

TeV window to grand unification: The Higgs boson's light color triplet partner

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The color-triplet partner of the Higgs doublet, called a T -particle, is a universal feature of grand unification. It was shown some time ago that this particle can be accessible for direct production in collider experiments. In this paper, we point out that the T -particle represents a simultaneous low-energy probe of baryon number violation as well as of the origin of the neutrino mass, linking the mediation of proton decay with oscillations of the neutron into a sterile neutrino. We point out a triple correlation between its collider signatures, proton decay measurements, and the searches for the magnetic resonance disappearance of free neutrons in cold neutron experiments. In this way, the T -particle can provide a diversity of correlated experimental windows into grand unification.

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I. INTRODUCTION

The goal of the present paper is to show that a light color-triplet partner of the Higgs doublet, predicted in Refs. [1,2], can serve as a low-energy window to grand unification via correlated signatures in different categories of experiment.

Grand unified theories (GUTs) provide a symmetry-based framework for the idea of the unification of forces. The minimal structure is provided by the $SU(5)$ model of Georgi and Glashow [3] which embeds the $G_{\text{SM}} = SU(3)_c \times SU(2)_w \times U(1)_Y$ symmetry of the Standard Model (SM) in its compact covering group $SU(5)$. This gauge symmetry is Higgsed (“broken”) down to the SM group at the GUT scale, M_G . The precise value of this scale depends on the low-energy particle content of the theory as well as on the threshold corrections from the physics at the scale M_G . Already in the minimal case, the running of the coupling constants shows that the GUT scale is typically very high, $M_G \sim 10^{15-16}$ GeV [4]. This is also in accordance with the phenomenological bounds on proton decay

mediated by the X and Y gauge bosons corresponding to the broken generators from $SU(5)/G_{\text{SM}}$.

The universal feature of grand unification is the presence of the color-triplet partner of the SM Higgs doublet, H . We shall denote this particle by T . Because of the GUT symmetry, the T -particle is forced to participate in the $SU(5)$ -invariant Yukawa couplings with quarks and leptons along with H . Correspondingly, in the limit of an unbroken $SU(5)$ -symmetry, it mediates processes that violate baryon and lepton numbers. A particularly important process is proton decay which is subject to very stringent experimental bounds, with a lifetime longer than $\sim 10^{34}$ yr.

To accommodate such bounds, ordinarily, the T -particle is assumed to have a mass $\sim M_G$. This goes in contrast with the mass of its Higgs partner which is around ~ 100 GeV. This mass difference is traditionally referred to as the “doublet-triplet splitting” problem.

However, this is only a naturalness issue, rather than a consistency one: since the GUT symmetry is broken, the mass splitting can be easily arranged via fine-tuning. The extensions of the theory can allow for more natural solutions either based on symmetry (e.g., within the context of supersymmetric extensions of GUTs) or a cosmological relaxation of the Higgs mass [5]. In this work, we shall not enter into the naturalness issues of the mass-splitting scenarios which generate a large mass for the T -particle.

From an observational point of view, a universal drawback of such scenarios is that they eliminate any

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chance of direct detection of the T -particle in collider experiments. However, an alternative realization of the theory exists in which such detection is possible. The idea, originally proposed in Ref. [1] and further refined in Refs. [2,6] and subsequent papers, is that, instead of the masses, the splitting can take place among the Yukawa couplings of H and T particles. As shown in this work, the breaking of GUT symmetry allows for a very strong relative suppression of the Yukawa couplings of the T -boson with ordinary quarks and leptons. Thus, in the low-energy effective theory, the Yukawa couplings of T and H must be treated as independent parameters. This is true already at the level of the $SU(5)$ GUT [6]. This mechanism eliminates the necessity of a large mass splitting, thereby liberating the T -boson to be nearly degenerate with the SM Higgs.

As a result, in this scenario, the T -boson emerges as a light particle, with mass $m_T \sim \text{TeV}$, but with suppressed baryon and lepton number violating couplings. This allows for the T -particle to be potentially accessible at colliders, without conflicting with other phenomenological bounds.

This setup opens up the possibility of a direct detection of GUT's most universal feature in particle collisions as well as the correlation of its collider signatures with other observations.

In the present paper, we point out some new implications for baryon and lepton number violating processes involving the neutrino sector. In particular, we point out a triple correlation between the signatures in the following three classes of experiments: 1) a direct production of the T -particle at high-energy accelerators; 2) the processes mediated by the T -boson in proton decay measurements; and 3) the resonant transitions of a free neutron into the sterile neutrino, induced by the T -particle, potentially accessible in experiments with cold neutrons.

Before proceeding, we would like to refer to two parallel complementary papers exploring implications of the T -particle for flavor physics [7] and for cosmology [8]. We would also like to point to various other implications of the T -particle with split couplings [1,2,6,9–19].

