Suplemento de la Revista Mexicana de Física **3** 020710 (2022) 1–8

One loop mechanism for neutrinoless double beta hyperon decay

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Received 31 December 2021; accepted 7 February 2022

Motivated by the large dataset to be accumulated of hyperon pairs produced in decays of $\mathcal{O}(10^{10})$ J/ψ and ψ' charmonia states in the BES-III collaboration, we revisited the predictions of $\Delta L = 2$ decays of hyperons $B_i^- \to B_f^+ \ell^- \ell'^-$ in the one-loop model mechanism involving Majorana neutrinos previously presented in Ref. [1]. Unlike the previous work, by modeling the momentum transfer dependence of the hyperon form factors in the computation we provide finite results for the loop integration. Furthermore, since we keep finite masses for the neutrinos throughout the calculations, we are able to consider the effects of heavy Majorana neutrinos. Thus, our results are applied to a simple model that involves two Majorana heavy neutrinos in the framework of a low-scale seesaw model. In order to provide and compare additional predictions, we study an alternative model where $\Delta L = 2$ decays are induced by the short-range effects of a scalar boson coupled to di-leptons.

Keywords: Neutrinos; hyperons; lepton number violation.

DOI: https://doi.org/10.31349/SuplRevMexFis.3.020710

1. Introduction

It is well know that the most extensive and sensitive laboratory to probe lepton number violation (LNV) and consequently the Majorana nature of neutrinos is neutrinoless double beta decay ($\beta\beta_{0\nu}$) of nuclei. If $\beta\beta_{0\nu}$ decays are generated by the so-called *mass mechanism* the amplitude is proportional to the "*effective Majorana neutrino mass*" m_{ee}, which is defined by

$$m_{ee} \equiv \sum_j m_{\nu_j} U_{ej} U_{ej}, \qquad (1)$$

where j runs over all the neutrino mass eigenstates, and U describes the mixing matrix in the leptonic sector. Because of the non-observation of $\beta\beta_{0\nu}$ decays in nuclei direct strong bounds have been set on m_{ee} at the sub-eV level [4]. On the other hand, nuclear transitions are limited in precision due to the model-dependent uncertainties of the nuclear matrix elements, and are sensitive only to the two electrons channel. This fact has motivated an extensive study of much less sensitive but complementary $\Delta L = 2$ transitions. For example, LNV transitions involving other lepton flavors are possible via semileptonic tau and meson decays, Λ_b baryons, di-muon production at colliders, or muon to positron/antimuon conversion in nuclei[i].

In this work, we are interested in $\Delta Q = \Delta L = 2$ hyperon decay ($B_i^- \to B_f^+ \ell^- \ell'^-$, $\ell^{(\prime)} = e$ or μ). As already mentioned, their study is complementary to those in nuclei, but with the advantage that the hadronic matrix elements involved are well known at low momentum transfer. Particularly, we aim to revisit and improve the estimation reported in Ref. [1], where this kind of decays are generated via a one-loop mechanism considering baryons as the relevant degrees of freedom (involving a loop with hyperons and Majorana neutrinos as intermediate states) as shown in Fig. 1. In Ref. [1], the authors neglected the momentum transfer dependence of the vector and axial form factors describing the weak vertices in their computation. Consequently, the obtained loop functions exhibit a logarithmic ultraviolet divergence which was regulated using a simple cut-off procedure.

There is in the literature another prediction for $\Delta L = 2$ hyperon decays (specifically for $\Sigma^- \to p e^- e^-$) presented by the same authors in Ref. [2] but this time based on the so called MIT bag model. The results of these two previous computations are rather different: the branching ratio of the $\Sigma^- \to p e^- e^-$ decay in the MIT bag model is of the order $O(10^{-23})$ [2], which is around ten orders of magnitude larger than its prediction based on the loop model $O(10^{-33})$ [1].

In the present work, we provide an improvement for the estimation of the ΔL=2 hyperon decays in the loop model mechanism and we help to understand the numerical difference found between the one-loop mechanism in and the MIT-bag model estimation by studying an alternative mechanism induced by a doubly charged Higgs boson H^{--}, using the current bounds on the search of such hypothetical particles. For a complete and extended description of this work we refer the reader to Ref. [3].

2. $\Delta L = 2$ hyperon decays (One loop-mechanism)

All the possible $\Delta L = 2$ decays (for 1/2-spin hyperon) channels can be classified according to their change in strangeness $\Delta S = 0, 1, 2$ (see Table I).

