

# Electron electric dipole moment in Inverse Seesaw models

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**ABSTRACT:** We consider the contribution of sterile neutrinos to the electric dipole moment of charged leptons in the most minimal realisation of the Inverse Seesaw mechanism, in which the Standard Model is extended by two right-handed neutrinos and two sterile fermion states. Our study shows that the two pairs of (heavy) pseudo-Dirac mass eigenstates can give significant contributions to the electron electric dipole moment, lying close to future experimental sensitivity if their masses are above the electroweak scale. The major contribution comes from two-loop diagrams with pseudo-Dirac neutrino states running in the loops. In our analysis we further discuss the possibility of having a successful leptogenesis in this framework, compatible with a large electron electric dipole moment.

**KEYWORDS:** Beyond Standard Model, CP violation, Neutrino Physics

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## 1 Introduction

The origin of neutrino masses, the relic dark matter abundance and the baryon asymmetry of the Universe (BAU) are pressing open questions calling for extensions of the Standard Model (SM). One of the minimal extensions aiming to provide at least an explanation for the neutrino oscillation phenomena consists in the addition of right-handed (RH) neutrinos, singlets under the SM gauge group, giving rise to Dirac mass terms. The fact that the RH neutrinos are sterile fermions implies that they can have a Majorana mass so that this simple extension of the SM corresponds to the embedding of the seesaw mechanism [1–7] into the SM. The possibility of probing this scenario depends on the mass regime of the additional sterile states and on the size of their Yukawa couplings to the active neutrinos. In the “standard” type I seesaw, the masses of the RH neutrinos are required to be very large for sizeable (natural) values of the Yukawa couplings, implying that any direct (collider observables) or indirect signals (low-energy or high-intensity observables) are likely impossible to be discovered. When the masses of the additional RH states are around or below the electroweak scale, these states can be directly produced in colliders and their contribution to low-energy observables can be important: this is why low-scale seesaw models [8–14] prove to be appealing. Among them, the Inverse Seesaw mechanism [8], the  $\nu$ MSM [11], the low-scale type-I seesaw [12, 13] and the Linear Seesaw [15, 16] are examples of models with a rich phenomenology.

Some of the latter scenarios may also provide a possible explanation to the dark matter relic density considering, putting forward a keV-scale sterile neutrino as a viable candidate [11, 17], and/or to the BAU through leptogenesis (via neutrino oscillations) [18–23].

Other than the above motivation, some of the current neutrino oscillation experiments (reactor [24–26], accelerator [27–30] and Gallium [31, 32]) suggest the existence of sterile fermions with masses in the eV range. This would imply that instead of the three-neutrino mixing scheme (in oscillation phenomena), one would have a 3 + 1-neutrino (or 3+more) mixing schemes (see, for instance, [33]).

Extensions of the SM with sterile fermions, which accommodate oscillation data, may also have an impact on other observables, such as charged lepton flavour violation (cLFV) in Higgs [34–36], neutral  $Z$  boson [37–39] and meson decays [40–44]. Sterile fermions may also contribute to lepton flavour conserving observables such as the dipole moments of charged leptons [45–47]. Coupling to the active neutrinos, the sterile fermions may also add new sources of CP violation to the already existing one of the SM, and can thus provide new contributions to different CP-odd observables, among them electric dipole moments (EDMs) [45].

In a recent work [47], we studied the impact of sterile fermions on the electric dipole moments of charged leptons in the context of the SM extended by an arbitrary number<sup>1</sup> of sterile neutrinos — without necessarily invoking a mechanism of neutrino mass generation — and we have shown that in order to have a non-vanishing contribution, the minimal number of sterile fermion states that must be added to the SM field content is two. In particular, in the framework of this ad-hoc construction, the latter states can give significant contributions to the charged lepton EDMs, some of them lying within future experimental sensitivity if their (non-degenerate) masses are both above the electroweak scale. The Majorana nature of the neutrino states is also an important ingredient in order to allow for significative contributions to the charged lepton EDMs.

In this work, we consider the electron EDM in a specific seesaw realisation, the Inverse Seesaw (ISS) in its minimal version, which offers the possibility of accommodating the smallness of the light (mostly active) neutrino masses for a comparatively low seesaw scale, but still with large values of the Yukawa couplings.

