Horizontal gauge symmetry and masses of neutrinos, Conf. Proc. C7902131 (1979) 95–99.
[25] S.L. Glashow, The future of elementary particle physics, NATO Sci. Ser. B 61 (1980) 687. https://doi.org/10.1007/978-1-4684-7197-7_15
[26] M. Gell-Mann, P. Ramond, R. Slansky, Complex spinors and unified theories, Conf. Proc. C 790927 (1979) 315–321. arXiv:1306.4669 [hep-th]
[27] R.N. Mohapatra, G. Senjanovic, Neutrino mass and spontaneous parity nonconservation, Phys. Rev. Lett. 44 (1980) 912. https://doi.org/10.1103/PhysRevLett.44.912
[28] M. Magg, C. Wetterich, Neutrino mass problem and gauge hierarchy, Phys. Lett. B 94 (1980) 61–64. https://doi.org/10.1016/0370-2693(80)90825-4
[29] J. Schechter, J.W.F. Valle, Neutrino masses in SU(2) x U(1) theories, Phys. Rev. D22 (1980) 2227. https://doi.org/10.1103/PhysRevD.22.2227
[30] G. Lazarides, Q. Shafi, C. Wetterich, Proton lifetime and fermion masses in an SO(10) model, Nucl. Phys. B 181 (1981) 287–300. https://doi.org/10.1016/0550-3213(81)90354-0
[31] R.N. Mohapatra, G. Senjanovic, Neutrino masses and mixings in gauge models with spontaneous parity violation, Phys. Rev. D 23 (1981) 165. https://doi.org/10.1103/PhysRevD.23.165
[32] E. Ma, Pathways to naturally small neutrino masses, Phys. Rev. Lett. 81 (1998) 1171–1174. arXiv:9805219 [hep-ph], https://doi.org/10.1103/PhysRevLett.81.1171
[33] I. Dorsner, P. Fileviez Perez, Unification without supersymmetry: neutrino mass, proton decay and light leptoquarks, Nucl. Phys. B723 (2005) 53–76. arXiv:0504276 [hep-ph], https://doi.org/10.1016/j.nuclphysb.2005.06.016
[34] I. Dorsner, P. Fileviez Perez, R. Gonzalez Felipe, Phenomenological and cosmological aspects of a minimal GUT scenario, Nucl. Phys. B 747 (2006) 312–327. arXiv:0512068 [hep-ph], https://doi.org/10.1016/j.nuclphysb.2006.05.006
[35] I. Dorsner, P. Fileviez Perez, G. Rodrigo, Fermion masses and the UV cutoff of the minimal realistic SU(5), Phys. Rev. D 75 (2007) 125007. arXiv:0607208 [hep-ph], https://doi.org/10.1103/PhysRevD.75.125007
[36] I. Dorsner, I. Mocioiu, Predictions from type II see-saw mechanism in SU(5), Nucl. Phys. B 796 (2008) 123–136. arXiv:0708.3332 [hep-ph], https://doi.org/10.1016/j.nuclphysb.2007.12.004
[37] S. Antusch, K. Hinze, S. Saad, Viable quark-lepton Yukawa ratios and nucleon decay predictions in SU(5) GUTs with type-II seesaw, Nucl. Phys. B 986 (2023) 116049. arXiv:2205.01120 [hep-ph], https://doi.org/10.1016/j.nuclphysb.2022.116049
[38] L. Calibbi, X. Gao, Lepton flavor violation in minimal grand unified type II seesaw models, Phys. Rev. D 106 (9) (2022) 095036. arXiv:2206.10682 [hep-ph], https://doi.org/10.1103/PhysRevD.106.095036
[39] S. Antusch, K. Hinze, S. Saad, Minimal SU(5) GUTs with vectorlike fermions, Phys. Rev. D 108 (9) (2023) 095010. arXiv:2308.08585 [hep-ph], https://doi.org/10.1103/PhysRevD.108.095010
[40] A. Kaladharan, S. Saad, Unified origin of inflation, baryon asymmetry, and neutrino mass, Phys. Rev. D 110 (11) (2024) 116012. arXiv:2409.02225 [hep-ph], https://doi.org/10.1103/PhysRevD.110.116012
[41] P.S.B. Dev, S. Goswami, D. Pachhar, S.K. Shukla, Scalar-induced Neutrinoless Double Beta Decay in $SU(5)$ (2025). arXiv:2507.16606 [hep-ph]
[42] R. Foot, H. Lew, X.G. He, G.C. Joshi, Seesaw neutrino masses induced by a triplet of leptons, Z. Phys. C 44 (1989) 441. https://doi.org/10.1007/BF01415558