At the hadron level, this process can be viewed as induced by the one-loop mechanism shown in Fig. 1 involving into the loop a neutral baryon η and a Majorana neutrino. In such case, the amplitude is given by

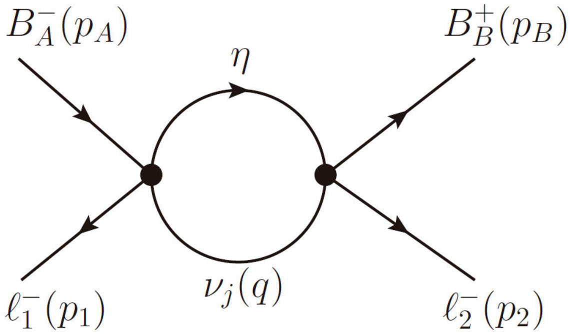

FIGURE 1. $\Delta L=2$ hyperon decays in the one-loop model and in the presence of Majorana neutrinos.

TABLE I. 1/2-spin hyperon $\Delta L=2$ decays allowed by kinematics.

$\Delta S=0$	$\Delta S=1$	$\Delta S=2$
$\Sigma^-\to\Sigma^+e^-e^-$	$\Sigma^-\to pe^-e^-$	$\Xi^-\to pe^-e^-$
	$\Sigma^-\to pe^-\mu^-$	$\Xi^-\to pe^-\mu^-$
	$\Sigma^-\to p\mu^-\mu^-$	$\Xi^-\to p\mu^-\mu^-$
	$\Xi^-\to\Sigma^+e^-e^-$	
	$\Xi^-\to\Sigma^+e^-\mu^-$	

$$i\mathcal{M}=-G^2\sum_j m_{\nu_j}U_{\ell_1 j}U_{\ell_2 j}\sum_\eta\int\frac{d^4q}{(2\pi)^4}\frac{L_1^{\alpha\beta}(p_1,p_2)}{[q^2-m_{\nu_j}^2]}\times\frac{h_{1\alpha\beta}(p_A,p_B)}{[Q_1^2-m_\eta^2]}-[\ell_1(p_1)\leftrightarrow\ell_2(p_2)]\,, \quad (2)$$

where m_{ν_j} are the masses of Majorana neutrinos and $U_{\ell j}$ their mixings connecting flavor ℓ and mass eigenstates. The overall constant $G^2=G_F^2\times(V_{ud}^2,V_{ud}V_{us},V_{us}^2)$ for $\Delta S=0,1,2$, respectively, with G_F the Fermi constant. Whereas $Q_i=p_A-p_i-q$ is the momentum carried by the neutral baryon state η with the appropriate quantum numbers to contribute as an intermediate state, and $\ell_1(p_1)\leftrightarrow\ell_2(p_2)$ stands for the contribution of a similar diagram interchanging the final external charged leptons. The hadronic and leptonic are given by

$$L_1^{\alpha\beta}(p_1,p_2)\equiv\bar{u}(p_2)\gamma^\alpha(1-\gamma_5)\gamma^\beta v(p_1), \quad (3)$$

$$h_{1\alpha\beta}(p_A,p_B)\equiv\bar{u}(p_B)\gamma_\alpha\left[f_{B\eta}(q''^2)+g_{B\eta}(q''^2)\gamma_5\right]\times(\not{Q}_1+m_\eta)\gamma_\beta\left[f_{A\eta}(q'^2)+g_{A\eta}(q'^2)\gamma_5\right]u(p_A). \quad (4)$$

$f_{(A,B)\eta}$ and $g_{(A,B)\eta}$ denote the vector and axial weak form factors, respectively, for transitions $A\to\eta$ and $\eta\to B$. They depend on the squared momentum transfer at each weak vertex, specifically, $f_{A\eta}$ and $g_{A\eta}$ depend on $q'^2=(p_1+q)^2$, whereas $f_{B\eta}$ and $g_{B\eta}$ depend on $q''^2=(p_2-q)^2$. Their values at zero momentum transfer have been calculated by different groups [34–36] with overall good agreement among them (in Table II we quote the values reported in Ref. [35]).

TABLE II. Vector and axial transition form factors for weak hyperon decays at zero momentum transfer. Here η stands for the intermediate baryon state, and the subscript A (B) represents the initial (final) baryon [35].