In [14], it was shown that it is possible to construct several minimal ISS realisations that can reproduce the correct neutrino mass spectrum while fulfilling all phenomenological constraints, each exhibiting distinct features. This allowed to identify a truly minimal ISS realisation denoted “(2,2) ISS” model, where the SM is extended by two RH neutrinos and two sterile states. This configuration leads to a 3-flavour mixing scheme in the normal hierarchy for the light neutrinos, and two pairs of (heavy) pseudo-Dirac mass eigenstates.

Compared to the ad-hoc construction previously mentioned, in the (2,2) ISS realisation, the additional CP phases and the peculiar heavy spectrum (two pairs of pseudo-Dirac states) may lead to different prospects concerning the charged lepton EDMs.

The present study shows that the pseudo-Dirac states can give significant contributions to the electron EDM, close to the future experimental sensitivity, should the pseudo-Dirac

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<sup>1</sup>Being SM gauge singlets, there is no constraint on their (generation) number from anomaly cancellation.

masses be above the electroweak scale. We have shown that, contrary to the ad-hoc model, the two-loops diagrams relying on the Majorana nature of the exchanged neutrinos turn out to be suppressed. Interestingly, due to the structure of the spectrum and of the lepton mixings, in the (2,2) ISS, the major contribution to the EDMs arises from the diagrams with Dirac-like fermions in the loops. Furthermore, we have estimated the maximal enhancement factor to the charged lepton EDMs in the case of a generic  $(N, N)$  ISS realisation, with  $N > 2$ . In this work we also discuss the possibility of having a successful (thermal) leptogenesis in this ISS framework in the regimes associated with significant contributions to the electron EDM.

The paper is organised as follows: in section 2, after describing the (2,2) ISS realisation, we discuss in detail the degrees of freedom associated to the CP-violating phases, the mass regimes of the sterile states and the active-sterile mixing angles. We also summarise the relevant constraints on these extensions of the SM. Section 3 is devoted to the charged lepton EDMs, including a thorough discussion of the several two-loop diagrams. The impact of sterile neutrino contributions to the EDMs within this minimal ISS realisation is numerically evaluated and presented in section 4, while the analytical determination of the EDM is summarized in the appendix. Our final remarks and discussion are collected in section 5. We summarize our results in section 6.

## 2 The (2,2) ISS model

### 2.1 The neutrino mass matrix

As argued in [14], the minimal realisation of the Inverse Seesaw model requires the addition of two right-handed neutrinos  $N_i$ , and two singlet fermions  $s_i$  to the SM field content. Assigning the same lepton numbers ( $L = +1$ ) to  $N_i$  and  $s_i$  allows for a small  $\Delta L = 2$  lepton number violating (LNV) mass parameter  $\mu$  and  $m$ , corresponding to Majorana masses in the sterile sector. This leads to the following neutrino mass terms in the Lagrangian

$$-\mathcal{L}_{m_\nu} = n_L^T C M n_L + \text{h.c.}, \tag{2.1}$$

where

$$n_L \equiv (\nu_L^1, \nu_L^2, \nu_L^3, N_1^c, N_2^c, s_1, s_2)^T, \text{ and } C = i\gamma^2\gamma^0. \tag{2.2}$$

The mass matrix  $\mathcal{M}$  can be written as:

$$M = \begin{pmatrix} 0 & 0 & 0 & d_{11} & d_{12} & 0 & 0 \\ 0 & 0 & 0 & d_{21} & d_{22} & 0 & 0 \\ 0 & 0 & 0 & d_{31} & d_{32} & 0 & 0 \\ d_{11} & d_{21} & d_{31} & m_{11} & m_{12} & n_{11} & n_{12} \\ d_{12} & d_{22} & d_{32} & m_{12} & m_{22} & n_{21} & n_{22} \\ 0 & 0 & 0 & n_{11} & n_{21} & \mu_{11} & \mu_{12} \\ 0 & 0 & 0 & n_{12} & n_{22} & \mu_{12} & \mu_{22} \end{pmatrix}. \tag{2.3}$$

Matrix	# of moduli	# of phases	Total
Diagonal and real $\mathbf{m}_\ell$	3	0	3
$d$ with one real column	6	3	9
$m$	3	3	6
Real and diagonal $n$	2	0	2
$\mu$ with real diagonal	3	1	4
Total	17	7	24

**Table 1.** Example of a basis in which the number of parameters matches the number of physical degrees of freedom ( $\mathbf{m}_\ell$  corresponds to the charged lepton masses).

One can always choose a basis in which the unphysical parameters are reabsorbed via appropriate redefinitions of the fields; one of the possible choices of basis — which we have adopted [14] — is summarised in table 1, leading to 24 physical parameters.