[43] B. Bajc, G. Senjanovic, Seesaw at LHC, JHEP 08 (2007) 014. arXiv:hep-ph/0612029 [hep-ph], https://doi.org/10.1088/1126-6708/2007/08/014
[44] I. Dorsner, P. Fileviez Perez, Upper bound on the mass of the type III seesaw triplet in an SU(5) model, JHEP 06 (2007) 029. arXiv:0612216 [hep-ph], https://doi.org/10.1088/1126-6708/2007/06/029
[45] P. Fileviez Perez, Renormalizable adjoint SU(5), Phys. Lett. B 654 (2007) 189–193. arXiv:0702287 [hep-ph], https://doi.org/10.1016/j.physletb.2007.07.075
[46] S. Antusch, K. Hinze, S. Saad, Quark-lepton Yukawa ratios and nucleon decay in SU(5) GUTs with type-III seesaw, Nucl. Phys. B 991 (2023) 116195. arXiv:2301.03601 [hep-ph], https://doi.org/10.1016/j.nuclphysb.2023.116195
[47] K.S. Babu, S. Nandi, Z. Tavartkiladze, New mechanism for neutrino mass generation and triply charged higgs bosons at the LHC, Phys. Rev. D80 (2009) 071702. arXiv:0905.2710 [hep-ph], https://doi.org/10.1103/PhysRevD.80.071702
[48] I. Doršner, S. Saad, Towards minimal $SU(5)$, Phys. Rev. D 101 (1) (2020) 015009. arXiv:1910.09008 [hep-ph], https://doi.org/10.1103/PhysRevD.101.015009
[49] I. Doršner, E. Džaferović-Mašić, S. Saad, Parameter space exploration of the minimal SU(5) unification, Phys. Rev. D 104 (1) (2021) 015023. arXiv:2105.01678 [hep-ph], https://doi.org/10.1103/PhysRevD.104.015023
[50] S. Antusch, I. Doršner, K. Hinze, S. Saad, Fully testable axion dark matter within a minimal SU(5) GUT, Phys. Rev. D 108 (1) (2023) 015025. arXiv:2301.00809 [hep-ph], https://doi.org/10.1103/PhysRevD.108.015025
[51] I. Doršner, E. Džaferović-Mašić, S. Fajfer, S. Saad, Gauge and scalar boson mediated proton decay in a predictive SU(5) GUT model, Phys. Rev. D 109 (7) (2024) 075023. arXiv:2401.16907 [hep-ph], https://doi.org/10.1103/PhysRevD.109.075023
[52] A. Zee, A theory of lepton number violation, neutrino majorana mass, and oscillation, Phys. Lett. B 93 (1980) 389. [Erratum: Phys.Lett.B 95, 461 (1980)]. https://doi.org/10.1016/0370-2693(80)90349-4
[53] P. Fileviez Perez, C. Murgui, Renormalizable SU(5) unification, Phys. Rev. D94 (7) (2016) 075014. arXiv:1604.03377 [hep-ph], https://doi.org/10.1103/PhysRevD.94.075014
[54] C. Klein, M. Lindner, S. Vogl, Radiative neutrino masses and successful $SU(5)$ unification, Phys. Rev. D100 (7) (2019) 075024. arXiv:1907.05328 [hep-ph], https://doi.org/10.1103/PhysRevD.100.075024
[55] I. Doršner, M. Matković, S. Saad, Nonrenormalizable SU(5) GUTs: leptoquark-induced neutrino masses, Phys. Rev. D 111 (11) (2025) 115039. arXiv:2504.16022 [hep-ph], https://doi.org/10.1103/3plx-vwrb
[56] T.P. Cheng, L.-F. Li, Neutrino masses, mixings and oscillations in SU(2) x u(1) models of electroweak interactions, Phys. Rev. D 22 (1980) 2860. https://doi.org/10.1103/PhysRevD.22.2860
[57] K.S. Babu, Model of 'Calculable' majorana neutrino masses, Phys. Lett. B 203 (1988) 132–136. https://doi.org/10.1016/0370-2693(88)91584-5
[58] K.S. Babu, S. Saad, Ultraviolet completion of a two-loop neutrino mass model, JHEP 03 (2025) 132. arXiv:2412.14562 [hep-ph], https://doi.org/10.1007/JHEP03(2025)132
[59] S. Saad, Origin of a two-loop neutrino mass from SU(5) grand unification, Phys. Rev. D99 (11) (2019) 115016. arXiv:1902.11254 [hep-ph], https://doi.org/10.1103/PhysRevD.99.115016