Transition	η	$f_{A\eta}$	$g_{A\eta}$	$f_{B\eta}$	$g_{A\eta}$
$\Sigma^-\to\Sigma^+$	Λ	0	0.656	0	0.656
	Σ^0	$\sqrt{2}$	0.655	$\sqrt{2}$	-0.656
$\Sigma^-\to p$	n	-1	0.341	1	1.267
	Σ^0	$\sqrt{2}$	0.655	$-1/\sqrt{2}$	0.241
	Λ	0	0.656	$-\sqrt{3/2}$	-0.895
$\Xi^-\to\Sigma^+$	Ξ^0	-1	0.341	1	1.267
	Σ^0	$1/\sqrt{2}$	0.896	$\sqrt{2}$	-0.655
	Λ	$\sqrt{3/2}$	0.239	0	0.656
$\Xi^-\to p$	Σ^0	$1/\sqrt{2}$	0.896	$-1/\sqrt{2}$	0.241
	Λ	$\sqrt{3/2}$	0.239	$-\sqrt{3/2}$	-0.895

After some redefinitions (see Ref. [3] for further details) the amplitude can be expressed as follows

$$i\mathcal{M}=-G^2\left(L_1^{\alpha\beta}(p_1,p_2)H_{1\alpha\beta}-L_2^{\alpha\beta}(p_1,p_2)H_{2\alpha\beta}\right), \quad (5)$$

with

$$H_{1\alpha\beta}=\sum_{\eta,j}m_{\nu_j}U_{\ell_1 j}U_{\ell_2 j}\Big\{\bar{u}(p_B)\gamma_\alpha\Big[\left(C_{1v_0}^{\eta j}+C_{1a_0}^{\eta j}\gamma_5\right)m_\eta+\left(C_{1v_1}^{\eta j}+C_{1a_1}^{\eta j}\gamma_5\right)\not{p}_1+\left(C_{1v_2}^{\eta j}+C_{1a_2}^{\eta j}\gamma_5\right)\not{p}_2+\left(C_{1v_A}^{\eta j}+C_{1a_A}^{\eta j}\gamma_5\right)\not{p}_A\Big]\gamma_\beta u(p_A)\Big\}, \quad (6)$$

and the $C_{1v_r}^{\eta j}$ and $C_{1a_r}^{\eta j}$ $(r=0,1,2,A)$ functions encode the effects of the strong interaction relevant in the loop computation. Note that they will depend on the neutrino masses and on the two independent Mandelstam variables $t\equiv(p_A-p_1)^2$, and $u\equiv(p_A-p_2)^2$.

The important point to stress here is that taking the form factors as constants in Eq. (4) is just an approximation that leads to a divergent amplitude and this bad behavior can be cured by a form factor that vanishes at large q^2.

In this way, the loop integration requires a proper modeling of hyperon form factors in all the range of momentum transfer scales. From neutrino and electron scattering off nucleons it has been found that the observed distributions can be described by a dipole parametrization and extrapolating to the timelike region leads to the dipole form factors given by

$$f_i(q^2)=f_i(0)\left(1-\frac{q^2}{m_{df_i}^2}\right)^{-2}, \quad (7)$$

$$g_i(q^2)=g_i(0)\left(1-\frac{q^2}{m_{dg_i}^2}\right)^{-2}, \quad (8)$$

with $m_{df_i}=0.84$ (0.97) GeV and $m_{dg_i}=1.08$ (1.25) GeV for the strangeness-conserving (strangeness-changing) form factors. On the other hand, SU(3) symmetry considerations are useful to fix the form factors at zero momentum transfer ($q^2 = 0$) (see Table II).

Now, note that in the limit of large momentum transfer, both the vector and axial dipole form factors behave as $\sim 1/q^4$. Then, considering the dipole approximation would require evaluating an integral with six propagators which, in general, are very difficult to evaluate even numerically, a better option is to take, in analogy with the case of meson form factors considered in Ref. [37] which behave as $\sim 1/q^2$ for large q^2, a monopolar approximation.

The "equivalence" of monopole m and dipole d form factors can be achieved by comparing their slopes at low momentum transfers; this leads to identify $m_d/\sqrt{2} \to m_m$ for the vector and axial poles to get the monopolar form factors in the $\Delta S = 0, 1$ cases. Using monopolar expressions for the form factors, the loop integrals become finite and, as a consecuence, the amplitudes are physical[ii].