II. SETUP

For simplicity, we shall work in $SU(5)$ theory with minimal field content and a single generation of fermions. In $SU(5)$, the SM fermions are arranged in a $\bar{5}$ -dimensional and a 10-dimensional representation in the following way:

$$\bar{5}_F = (d_r^c \quad d_g^c \quad d_b^c \quad e^- \quad -\nu) \quad (1)$$

and

$$10_F = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & u_b^c & -u_g^c & -u_r & -d_r \\ -u_b^c & 0 & u_r^c & -u_g & -d_g \\ u_g^c & -u_r^c & 0 & -u_b & -d_b \\ u_r & u_g & u_b & 0 & e^+ \\ d_r & d_g & d_b & -e^+ & 0 \end{pmatrix}. \quad (2)$$

Here, the superscript c stands for denoting antiparticles, and r, g, b are the three colors of QCD. All fermions are written in the left-handed basis.

In addition, for generating the mass of the neutrino, we introduce an $SU(5)$ -singlet fermion ν^c , also written in the left-handed basis. In the left-right basis, this would play the role of a right-handed (sterile) partner of the active left-handed neutrino ν .

The SM Higgs is embedded in a fundamental representation,

$$5_H = \begin{pmatrix} T_r \\ T_g \\ T_b \\ H^+ \\ H^0 \end{pmatrix}, \quad (3)$$

where the SM Higgs doublet H sits in the lower two components and a new triplet T in the upper three.

As usual, the breaking of the $SU(5)$ -symmetry down to the SM group is achieved by a Higgs scalar in the adjoint representation which we denote by 24_H . This field acquires a vacuum expectation value (VEV),

$$\langle 24_H \rangle \propto M_G \text{diag} \left(1, 1, 1, -\frac{3}{2}, -\frac{3}{2} \right). \quad (4)$$

This VEV reduces the symmetry of the vacuum to G_{SM} . The second stage of the symmetry breaking is achieved by the VEV of the Higgs doublet $\langle H^0 \rangle = v$, living in the 5_H -plet, $\langle 5_H \rangle = (0 \ 0 \ 0 \ 0 \ v)$.

This field generates masses of fermions via the $SU(5)$ -invariant Yukawa couplings. The minimal set of renormalizable couplings that can provide masses to all the fermions, including neutrinos, is given by the structure

$$\begin{aligned} \mathcal{L}_{\text{fermi}} = & Y_d 5_H^* 10_F \bar{5}_F + Y_u 10_F 10_F 5_H \\ & + Y_\nu 5_H \bar{5}_F \nu^c + m_R \nu^c \nu^c, \end{aligned} \quad (5)$$

where Y_d, Y_u, Y_ν are the Yukawa couplings. The structure of $SU(5)$ and Lorentz indices is obvious and is not shown explicitly. The first row is responsible for the masses of the charged fermions, whereas the second row generates the mass of the neutrino. Depending on the values of m_R and Y_ν , the mass of ν can be either Dirac or Majorana type. In particular, for $m_R \gg Y_\nu v$, the Majorana mass of the active

neutrino, $m_\nu \sim (Y_\nu v)^2/m_R$, is generated via the seesaw mechanism.

Decomposing the above operators in terms of SM particles, we immediately discover that the couplings of T and H are related. In such a case, the mass of T must be very large in order not to conflict with the experimental bounds on baryon and lepton number violating processes. This is the logic behind the standard assumption of the doublet-triplet mass splitting.

However, notice that the masses of ordinary fermions generated from the above Yukawa couplings are not phenomenologically viable. In particular, there is a wrong prediction of equal masses for down quarks and electrons [20]. This rules out the simplest theory.

To obtain viable masses, one necessarily needs to go beyond the minimal structure by adding higher-dimensional operators [20]. These operators must involve the 24_H -Higgs. This is necessary for communicating the message about the breaking of the $SU(5)$ -symmetry to the Yukawa sector. In a similar spirit, the effective non-renormalizable operators can affect the unification of the gauge couplings [21] (for a recent analysis in this direction, see Ref. [22]).

Now, the idea of Refs. [1,2,6] is that the high-dimensional operators with 24_H necessarily undo the $SU(5)$ -relations among the couplings of H and T . As a result, in the low-energy theory these couplings appear as independent parameters. In particular, the couplings of the color-triplet T with the fermions can be arbitrarily suppressed. This liberates the mass of the T -particle, allowing it to be within the range of direct experimental accessibility at colliders.

Obviously, the above mechanism is operative already at the level of minimal $SU(5)$ [6]. For example, to the lowest order in 24_H , there exist the following operators that correct the structure of the Yukawa couplings,

$$\begin{aligned} \mathcal{L}_{5\text{-dim}} \supset & \frac{g_1}{\Lambda} 5_H^* 10_F 24_H \bar{5}_F + \frac{g_2}{\Lambda} 5_H^* 24_H 10_F \bar{5}_F \\ & + \frac{g_3}{\Lambda} 10_F 10_F 24_H 5_H + \frac{g_4}{\Lambda} 24_H 10_F 10_F 5_H \\ & + \frac{g_5}{\Lambda} 5_H 24_H \bar{5}_F \nu^c, \end{aligned} \quad (6)$$

where Λ is a cutoff scale and the g_i -s ($i = 1, 2, \dots, 5$) are numerical coefficients. It is clear that upon inserting the VEV of the 24_H (4) in the above operators, we get the effective low-energy theory in which the $SU(5)$ relations among the Yukawa couplings are completely shattered.