In the following we will use this theoretical framework, denoted “(2,2) ISS” model. Moreover, we will neglect the mass parameters  $m_{ij}$ , which induce subdominant effects when compared to the entries  $\mu_{ij}$  in the mass matrix of eq. (2.3). Indeed, both these lepton number violating mass matrices can be dynamically generated as done in the general original formulation of the Inverse Seesaw mechanism [8], in which the smallness of the  $\mu$  matrix elements was attributed to supersymmetry breaking effects in a (superstring inspired)  $E_6$  scenario. For instance, in the context of a non-supersymmetric SO(10) model, which contains the remnants of a larger  $E_6$  group, the mass matrix  $\mu$  is generated at two-loop level while the matrix  $m$  is generated at higher order, thus justifying the smallness of its entries compared to the ones of  $\mu$  [48]. Once the lepton number violating terms  $m_{ij}$  are neglected, the mass matrix is reduced to

$$M = \begin{pmatrix} 0 & 0 & 0 & d_{11} & d_{12} & 0 & 0 \\ 0 & 0 & 0 & d_{21} & d_{22} & 0 & 0 \\ 0 & 0 & 0 & d_{31} & d_{32} & 0 & 0 \\ d_{11} & d_{21} & d_{31} & 0 & 0 & n_1 & 0 \\ d_{12} & d_{22} & d_{32} & 0 & 0 & 0 & n_2 \\ 0 & 0 & 0 & n_1 & 0 & \mu_{11} & \mu_{12} \\ 0 & 0 & 0 & 0 & n_2 & \mu_{21} & \mu_{22} \end{pmatrix}, \tag{2.4}$$

where the sub-matrix  $d$  is parametrised by 6 moduli and 3 CP phases,  $n$  via 2 moduli, while  $\mu$  includes 3 moduli and 1 CP phase. Thus the mass matrix  $M$  totally is defined by 11 moduli and 4 CP phases. Since the determinant of the mass matrix in eq. (2.4) vanishes, the lightest mass eigenvalue is zero in the minimal Inverse Seesaw model. The diagonalisation of the mass matrix in eq. (2.4) leads at leading order to three light (almost active) neutrinos (systematically in the normal hierarchy ordering, as found in [14]), and to two pseudo-Dirac pairs containing the mostly sterile eigenstates, with mass differences of the order of the LNV entries of the  $\mu$  sub-matrix.<sup>2</sup>

<sup>2</sup>In the limit in which lepton number is conserved (i.e.  $\mu \rightarrow 0$ ), these states become Dirac particles.

The weak charged current Lagrangian for the leptons is modified as

$$- \mathcal{L}_{cc} = \frac{g}{\sqrt{2}} U_{\alpha i} \bar{\ell}_\alpha \gamma^\mu P_L \nu_i W_\mu^- + \text{h.c.}, \quad (2.5)$$

where  $U_{\alpha i}$  is the unitary lepton mixing matrix,  $i = 1, \dots, 7$  denotes the physical neutrino states and  $\alpha = e, \mu, \tau$  the flavour of the charged leptons. In the case of three neutrino generations,  $U$  would correspond to the ( $3 \times 3$  unitary) PMNS matrix,  $U_{\text{PMNS}}$ . The mixing between the left-handed leptons, here denoted by  $\tilde{U}_{\text{PMNS}}$ , now corresponds to a  $3 \times 3$  block of the  $7 \times 7$  unitary matrix  $U$ , which can be parametrised as

$$U_{\text{PMNS}} \rightarrow \tilde{U}_{\text{PMNS}} = (\mathbb{1} - \eta) U_{\text{PMNS}}, \quad (2.6)$$

where the matrix  $\eta$  encodes the deviation of  $\tilde{U}_{\text{PMNS}}$  from unitarity [6, 49]. It is also convenient to introduce the invariant quantity  $\tilde{\eta} = 1 - |\text{Det}(\tilde{U}_{\text{PMNS}})|$ , particularly useful to illustrate the effect of the active-sterile mixings.