2.1. Numerical analysis

The relevant hadronic matrix element defined in Eqs. (5), (6) depends on the effective total form factors $C^{\eta}_{v_r,a_r}$; for each intermediate hadronic state η in the loop, they can be written as

$$C^{\eta}_{v_r} \equiv \sum_j m_{\nu_j} U_{\ell_1 j} U_{\ell_2 j} C^{\eta j}_{v_r}, \quad (v_r = v_0, v_1, v_2, v_A),$$
$$C^{\eta}_{a_r} \equiv \sum_j m_{\nu_j} U_{\ell_1 j} U_{\ell_2 j} C^{\eta j}_{a_r}, \quad (a_r = a_0, a_1, a_2, a_A), \quad (9)$$

where the individual $C^{\eta j}_{v_r,a_r}$ factors, determined from the computation, depend in general upon the neutrino mass m_{ν_j}

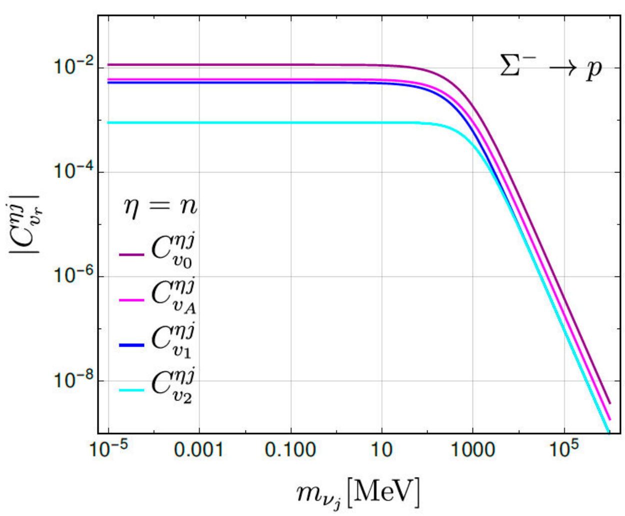

FIGURE 2. Individual $C^{\eta j}_{v_r}$ loop-factors as function of the neutrino mass in the monopole form factors model for the decay chain $\Sigma^- \to n^* \to p$. Note that these functions depend on the t and u Maldemstam variables. For illustration purposes, we have used the maximum allowed values for t and u, whereas $m_1 = m_2 = m_e$.

involved in the neutrino propagator. Figure 2 shows the absolute value of the $C^{\eta j}_{v_r}$ (with $\eta = n$) form factors as a function of the intermediate neutrino mass m_{ν_j} in the case of the $\Sigma^- \to p$ transition, using the monopolar approximation (similar results are obtained for the rest of the decay channels listed in Table II as well as for the analysis of the axial $C^{\eta j}_{a_r}$ form factors). From this plot we observe that the dominant contribution arises from the $C^{\eta j}_{v_0}$ coefficient. Also, for light neutrinos ($m_{\nu_j} \lesssim 100$ MeV), all the $C^{\eta j}_{v_r}$ factors are insensitive to the neutrino mass value. However, for the case of heavy neutrinos, the $C^{\eta j}_{v_r}$ one-loop functions describing the $\Delta L = 2$ hyperon decays become strongly-dependent on the neutrino mass. This led us to consider two different scenarios: we call *scenario A* to the familiar case where only three light Majorana neutrinos are present; the second *scenario B* corresponds to the contributions of heavy Majorana neutrinos such as those appearing in the so-called low scale seesaw models.

Scenario A: Assuming that only very light neutrino states exist, the total form factors in Eq. (9) can be approximated by

$$C^{\eta}_{v_r} \equiv m_{\ell_1\ell_2} C^{\eta 0}_{v_r}, \quad C^{\eta}_{a_r} \equiv m_{\ell_1\ell_2} C^{\eta 0}_{a_r}, \qquad (10)$$

where $m_{\ell_1\ell_2}$ is the effective Majorana mass parameter, and $C^{\eta 0}_{v_r}$ ($C^{\eta 0}_{a_r}$) are the one-loop functions evaluated at $m_{\nu_j} = 0$. Using the direct upper limits for $m_{\ell\ell'}$ reported in Ref. [38]:

$$m_{ee} = 0.36\,\text{eV},$$
$$m_{e\mu} = 90\,\text{GeV},$$
$$m_{\mu\mu} = 480\,\text{GeV}, \qquad (11)$$