For example, already the addition of the g_2 -operator splits the $SU(5)$ invariant coupling $5_H^* 10_F \bar{5}_F$ between T and H as $Y_d + g_2 \frac{M_G}{\Lambda}$ and $Y_d - g_2 \frac{3M_G}{2\Lambda}$, respectively. Other couplings undergo similar splittings. This shows that the effective coupling of the T -particle can suppress proton decay to the observable level while accommodating the required Yukawa couplings of the Higgs.

In general, one has to include all possible higher-order operators. Their precise structure depends on the cutoff Λ and the specifics of the fundamental theory operating above it. However, there are some general consistency arguments indicating that the cutoff cannot be arbitrarily far from M_G .

In particular, in any theory propagating N particle species, gravity imposes a universal nonperturbative upper bound on the cutoff scale given by

$$\Lambda < \frac{M_P}{\sqrt{N}}, \quad (7)$$

where M_P is the Planck mass [23,24].¹ Notice that already within the SM $N \sim 100$, and in GUTs, it is naturally even higher. This brings the upper bound on Λ at least an order of magnitude below the Planck scale. This fact is important as it tells us that, already from pure gravity, the corrections to the low-energy couplings cannot be ignored.

In practice, Λ can be even lower than the species scale and not far from the GUT scale. In particular, the operators of the type (6) can be obtained from renormalizable interactions after integrating out multiplets with masses around M_G . The explicit realizations of such scenarios can be found in Refs. [1,2]. However, specifying the structure of the high-energy theory is beyond the scope of the present paper. Instead, we shall treat the Yukawa couplings of the T -particle as free parameters of the low-energy theory and study certain interesting ranges of the parameter space.

Following Ref. [6], with all generality, we parametrize the interactions of the T -particle and the ordinary Higgs doublet, H , with the SM fermions by the following effective Lagrangian:

$$\begin{aligned} \mathcal{L} \supset & g_{Tud} T^* u^c d^c + g_{TQL} T^* Q L \\ & + g_{Tue} T u^c e^c + g_{TQQ} T Q Q + g_{Tdv} T d^c \nu^c \\ & + Y'_d H^* Q d^c + Y'_e H^* L e^c + Y'_u H Q u^c + Y'_\nu H L \nu^c. \end{aligned} \quad (8)$$

In the limit of an exact $SU(5)$ -symmetry, we would have $g_{Tud} = g_{TQL} = Y'_d = Y'_e$, $g_{Tue} = g_{TQQ} = Y'_u$, and $g_{Tdv} = Y'_\nu$. However, through high-dimensional operators of the type (6), the VEV of 24_H splits them arbitrarily.

The couplings of the doublet are fixed by the Higgs VEV (v) and the masses of corresponding fermions in the usual way: $Y'_d = m_d/v$, $Y'_e = m_e/v$, $Y'_u = m_u/v$. The analogous relation for the neutrino coupling/mass depends on the value of m_R and will be discussed below.

¹This bound originates from the consistency of black hole physics and is fully nonperturbative in nature. It is therefore exact and cannot be removed by any resummation of the series. The perturbative arguments supporting it can also be given (e.g., see Refs. [25,26]).

The couplings of the triplet shall be constrained by various processes with baryon and lepton number violation.

III. T -PHENOMENOLOGY

A. Proton and neutron decays

The exchange of the T -particle can mediate baryon and lepton number violating processes. Let us first consider the processes of baryon (proton or neutron) decay. Such decays require an exchange that on one hand converts a quark into a quark and on the other hand a quark into a lepton. Correspondingly, the amplitudes of such processes are governed by the different cross products of the couplings of the type T -quark-quark (g_{Tud}, g_{TQQ}) and T -quark-lepton ($g_{Tue}, g_{TQL}, g_{Td\bar{\nu}}$).

For a clearer visualization, it is useful to write down the set of corresponding baryon and lepton number violating effective four-fermion operators obtained by the tree-level exchanges of the T -particle,

$$\begin{aligned} \mathcal{L}_{\text{eff}} \supset & \frac{g_{TQQ}g_{Tue}}{m_T^2} QQ u^{c*} e^{c*} + \frac{g_{Tud}g_{Tue}}{m_T^2} u^c d^c u^c e^c \\ & + \frac{g_{Tud}g_{Td\bar{\nu}}}{m_T^2} u^c d^c d^c \nu^c + \frac{g_{TQQ}g_{TQL}}{m_T^2} QQ QL \\ & + \frac{g_{Tud}g_{TQL}}{m_T^2} u^c d^c Q^* L^* + \frac{g_{TQQ}g_{Td\bar{\nu}}}{m_T^2} QQ d^{c*} \nu^{c*}, \quad (9) \end{aligned}$$

where m_T is the mass of the T -particle and the star stands for conjugation. These operators can mediate proton and neutron decay in various channels. For example, consider a proton decay via the channel $p \rightarrow e^+ \pi^0$. A representative exchange diagram is given in Fig. 1.