## 2.2 Constraints

Depending on their masses and on the mixings with the active (light) neutrinos, the sterile states are severely constrained from several observations. Firstly, these extensions should account for oscillation data. In our analysis we have required compatibility with the best fit intervals for a normal hierarchical light spectrum [50],

$$0.270 \leq \sin^2 \theta_{12} \leq 0.344, \quad 0.382 \leq \sin^2 \theta_{23} \leq 0.643, \quad 0.0186 \leq \sin^2 \theta_{13} \leq 0.0250, \quad (2.7)$$

$$7.02 \leq \frac{\Delta m_{21}^2}{10^{-5} \text{ eV}^2} \leq 8.09, \quad 2.317 \leq \frac{\Delta m_{31}^2}{10^{-3} \text{ eV}^2} \leq 2.607. \quad (2.8)$$

(As discussed in [14], for such a minimal realisation of the Inverse Seesaw model, an inverted hierarchy is strongly disfavoured.) We also notice that the non-unitarity of the  $\tilde{U}_{\text{PMNS}}$  (sub-matrix) is constrained by a number of observations, as discussed in [51, 52].

Further constraints on the active-sterile mixings and on the sterile neutrino masses can be inferred from current bounds arising from neutrinoless double beta decays [53]. In the present scenario, the relevant effective neutrino mass is given by [46, 54]

$$m_{ee} = \sum_{i=1}^7 U_{ei}^2 \frac{p^2 m_i}{p^2 - m_i^2}, \quad (2.9)$$

where  $p^2 = -(125 \text{ MeV})^2$ .

Several experiments (like GERDA [55], EXO-200 [56, 57], and KamLAND-ZEN [58]) have put constraints on  $|m_{ee}|$ , which translate into bounds on combinations of  $U_{ei}^2 m_i$ ,  $i = 4, \dots, 7$ . In our numerical study, we have imposed that our solutions always comply with the (conservative) experimental bound  $|m_{ee}| \lesssim 0.01 \text{ eV}$ .

Particularly relevant for the case of a large sterile mass regime is the perturbative unitarity bound: if the additional sterile fermions are sufficiently heavy to decay into a  $W$  boson and a charged lepton, or into an active neutrino and either a  $Z$  or a Higgs boson, their decay widths should comply with the perturbative unitarity condition [38, 59–64]. In

this case, and since the dominant decay mode of the (mostly) sterile neutrinos,  $\nu_{4\dots 7}$ , would be  $\nu_i \rightarrow \ell_\alpha^\mp W^\pm$ , their decay width should comply with the perturbative unitary bound:<sup>3</sup>

$$\frac{\Gamma_{\nu_i}}{m_i} < \frac{1}{2} \quad \text{where} \quad \Gamma_{\nu_i} \approx \frac{g_2^2 m_i^3}{16\pi m_W^2} \sum_\alpha |U_{\alpha i}|^2, \quad (i = 4, \dots, 7), \quad (2.10)$$

which translates into an upper bound on the sterile neutrino masses as follows,

$$m_i \lesssim 873 \text{ GeV} \left( \sum_\alpha |U_{\alpha i}|^2 \right)^{-1/2}. \quad (2.11)$$

Important bounds arise from electroweak precision tests; the active-sterile mixings are constrained from observables such as the  $W$  boson decay width, the  $Z$  invisible decay, meson decays and the non-unitarity of the  $3 \times 3$  sub-matrix ( $\tilde{U}_{\text{PMNS}}$ ) of  $U_{ij}$ . For  $m_i < m_W, m_Z$  ( $i = 4 - 7$ ), the most important constraints arise from the  $W$  decay and the  $Z$  invisible decay; for the neutrino mass range  $3 \text{ GeV} \lesssim m_i \lesssim 90 \text{ GeV}$ , the strongest constraints are those of the DELPHI [65] and L3 [66] Collaboration.

Additional sterile states might also lead to the violation of lepton flavour universality, as arising from meson decays such as  $\pi^+ \rightarrow \ell_\alpha^+ \nu_\alpha$  and  $K^+ \rightarrow \ell_\alpha^+ \nu_\alpha$  [43, 44, 67]. Lepton flavour violating processes also provide important constraints on the sterile fermion parameter space; in particular, the bounds from the muon-electron sector lead to the strongest constraints, which are

$$\text{Br}(\mu \rightarrow e\gamma) \leq 5.7 \times 10^{-13}, \quad \text{Br}(\mu \rightarrow e\bar{e}e) \leq 1.0 \times 10^{-12}, \quad \text{Cr}(\mu - e, \text{Au}) \leq 7 \times 10^{-13}, \quad (2.12)$$

as obtained, respectively, by [68, 69] and [70].