and by computing numerically the form factors in Eq. (10), we obtained the rates listed in the second column of Table III. Note that channels involving two electrons are strongly suppressed due to the strong limits imposed from neutrinoless nuclear $\beta\beta_{0\nu}$ decay. Contrary, channels with $e\mu$ or $\mu\mu$ in the final states have loosely bounds due to the very poor direct limits reported in Eq. (11). In particular, the channel $\Sigma^- \to p\mu\mu$ appears to be close to the future sensitivity projected by BES-III. Nevertheless, those numbers should be taken with care since the upper limits used for $m_{e\mu}$ and $m_{\mu\mu}$ lie beyond the range of validity of *scenario A*. If we assume much lower values of these two parameters as expected if the three neutrinos were very light, these channels are also very suppressed (see Table III for a comparison with the rates from *scenario B*).

***Scenario B*:** In order to assess the effects of LNV in this case we will consider the minimal parametrization presented in references [41, 42].

Figure 3 illustrates the behaviour of the total dominant $C^{\eta}_{v_0}$ coefficient in Eq. (9) associated to the $\Sigma^- \to p\ell\ell'$ ($\eta = n$; $\ell\ell' = ee, e\mu$) decay as a function of the r parameter. A similar behavior is obtained for the rest of the total vector $C^{\eta}_{v_r}$ ($v_r = v_1, v_2, v_A$) form factors[iii]. The following comments are worth being pointed out:

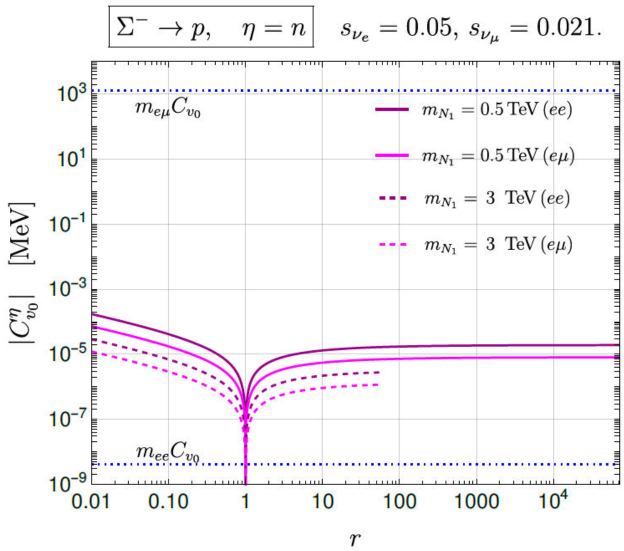

FIGURE 3. Behavior of the total $C^{\eta}_{v_0}$ form factor as a function of r for the $\Sigma^- \to p\ell\ell'$ decay (with intermediate state $\eta = n$) in the low-scale seesaw model of Ref. [41, 42]. Lepton number is conserved when the two heavy neutrino states form a heavy Dirac neutrino singlet ($r = 1$). Purple and magenta lines are associated with the $\Sigma^- \to pee$ and $\Sigma^- \to pe\mu$ decays, respectively (plots for the $\mu\mu$ channel are not shown as they are almost identical to the $e\mu$ case). The lines stop in the non-perturbatively allowed range (see text). The horizontal blue dotted lines stand for comparison with *scenario A* with $m_{ee} = 0.36$ eV and $m_{e\mu} = 90$ GeV. Finally, the lower (upper) horizontal blue dotted line represents the $C^{\eta}_{v_0}$ form factor associated to the $\Sigma \to pee$ ($\Sigma \to pe\mu$) decay in *scenario A*.

- We focus on heavy neutrino masses around a few TeV where the direct limits on the heavy-light mixings are less restrictive. Nevertheless, we consider as a benchmark the mass-independent indirect limits coming from the latest global fits to electroweak precision observables given in Refs. [43, 44]:

$$s_{\nu_e} < 0.050, \quad s_{\nu_\mu} < 0.021, \quad s_{\nu_\tau} < 0.075. \tag{12}$$

- Using all the values presented in Fig. 3 we have that: the contributions of heavy neutrinos to the total $C^{\eta}_{v_0}$ form factor of the $\Sigma \to pee$ channel (purple-lines) can be increased in the *scenario B* by up to a factor $\sim 10^5$ ($\sim 10^{10}$ at the Branching ratio level see Table III). However, these effects are still well below the projected sensitivities at current experiments like BESIII [39]. In contrast, the $\Sigma^- \to pe^-\mu^-$ and $\Sigma^- \to p\mu^-\mu^-$ channels are far more restrictive than the ones obtained using the direct limits in the *scenario A* ('naive' approximation). Notice that the $C^{\eta}_{v_0}$ form factor of the $e\mu$ channel (magenta-lines) in *scenario B* is far below the one in *scenario A*. Table III shows the branching ratios for all the channels comparing both scenarios.