Taking the mass of the triplet around its current collider bound, $m_T \sim \text{TeV}$, the experimental bound on the proton lifetime, $\tau_p > 10^{34}$ yr [27], implies the following constraint on the cross products of T -quark-quark and T -quark-lepton couplings,

$$(\text{T-quark-quark}) \cdot (\text{T-quark-Lepton}) < 10^{-26}, \quad (10)$$

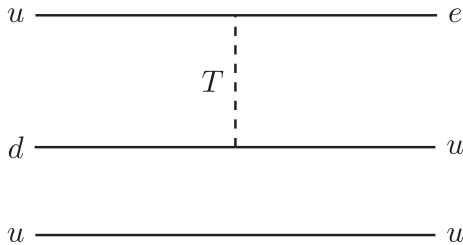


FIG. 1. An example of a proton decay process mediated via the T -particle.

where the product runs over $(g_{Tud} \cdot g_{Tue}), (g_{TQQ} \cdot g_{TQL}), (g_{Tud} \cdot g_{TQL}), (g_{TQQ} \cdot g_{Tue})$.²

Other processes of the proton decay into a charged lepton, such as $p \rightarrow e^+ \gamma$, give weaker bounds.

Likewise, the above operators generate the process of neutron decay into a charged lepton $n \rightarrow e^+ \pi^-$. These are within experimental bounds provided the couplings satisfy the proton decay constraints (10).

We now turn to the processes that involve the neutrino sector. For this, let us first identify the neutrino mass eigenstates. The neutrino mass originates from the following Yukawa coupling with H ,

$$Y'_\nu H L \nu^c + m_R \nu^c \nu^c. \quad (11)$$

For a nonzero Majorana mass, m_R , the ν^c entering this coupling is not a mass eigenstate. Instead, it is a superposition of the mass eigenstates $\nu^{c'}$ and ν' of the sterile and active neutrino flavors. For $m_R \gg Y'_\nu v$, the mixing angle can be approximated as $\vartheta \simeq Y'_\nu v / m_R$ and can be expressed via mass eigenvalues as $\vartheta \simeq \sqrt{m_\nu / m_R}$, where $m_\nu \simeq (Y'_\nu v)^2 / m_R$ is the mass of the active SM neutrino. The interaction and mass eigenstates are therefore related as

$$\nu^c \simeq \nu^{c'} + \sqrt{\frac{m_\nu}{m_R}} \nu', \quad \nu \simeq \nu' - \sqrt{\frac{m_\nu}{m_R}} \nu^{c'}. \quad (12)$$

Next, let us turn to the couplings of the T -particle that can mediate proton decay with a neutrino in the final state:

$$\begin{aligned} & g_{TQQ} T QQ + g_{TQL} T^* QL \\ & + g_{Tud} T^* u^c d^c + g_{Td\bar{\nu}} T d^c \nu^c. \end{aligned} \quad (13)$$

The corresponding four-fermion operators obtained by T -exchange come from the second and third rows of (9). For clarity, we shall present them here explicitly in terms of their neutrino entries,

$$\begin{aligned} & \frac{g_{Tud}g_{Td\bar{\nu}}}{m_T^2} u^c d^c d^c \nu^c + \frac{g_{TQQ}g_{TQL}}{m_T^2} udd\nu \\ & + \frac{g_{Tud}g_{TQL}}{m_T^2} u^c d^c d^{c*} \nu^* + \frac{g_{TQQ}g_{Td\bar{\nu}}}{m_T^2} udd^{c*} \nu^{c*}. \end{aligned} \quad (14)$$

The participation of the neutrino mass eigenstates in these couplings goes via the substitutions according to Eq. (12).

For $m_R > m_p$, the sterile neutrino participates in the proton decay only as a virtual state. We shall focus on such a case. The couplings in (13) and the corresponding

²We note that throughout the paper the displayed bounds on the parameters of the theory must be understood as order of magnitude wise relations with numerical factors not shown explicitly.