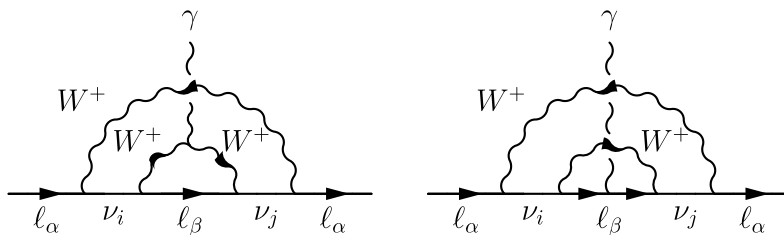
Direct searches at LEP have put strong constraints on sterile neutrinos whose masses are  $m_i \lesssim \mathcal{O}(100) \text{ GeV}$ . The relevant process is  $e^+e^- \rightarrow \nu_i \nu_j^* \rightarrow \nu_i e^\pm W^\mp$  where  $i \leq 3$  and  $j \geq 4$ , which violates lepton number conservation. This has allowed to exclude certain regimes of the mixing angles  $|U_{\alpha i}|$  [35]. Searches at the LHC for a same sign di-lepton channel  $pp \rightarrow W^{\pm*} \rightarrow \ell^\pm \nu_i \rightarrow \ell^\pm \ell^\pm jj$  (where  $i \geq 4$  and  $j$  denotes a jet), have led to further bounds for  $m_i \gtrsim \mathcal{O}(100 \text{ GeV})$ : for values of the integrated luminosity of  $20 \text{ fb}^{-1}$  at  $\sqrt{s} = 8 \text{ TeV}$ , LHC data already allows to constrain the mixing angle  $|U_{\alpha i}|$  for sterile neutrino masses up to  $500 \text{ GeV}$  [71, 72].

### 3 Electric Dipole Moments

In Inverse Seesaw models, the charged lepton EDMs are induced at two-loop level as shown in refs. [45, 47]. In general, there is a very large number ( $\sim 100$ ) of diagrams contributing to the charged lepton EDM;<sup>4</sup> however and noticing that the heavy (sterile) neutrinos providing the dominant contributions form pseudo-Dirac pairs, the number of diagrams that must be

<sup>3</sup>Another common criterion of perturbativity is that the couplings should be less than  $\sqrt{4\pi}$ . This criterion also gives a bound similar to eq. (2.11).

<sup>4</sup>The number of relevant diagrams to EDMs depends on the chosen gauge. If a non-linear gauge is taken, the number of diagrams is considerably reduced [73, 74].



**Figure 1.** Diagrams for charged lepton EDMs.

evaluated can be significantly reduced, as we will presently discuss. In a previous work [47], we have discussed the several possible contributing diagrams in the case of an ad-hoc model corresponding to the Standard Model extended by  $N$  sterile states, considering both Majorana and Dirac fermion contributions. There was, in the Feynman gauge, in total 44 diagrams of type of figure 1, as well as 96 additional diagrams corresponding to  $Z$  and to the Higgs bosons mediation. In this previous study we have computed the diagrams and discussed the relevance of each contribution, distinguishing also the Dirac contribution from the Majorana one. It is worth noticing that in the Inverse seesaw framework, the number of diagrams that must be evaluated is significantly reduced, as we will presently discuss.

In “effective” neutrino models (corresponding to minimalistic ad-hoc constructions where the SM is extended by  $N$  sterile fermions), the EDM of the electron can be formally written as [47]

$$d_e = -\frac{g_2^4 e m_e}{4(4\pi)^2 m_W^2} \sum_{\beta} \sum_{i,j} \left[ J_{ije\beta}^M I_M(x_i, x_j) + J_{ije\beta}^D I_D(x_i, x_j) \right], \quad (3.1)$$

where  $x_{i,j} \equiv m_{i,j}^2/m_W^2$  and  $m_{\alpha,\beta}^2/m_W^2 \ll 1$  ( $\alpha, \beta = e, \mu, \tau$ ),  $I_M$  and  $I_D$  are the loop functions whose analytic expressions (for the dominant Majorana contributions) can be found in [47]. The (CP-odd) factors  $J_{ij\alpha\beta}^M$  and  $J_{ij\alpha\beta}^D$  are defined by

$$J_{ij\alpha\beta}^M \equiv \text{Im} (U_{\alpha j} U_{\beta j} U_{\beta i}^* U_{\alpha i}^*), \quad J_{ij\alpha\beta}^D \equiv \text{Im} (U_{\alpha j} U_{\beta j}^* U_{\beta i} U_{\alpha i}^*). \quad (3.2)$$