TABLE III. Branching ratios of $\Delta L = 2$ hyperon decays induced by Majorana neutrinos. *Scenario A*: we consider the upper limits of the effective Majorana masses given in Eq. (11). *Scenario B*: in addition to the limits in Eq. (12), we use the representative values $m_{N_1} = 1$ TeV and $r = 0.01$ ($m_{N_2} = 100$ GeV).

Transition	Branching Ratio	Branching Ratio
	Scenario A	*Scenario B*
$\Sigma^- \to \Sigma^+ ee$	7.6×10^{-41}	3.6×10^{-32}
$\Sigma^- \to pee$	1.0×10^{-33}	4.3×10^{-25}
$\Sigma^- \to p\mu\mu$	1.7×10^{-10}	1.2×10^{-27}
$\Sigma^- \to p\mu e$	1.6×10^{-12}	6.8×10^{-26}
$\Xi^- \to \Sigma^+ ee$	9.9×10^{-36}	3.7×10^{-27}
$\Xi^- \to \Sigma^+ \mu e$	1.8×10^{-14}	8.6×10^{-32}
$\Xi^- \to pee$	3.4×10^{-35}	2.5×10^{-26}
$\Xi^- \to p\mu\mu$	2.5×10^{-11}	3.3×10^{-28}
$\Xi^- \to p\mu e$	2.3×10^{-12}	6.1×10^{-27}

3. $\Delta L = 2$ hyperon decays induced by a doubly charged scalar boson

Majorana neutrinos are the most appealing but not unique mechanism to generate $\Delta L = 2$ transitions in hyperons. We also explore the possible effects that can arise in the presence of doubly charged scalar bosons coupled to dileptons, particularly, in the so-called Higgs Triplet Model (HTM) [45].

Here, we only focus on the phenomenology of the doubly charged states; a complete list of all the new vertices and the corresponding Feynman rules for the HTM can be found in Ref. [46]. The amplitude of the dominant diagram depicted in Fig. 4 is given by

$$i\mathcal{M} = 2\sqrt{2}G^2 h_{\ell_1\ell_2}\frac{v_\Delta}{M^2_{H^{\pm\pm}}} g_{\mu\nu}\,\bar{u}(p_2)(1-\gamma_5)v(p_1) \times \langle B^+_B(p_B)|\Gamma_{\mu\nu}|B^-_A(p_A)\rangle - (p_1 \leftrightarrows p_2), \tag{13}$$

where $\langle B^+_B(p_B)|\Gamma_{\mu\nu}|B^-_A(p_A)\rangle$ are the hadronic matrix elements describing the transition from the B^-_A hyperon to B^+_B. The hadronic current tensor $\Gamma_{\mu\nu}$ contains four quark operators. This model provides a concrete realization of the local six-fermion effective Lagrangian proposed in Ref. [40].

Since the doubly charged scalar couples to quarks via intermediate W gauge bosons as shown in Fig. 4, the current tensor is the product of two bilinear $V - A$ quark currents. Therefore, the hadronic matrix elements are given by [2]

$$\langle B^+_B(p_B)|\Gamma_{\mu\nu}|B^-_A(p_A)\rangle \equiv \langle B^+_B(p_B)|\,(\bar{u}\gamma_\mu(1-\gamma_5)D) \times (\bar{u}\gamma_\nu(1-\gamma_5)D')\,|B^-_A(p_A)\rangle, \tag{14}$$

where D and D' stand for down-type quarks d or s. In order to estimate the rate of $\Delta L = 2$ hyperon decays due to the amplitude (13), we will consider the results of Ref. [2] where the hadronic matrix elements were computed in the framewrok of the so-called MIT bag model [47]. In the non-relativistic