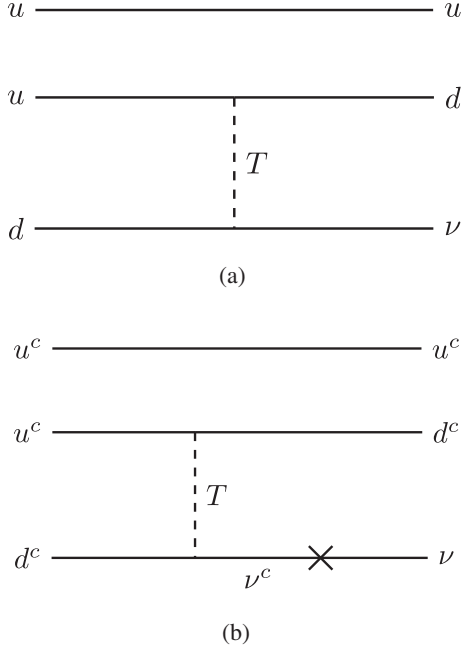


FIG. 2. Examples of diagrams responsible for proton decay into a pion and an active neutrino via different couplings. (a) via g_{TQL} . (b) via $g_{Td\nu}$.

effective operators (14) lead to the processes of proton decay in the channel $p \rightarrow \pi^+ \nu'$.

Some examples of representative diagrams are given in Figs. 2(a) and 2(b). The effective operators in (14) involving the coupling g_{TQL} lead to the lifetimes of the proton

$$\tau'_p \sim C \frac{m_T^4}{m_p^5}, \quad (15)$$

where the parameter C is set by the coefficient of the corresponding operator as $C = (g_{TQQ}g_{TQL})^{-2}$ or $C = (g_{Tud}g_{TQL})^{-2}$. Analogously, the operators involving $g_{Td\nu}$ give the lifetimes,

$$\tau_p \sim C \frac{m_R}{m_\nu} \frac{m_T^4}{m_p^5}, \quad (16)$$

with $C = (g_{TQQ}g_{Td\nu})^{-2}$ or $C = (g_{Tud}g_{Td\nu})^{-2}$.

Here, we distinguish the would-be lifetimes by a prime in order to differentiate between the processes involving the couplings g_{TQL} and $g_{Td\nu}$. Obviously, to have a neutrino in the final state, one of these couplings is necessary. The processes that go through $g_{Td\nu}$ exhibit an extra enhancement of the proton lifetime by a factor m_R/m_ν which comes from the mixing angle in (12).

The experimental bound on the proton lifetime in the decay channel $p \rightarrow \pi^+ \nu'$, $\tau_p > 10^{32}$ yr [28], translates into the following bounds on the coefficients in (15) and (16), respectively:

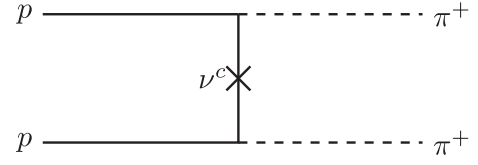


FIG. 3. Double proton decay.

$$C > 10^{50} \quad \text{and} \quad C \frac{m_R}{m_\nu} > 10^{50}. \quad (17)$$

Notice that if we assume the $SU(5)$ -relation, $g_{Td\nu} = Y'_\nu$, in the expression (16) the coupling and m_R can be traded for the known parameters of the SM such as the Higgs VEV, v , and the neutrino mass m_ν . This gives

$$\tau_p \sim (C g_{Td\nu}^2) \left(\frac{v}{m_\nu} \right)^2 \frac{m_T^4}{m_p^5}. \quad (18)$$

Taking $m_\nu \sim 10^{-2}$ eV, the experimental bound $\tau_p > 10^{32}$ yr [28] implies

$$g_{Tud}, g_{TQQ} < 10^{-12}. \quad (19)$$

The above constraints also satisfy the bounds on neutron decay.

The Majorana mass of ν^c can also lead to a process of double proton decay $pp \rightarrow \pi^+ \pi^+$ (see Fig. 3). However, the experimental bound on this process is 10^{31} yr [29], which puts weaker bounds on our parameters as compared to the bounds discussed above.

B. Neutron oscillations into the sterile neutrino

We now wish to point out that the light T -particle can induce the interesting process of neutron oscillations into the sterile neutrino. This is possible if m_R is close to the mass of the neutron m_n . The triplet exchange generates a mixing mass term between n and $\nu^{c'}$, which we shall denote by $m_{n\nu^c}$. This mixing term induces the oscillations of a free neutron into the sterile neutrino with the period $\tau_{n\nu^c} = m_{n\nu^c}^{-1}$. The remarkable thing is that the transition process is correlated with the process of proton decay in the channel $p \rightarrow \pi^+ \nu'$, essentially with no free parameters involved.

Indeed, all processes leading to this decay have counterparts that generate the mixing term $m_{n\nu^c}$. For example, a pair of corresponding representative diagrams is depicted in Fig. 4.

The induced oscillation periods are in one-to-one correspondence with the expressions for the proton lifetimes. The difference is that the processes that go through the coupling g_{TQL} exhibit a relative enhancement by a factor $\sqrt{\frac{m_R}{m_\nu}}$ in neutron oscillation period versus the proton lifetime, as compared to the processes that go via $g_{Td\nu}$.