[60] A.M. Sirunyan, et al., CMS Collaboration, Search for high mass dijet resonances with a new background prediction method in proton-proton collisions at $\sqrt{s} = 13$ TeV, JHEP 05 (2020) 033. arXiv:1911.03947 [hep-ex], https://doi.org/10.1007/JHEP05(2020)033
[61] R.M. Harris, E.G. Guler, Y. Guler, Sensitivity to dijet resonances at proton-proton colliders, in: Snowmass 2021, 2022. arXiv:2202.03389 [hep-ex]
[62] G. Bernardi, et al., The Future Circular Collider: a Summary for the US 2021 Snowmass Process (2022). arXiv:2203.06520 [hep-ex]
[63] ATLAS Collaboration, G. Aad, et al., A search for R-parity-violating supersymmetry in final states containing many jets in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, JHEP 05 (2024) 003. arXiv:2401.16333 [hep-ex], https://doi.org/10.1007/JHEP05(2024)003
[64] I. Esteban, M.C. Gonzalez-Garcia, M. Maltoni, I. Martinez-Soler, J.a.P. Pinheiro, T. Schwetz, NuFit-6.0: updated global analysis of three-flavor neutrino oscillations, JHEP 12 (2024) 216. arXiv:2410.05380 [hep-ph], https://doi.org/10.1007/JHEP12(2024)216
[65] KamLAND-Zen Collaboration, S. Abe, et al., Search for Majorana Neutrinos with the Complete KamLAND-Zen Dataset (2024). arXiv:2406.11438 [hep-ex]
[66] nEXO Collaboration, G. Adhikari, et al., nEXO: neutrinoless double beta decay search beyond $10^{28}$ year half-life sensitivity, J. Phys. G 49 (1) (2022) 015104. arXiv:2106.16243 [nucl-ex] https://doi.org/10.1088/1361-6471/ac3631
[67] I. Dorsner, A scalar leptoquark in SU(5), Phys. Rev. D86 (2012) 055009. arXiv:1206.5998 [hep-ph], https://doi.org/10.1103/PhysRevD.86.055009
[68] I. Doršner, S. Saad, Is doublet-triplet splitting necessary?, Phys. Rev. D 110 (7) (2024) 075025. arXiv:2404.09021 [hep-ph], https://doi.org/10.1103/PhysRevD.110.075025
[69] M.I. Buchoff, M. Wagman, Perturbative renormalization of neutron-antineutron operators, Phys. Rev. D 93 (1) (2016) 016005. [Erratum: Phys. Rev. D 98 (2018) 079901]. arXiv:1506.00647 [hep-ph], https://doi.org/10.1103/PhysRevD.93.016005
[70] E. Rinaldi, S. Syritsyn, M.L. Wagman, M.I. Buchoff, C. Schroeder, J. Wasem, Neutron-antineutron oscillations from lattice QCD, Phys. Rev. Lett. 122 (16) (2019) 162001. arXiv:1809.00246 [hep-lat], https://doi.org/10.1103/PhysRevLett.122.162001
[71] E. Rinaldi, S. Syritsyn, M.L. Wagman, M.I. Buchoff, C. Schroeder, J. Wasem, Lattice QCD determination of neutron-antineutron matrix elements with physical quark masses, Phys. Rev. D 99 (7) (2019) 074510. arXiv:1901.07519 [hep-lat], https://doi.org/10.1103/PhysRevD.99.074510
[72] K. Fridell, J. Harz, C. Hati, Probing baryogenesis with neutron-antineutron oscillations, JHEP 11 (2021) 185. arXiv:2105.06487 [hep-ph], https://doi.org/10.1007/JHEP11(2021)185
[73] M. Thomas Arun, S. M, R. Pal, RG evolution and effect of intermediate new physics on $\Delta B = 2$ six-quark operators, JHEP 10 (2025) 032. arXiv:2506.10105 [hep-ph], https://doi.org/10.1007/JHEP10(2025)032
[74] J.M. Frere, D.R.T. Jones, S. Raby, Fermion masses and induction of the weak scale by supergravity, Nucl. Phys. B 222 (1983) 11–19. https://doi.org/10.1016/0550-3213(83)90606-5
[75] L. Alvarez-Gaume, J. Polchinski, M.B. Wise, Minimal low-energy supergravity, Nucl. Phys. B 221 (1983) 495. https://doi.org/10.1016/0550-3213(83)90591-6
[76] J.A. Casas, S. Dimopoulos, Stability bounds on flavor violating trilinear soft terms in the MSSM, Phys. Lett. B 387 (1996) 107–112. arXiv:9606237 [hep-ph], https://doi.org/10.1016/0370-2693(96)01000-3