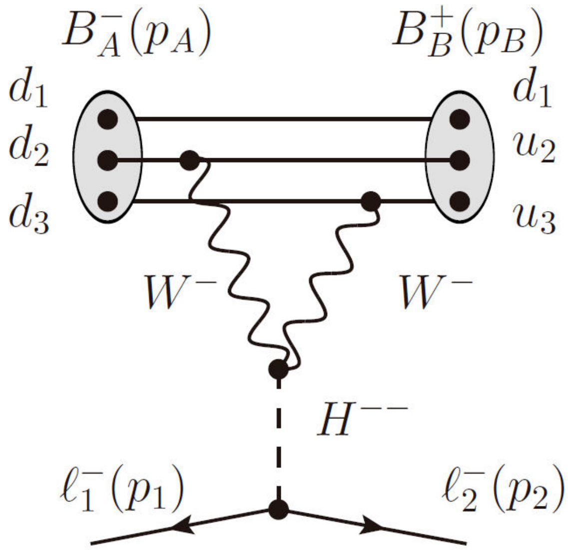

FIGURE 4. $\Delta L = 2$ hyperon decay mediated by a doubly charged scalar in the HTM. Similar contributions replacing each of the weak W^- bosons with a singly charged H^- scalar are suppressed due to small couplings $H^- q_d \bar{q}_u$ proportional to the light quark masses.

approximation, these hadronic matrix elements can be expressed in terms of only two A and B functions, in such a way that after the contraction of Lorentz indices in Eq. (13) the amplitude can be written in the simple form

$$i\mathcal{M} = 2\sqrt{2}G^2 \nu_\Delta \frac{h_{\ell_1\ell_2}}{M^2_{H^{\pm\pm}}} \bar{u}(p_2)(1-\gamma_5)v(p_1) \times \bar{u}(p_B)\,[A + B\gamma_5]\,u(p_A) - (p_1 \leftrightarrows p_2), \quad (15)$$

where $u(p_A)$ and $u(p_B)$ denotes the initial and final hyperon spinors, respectively. The A and B functions were obtained using the eigenfunctions of quarks confined within a baryon in the MIT bag model [2].

In addition to the overall decay parameters and the A and B form factors, the strenght of the decay amplitude in Eq. (13) is determined by the factor $v_\Delta h_{\ell_1\ell_2}/M^2_{H^{--}}$. The value of v_Δ is constrained from the correction to the ρ parameter, which after the introduction of the Higgs triplet becomes

$$\rho = M_W^2/M_Z^2 \cos^2\theta_W = \frac{1 + 2v_\Delta^2/v^2}{1 + 4v_\Delta^2/v^2}, \quad (16)$$

where $v = 246$ GeV is the v.e.v. of the SM doublet. Considering the experimental value $\rho^{\rm exp} = 1.00038(20)$ [4] one is lead to the upper limit $v_\Delta \lesssim \mathcal{O}(1)$ GeV [51, 53].

On the other hand, the mass of the doubly charged Higgs boson is constrained indirectly as a function of the product of leptonic Yukawa couplings from several processes [48–53], including Bhabha scattering, LFV violating transitions, muonic oscillation, and the electron and muon $(g-2)$ observables (see Table IV in Ref. [3]). With all these input parameters we can estimate upper bounds on the rates of $\Delta L = 2$ hyperon decays induced by a doubly charged Higgs boson. Let us consider the specific example of the $\Sigma^- \to p\ell\ell'$ decays for which the values $A = 3.56 \times 10^5$ MeV^3 and $B = 0$ have been calculated [2] using the MIT bag model (the non-relativistic approximation for moving hyperons was assumed in calculations).

In order to present our estimates, we consider for simplicity non τ-flavored interactions, that is, $h_{\tau i} = 0$ $(i = e, \mu, \tau)$. Furthermore, we adopt a conservative benchmark considering that $v_\Delta = 3$ GeV, and $h_{mm} \simeq 0.1$ $(m = e,\ \mu)$ for the rest of diagonal Yukawa couplings. If we now consider the limits from $\ell\ell \to \ell\ell$ $(\ell = e, \mu)$ data which only involve diagonal couplings h_{ee} and $h_{\mu\mu}$, then $m_{H^{\pm\pm}} \gtrsim 395$ GeV. Choosing the lowest value for $m_{H^{\pm\pm}}$, we obtain

$$\mathrm{BR}_{H^{--}}(\Sigma^- \to pee) = 1.1 \times 10^{-30},$$
$$\mathrm{BR}_{H^{--}}(\Sigma^- \to p\mu\mu) = 1.0 \times 10^{-31}. \quad (17)$$

For smaller Yukawa couplings h_{mm}, the above upper limit increases by a factor $1/h^2_{mm}$ if we still assume the lower bound on $m_{H^{\pm\pm}}$ from $ee \to \mu\mu$ data quoted in Table IV in Ref. [3]. Notice that, in this case, it is necessary to consider $h_{e\mu} \lesssim 3.5 \times 10^{-6}$ in order to obey the strong constraint coming from $\mu \to eee^+$. In such a case, we have

$$\mathrm{BR}_{H^{--}}(\Sigma^- \to pe\mu) = 1.3 \times 10^{-39}. \quad (18)$$

Finally, bounds for the other hyperon decays in this HTM model can be computed in a similar way if matrix elements of four-quark operators for other channels become available.