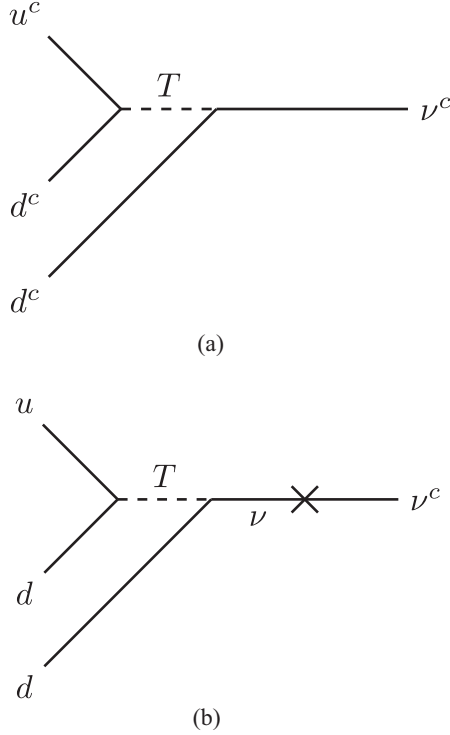


FIG. 4. Examples of diagrams mixing neutron with sterile neutrino via different couplings. (a) via $g_{Td\nu}$. (b) via g_{TQL} .

To be more explicit, for processes mediated via g_{TQL} , we have the following relation between the two timescales,

$$\tau_{n\nu^c} \simeq \sqrt{\frac{\tau'_p}{m_\nu}}, \quad (20)$$

whereas for the processes mediated via $g_{Td\nu}$, the analogous relation is

$$\tau_{n\nu^c} \simeq m_n^{-1} \sqrt{\tau_p m_\nu}. \quad (21)$$

In both expressions, we took $m_R = m_n = m_p$. It is clear that for τ'_p and τ_p at the current experimental bound on the proton lifetime the expression (21) gives a shorter oscillation time. Taking $\tau_p \sim 10^{32}$ yr [28], we get for the oscillation period of a free neutron into the sterile neutrino $\tau_{n\nu^c} \sim 10^2 - 10^3$ s.

This timescale is within the range of interest for the experiments that measure the disappearance of cold neutrons into hidden species [30,31]. Both experiments are set up for testing the transition of a neutron into a nearly degenerate fermion. Therefore, they are ready made for probing the proposed transition of the neutron into the sterile neutrino. These experiments place the ultracold neutrons in a magnetic field B which changes the energy

of the neutron proportional to its magnetic moment μ_n , $\Delta E = \mu_n B$. When the energies of the neutron and the partner particle are nearly degenerate, a resonance increase in the neutron disappearance probability takes place. Ban *et al.* [31] performed a scanning of the magnetic field in steps of $\Delta B = 3 \mu\text{T}$ within the range of $50 - 1100 \mu\text{T}$. Correspondingly, they were able to put a bound on the neutron disappearance time $\tau_{nn'} > 1$ s for a mass splitting $2 - 69$ peV. In the present context, this bound translates as the corresponding constraint on $\tau_{n\nu^c}$ for the “resonant” regime $m_n + \Delta E \simeq m_R$.

If the splitting between m_n and m_R is significant, the transitions of the free neutron into ν^c are suppressed. However, for certain values $m_n > m_R$, the mass m_R can come in resonance with some energy levels E_n of a bound neutron inside the nucleus. In this case, the bounds on the couplings of T become more stringent in order to accommodate the neutron disappearance bounds from the nuclei, which gives the timescale $> 10^{30}$ yr [32,33].

Notice that for having a resonant transition it is sufficient that the difference between m_R and the neutron energy level E_n is within the binding energy of the neutron $\Delta E_n \equiv m_n - E_n$, which is $\Delta E_n \sim \text{MeV}$. The reason is that in such a case the neutron can transition into a sterile neutrino with a momentum $\mathbf{p}$ such that $E_n = \sqrt{m_R^2 + |\mathbf{p}|^2}$. This is because, for $|\mathbf{p}| < \sqrt{\Delta E_n m_R}$, the momentum can be absorbed by the recoil of the nucleus.

Since the would-be oscillation period is much longer than the nuclear deexcitation time as well as the escape time of ν^c , the inverse transition is not possible. Correspondingly, the outcome of the process will be a transition of the “host” atom into an isotope with fewer neutrons. This transition will be accompanied by the emission of photons of nuclear energy. The nonobservation of this effect puts the bound on the disappearance time $\sim 10^{30}$ yr [32,33].

Notice that once the resonant $n - \nu^c$ transition for a nuclear neutron satisfies this bound the corresponding two-body decays of the neutron and the proton, such as $p \rightarrow \pi^+ \nu$ and $n \rightarrow \pi^0 \nu$, are well below the experimental limits.

To summarize, for various resonant values of m_R with neutron energy levels, the T -particle can mediate transitions between the neutron and ν^c for a free as well as for a bound neutron. This provides a new set of low-energy probes of grand unification linked with the neutrino mass. In principle, the resonant values can be scanned by the external magnetic field. One can say that in this setup the neutron energy levels provide a spectroscopy of the sterile neutrino.