4. Conclusions

We have studied all the $\Delta L = 2$ decays of spin-1/2 hyperons $B_A^- \to B_B^+ \ell^- \ell'^-$ within a model involving a one loop mechanism with baryons and Majorana neutrinos as intermediate states. Our results improve previous estimates reported in [1] in several ways. We have included the momentum dependence of hyperon form factors using a monopolar model which allows to cure the bad ultraviolet behavior encountered in Ref. [1]. Furthermore, we kept finite values for the Majorana neutrino masses to cover the case of heavier neutrinos. In addition to the scenario with light Majorana neutrinos (*scenario A*), where we use current bounds available for the effective Majorana masses, we have also considered a minimal seesaw model that involves two heavier neutrinos (*scenario B*).

We applied our results to a specific low-scale seesaw model which includes two Majorana neutrinos with masses in the TeV range and not very suppressed mixings [41, 42], and three light active neutrinos. In this model, the lepton number violating effects are encoded in the mass splitting of the heavy neutrinos $(r \neq 1)$, while the heavy-light mixing angles are bounded from the perturbative unitarity condition. The predicted upper limits for the branching ratios in this *scenario B* turn out to be eleven orders of magnitude larger for the two electron channels although very far from expected sensitivities at BESIII.

A different mechanism for $\Delta L = 2$ in hyperon decays is also explored. We consider a Higgs Triplet Model, which can generate neutrino masses through the type-II seesaw mechanism, and contains a doubly charged scalar that couples to equal-sign leptons. Using current bounds on the parameters of the model relevant for $\Delta L = 2$ decays, we find the branching fraction for $\Sigma^- \to pee$ to be less suppressed than in *scenario A* discussed above, but still far below the current and expected sensitivities at BESIII.

It is important to mention that a more appropriate framework to deal with the contributions of heavy Majorana neutrinos ($m_\nu \sim$ few TeV) corresponds to integrate out the heavy states considering the perspective of an effective field theory. In this regard, the discussion about *scenario B* presented here in the loop mechanism where we kept the dependence on the neutrino masses turns out useful to see that the behavior of the hadronic form-factors has a strong dependence for neutrino states heavier than a few tens or hundreds MeVs (such as is depicted in Fig. 2). Note that if the validity of the loop mechanism is extended for states with masses around 1 GeV, this is an important point to take into account. Nevertheless, this is not the most appropriate framework to deal with the contributions of heavy neutrinos around/above the electroweak scale. Therefore, the results of the third column in Table III should be taken with care and just like a simple try of a smooth transition describing the contributions from light to heavy neutrino effects. The effects of heavy neutrinos to ΔL=2 hyperon decays from an effective field theory perspective will be presented in a new version of Ref. [3].

Acknowledgments

I wish to thank G. López-Castro and D. Portillo-Sánchez. This work has been done in collaboration with them. I also like to thank PROGRAMA DE BECAS POSDOCTORALES DGAPA-UNAM for economical support.

$i.$ The study of $\Delta L = 2$ processes in decays of tau leptons [5–10, 10–12], pseudoscalar mesons [13–29], and Λ_b baryons [30–33] is mainly motivated by the resonant effect produced by an intermediate Majorana neutrino and their study in flavor-factories experiments.

$ii.$ In order to avoid lengthy expressions, we omit here the expressions of the relevant $C_{v_r}^{nj}$ form factors in terms of Passarino-Veltman functions for the monopolar approximation. These expressions can be found in [3]. The one loop integration have been performed using the package *Package-X* [54] and evaluated numerically with *Collier* [55].

$iii.$ The behavior of the total axial $C_{a_r}^{\eta}$ form factors is very similar to the vector $C_{v_r}^{\eta}$ ones. Our numerical results include all the factors, although for $r = 1, 2, A$ they are sub-dominant.

$iv.$

$v.$

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