Of course, in this setup, we treat the mass of the sterile neutrino as a parameter of the effective theory and derive its observable consequences. The exciting feature is that it gives a chance to probe GUT physics in low-energy measurements with some spectacular signatures and correlate it with other observations.

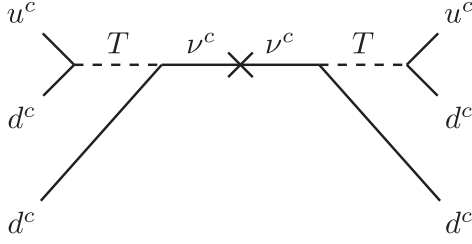


FIG. 5. An example of a diagram generating the neutron-antineutron transition.

C. Neutron-antineutron transitions

Since in the above parameter space ν^c is a Majorana fermion, its mixing with the neutron automatically induces a mixing between the neutron and the anti-neutron; see Fig. 5. The corresponding transition time for a free neutron is

$$\tau_{n\bar{n}} \sim \tau_{n\nu^c}^2 m_n \sim 10^{30} \text{ s}, \quad (22)$$

which is very far from the latest experimental limits, $\tau_{n\bar{n}} > 0.86 \times 10^8 \text{ s}$ [34].

For a neutron that is in a bound state inside a nucleus, taking into account the additional suppression due to a level splitting between n and $\bar{n}$, this gives a transition time $\tau_{n\bar{n}}^{(\text{nucl})} \sim 10^{75} \text{ yr}$, which by many orders of magnitude exceeds the current experimental bound $\tau_{n\bar{n}}^{(\text{nucl})} \sim 3.6 \times 10^{32} \text{ yr}$ [35].

Thus, for the resonant regime $m_n \simeq m_R$, the dominant process allowed by the proton stability constraint is the oscillation between a neutron and a sterile neutrino.³ In general, in this regime, the $n - \bar{n}$ oscillation period $\tau_{n\bar{n}}$ is additionally enhanced by the factor $\tau_{n\nu^c} m_n$ relative to $\tau_{n\nu^c}$. This is because the $n - \bar{n}$ transition has to go via an intermediate ν^c state.

D. Correlations with T -production at LHC

The T -particle can be pair produced at the LHC in the collision of ordinary hadrons. Early discussions of its production mechanisms and collider signatures were given in Refs. [1,2,10,15,37,38]. More recent studies of bounds on new colored particles in the light of LHC data [39–44] set the lower bound on the mass of the T -particle around TeV.

For this mass, the correlations between various baryon number violating processes were discussed above. These processes constrain the lifetime of T , which is important for its LHC phenomenology.

³The correlation with proton decay is an important constraint on $n - \bar{n}$ versus $n - \nu^c$ transitions exhibited by the present GUT framework. This goes in sharp difference with alternative extensions of the SM such as the $n - \bar{n}$ oscillations via mirror copies of the neutron [36] where such correlations are absent.

For example, assuming that the dominant couplings are given by the second row of (13), the T -triplet can decay either into an up and a down quark or into a down quark and a sterile neutrino. The dominance of the channel depends on the relative strength of the couplings, g_{Tud} versus g_{Tdv} . In the simplest case, in which the couplings of the doublet and triplet to the sterile neutrino are given by the $SU(5)$ -relation, we can estimate

$$g_{Tdv} \sim \frac{\sqrt{m_\nu m_R}}{v} \sim 10^{-8}. \quad (23)$$

The T -quark–quark coupling is then constrained by the proton decay as

$$g_{Tud} < 10^{-12}. \quad (24)$$

The corresponding lifetime of the T -particle is

$$\tau_T \sim 10^{-11} \text{ s} \sim 1 \text{ mm}. \quad (25)$$

Since this timescale is much longer than the QCD length, the T -triplet will have plenty of time to hadronize and to form a color-singlet state. In particular, it can form a “ T -baryon” state with two fermions of the same color ($Tqq, T^*q^c q^c$) or a “ T -meson” state with a fermion of an anticolor (Tq^c, T^*q). Notice that, since the T -particle is a scalar, with such a definition, the T -baryon is a boson, whereas T -meson is a fermion.

Assuming that upon its production the T -particle is moderately relativistic, the hadronized T -particle will decay via a displaced vertex. Depending on the precise values of the parameters, the decay can take place within the detector. For example, for a T -meson the decay products can be an ordinary QCD meson and a sterile neutrino, e.g.,

$$(Td^c) \rightarrow \pi^0 + \nu^c. \quad (26)$$

Analogously, for a T -baryon, the decay can take place into an ordinary baryon and a sterile neutrino, e.g.,

$$(T^*u^c d^c) \rightarrow \bar{n} + \nu^c. \quad (27)$$

In both cases, the decay of the T -hadron produces one jet in the detector. If instead the T -particle decays into two quarks, the decay of the T -hadron will be accompanied by two jets. Of course, for more suppressed values of its couplings, the decay can take place outside of the detector, in which case the T -hadron will appear as a stable particle. The various parameter regimes require more precision studies. However, the general message is that for the light GUT color-triplet scenario [1,2] the searches for proton decay on one hand, and the cold neutron oscillations/disappearance experiments on the other hand become correlated with each other as well as with the production of a long-lived T -hadron states in colliders (see Fig. 6).

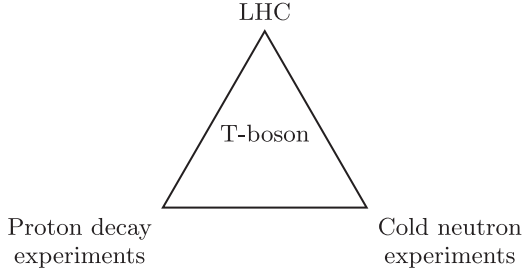


FIG. 6. A schematic representation of correlations between various experimental searches.

IV. CONCLUSION

The color-triplet partner of the Higgs doublet is a model-independent property of any GUT. Therefore, an experimental detection of the T -boson would be direct evidence of grand unification.

As originally pointed out in Refs. [1,2], the T -particle can be part of the low-energy spectrum of the theory with the mass potentially accessible at collider experiments, such as the LHC. The reason is that already the lowest-order operators involving the VEV of 24_H are capable of suppressing the baryon and lepton number violating couplings of the T -boson to the levels compatible with the current experimental bounds.

This opens up an avenue for simultaneous searches of the signatures of the T -particle at high-energy colliders as well as in the experiments looking for baryon and lepton number violating low-energy processes.

In particular, in the present paper, we have pointed out the potential correlations among the searches at three very different sets of experiments. The first one comes from the link between proton decay and the oscillations of a cold neutron into a sterile neutrino. Namely, the current bounds on the proton lifetime put the T -mediated neutron-to-sterile neutrino oscillation processes within the ballpark of sensitivity of cold neutron experiments [30,31].

In general, due to its vanishing SM gauge charges, the neutron represents an interesting probe of feebly coupled new physics. The early examples of such effects include the neutron transitions into its hidden copies [45–47]. It was shown recently [48] that within the framework of large extra dimensions [49], the “disappearance” or oscillations of a cold neutron into the extra dimensions is a rather generic phenomenon. Moreover, this effect can be linked with the extra-dimensional origin of the neutrino mass [50,51].

The analysis of the present work shows that a similar phenomenon can take place in GUTs. Moreover, here, too, the cold neutron transition is linked with the origin of the neutrino mass. Namely, via T -exchange, a neutron can oscillate into a sterile neutrino which generates the Majorana mass for the SM neutrino. Thus, the GUT framework with a light T -particle can correlate the T -mediated proton and neutron decay processes with

the oscillations of cold neutrons into the sterile neutrino species. In this way, the cold neutron experiments can potentially turn out to be simultaneous probes of grand unification and of the neutrino mass.

In addition, the parameter space predicting the signatures for both the above types of measurements, at the same time, is linked with the direct production of the T -particle at the LHC. Thus, the T -boson, which represents an intrinsic part of any GUT, can manifest itself in the correlated signatures in three different types of experiments.

We wish to make some comments about the influence of higher-dimensional operators on other predictions of GUTs. First, as already mentioned in the Introduction, the inclusion of higher-dimensional operators affects the low-energy predictions of grand unification. This should not be taken as a drawback of the concept of GUTs but as an inevitable consequence of its scale being not far from the cutoff of the low-energy theory such as the species scale in Eq. (7). The proximity to the cutoff is expected to produce threshold corrections, which can only be resolved within the UV theory. This is irrespectively of the effect of the light triplet on the perturbative unification scale [22].

Although we have illustrated the effect within the minimal $SU(5)$ with a single generation of fermions, it is straightforward to see that it persists in the generalizations to higher GUTs such as $SO(10)$ or $SU(6)$, as well as in their supersymmetric extensions which have been formulated previously [1,2,6,16,17]. Moreover, the inclusion of couplings of the T -particle with heavier generations of fermions provides additional baryon decay channels and correlated transitions of the neutron into the corresponding flavors of the sterile neutrinos.

Another important point concerns the top Yukawa coupling. This Yukawa coupling is order 1,⁴ and so is the renormalizable coupling of the Higgs triplet to the third generation of fermions originating from $Y_u 10_F 10_F 5_H$. Therefore, its corrections by higher-dimensional operators raises the question of calculability since the corresponding higher-dimensional operator must have order one effect at the GUT scale. We would like to stress that, while giving up the perturbative calculability for the top quark at the expense of cutoff physics remains a consistent option within our philosophy, this is not the only option. In fact, effective higher-dimensional operators which determine the top quark Yukawa coupling can originate from renormalizable physics. A transparent and predictive realization of this mechanism can be found in Ref. [54].

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⁴For implications in SUSY unification, see, e.g., Refs. [52,53].

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DATA AVAILABILITY

No data were created or analyzed in this study.

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