Similar expressions hold for the EDMs of the other charged leptons ( $\mu$  and  $\tau$ ); however we do not consider them here since the predicted EDMs for  $\mu$  and  $\tau$  are extremely tiny compared to the sensitivities of future experiments. The first term in the brackets in eq. (3.1) represents a contribution which reflects the Majorana nature of the sterile fermions, while the second term corresponds to a generic Dirac fermion contribution. As one can see from the definition of the factors  $J_{ij\alpha\beta}^M$  and  $J_{ij\alpha\beta}^D$ , these are totally anti-symmetric in terms of  $i \leftrightarrow j$ , implying that a non-vanishing EDM requires contributions from two neutrino states with  $i \neq j$ . The loop functions  $I_M(x_i, x_j)$  and  $I_D(x_i, x_j)$  should also be anti-symmetric under  $i \leftrightarrow j$ . As a result, one can see that the EDM itself is fully symmetric under  $i \leftrightarrow j$ .

Such a formulation is valid for any extension of the SM involving sterile fermions, as is the case of type-I, Inverse and Linear Seesaw models. In Inverse Seesaw models (with equal number of  $N_i$  and  $s_i$ ), the heavy sterile neutrinos form pseudo-Dirac fermion pairs. Moreover, within a pseudo-Dirac pair, the states are highly degenerate in mass (recall



that their non-degeneracy is proportional to the entries of the  $\mu$  matrix). The Majorana contribution in eq. (3.1) is thus very suppressed when compared to the Dirac one.<sup>5</sup> Taking into account this fact, the expression of the electron EDM for the case of Inverse Seesaw models is simplified to

$$d_e \approx -\frac{g_2^4 e m_e}{4(4\pi)^2 m_W^2} \sum_{\beta} \sum_{i,j} J_{ije\beta}^D I_D(x_i, x_j). \quad (3.3)$$

In addition, in the case of the “(2,2)” minimal Inverse Seesaw model, the two pairs of heavy sterile neutrinos,  $(\nu_4, \nu_5)$  and  $(\nu_6, \nu_7)$ , are nearly degenerate respectively, with  $m_{4,5} < m_{6,7}$ . Taking into account this fact and unitarity of the mixing matrix ( $UU^\dagger = \mathbb{1}$ ), the formula of the electron EDM can be further simplified by

$$d_e \approx -\frac{g_2^4 e m_e}{2(4\pi)^2 m_W^2} J^D I'_D(x_4, x_6), \quad (3.4)$$

where the loop function  $I'_D$  and phase factor  $J^D$  are defined by

$$I'_D(x_4, x_6) \equiv I_D(0, x_4) - I_D(0, x_6) + I_D(x_4, x_6), \quad (3.5)$$

$$J^D \equiv \sum_{\beta} \left[ J_{46e\beta}^D + J_{47e\beta}^D + J_{56e\beta}^D + J_{57e\beta}^D \right]; \quad (3.6)$$

and the factor 2 difference between eq. (3.3) and (3.4) arises from having the EDM expression totally symmetric under  $i \leftrightarrow j$ .

The electron EDM has been experimentally searched for by ACME Collaboration [75], and the current upper bound is given by

$$|d_e|/e \leq 8.7 \times 10^{-29} \text{ cm}. \quad (3.7)$$

The comparison of the above bounds with eq. (3.4) allows to set the limit  $|J^D I'_D(x_4, x_6)| \lesssim 9.7 \times 10^{-5}$ . The upper bound is expected to be improved to  $|d_e|/e \lesssim 10^{-30}$  cm by the upgraded ACME Collaboration [76].

## 4 Numerical results

The most difficult part of the computation of the charged lepton EDMs is the evaluation of the loop function  $I_D(x_i, x_j)$  in eq. (3.3). This can be done with FeynCalc [77] and the analytical expressions that we have obtained are given in the appendix, where the derived analytical formulas are written by multiple integrals. The loop function  $I'_D$  of eq. (3.5) has been numerically evaluated, and some illustrative examples have been collected in the left panel of figure 2, in which we display  $I'_D$  as a function of  $m_4$  (nearly degenerate with  $m_5$ ,  $m_5 \approx m_4$ ) for several fixed values of  $m_6$  ( $m_7 \approx m_6$ ). Recall that this degeneracy is a consequence of the pseudo-Dirac nature of the heavy spectrum. The sign of the loop

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<sup>5</sup>We have numerically confirmed that the phase factor for the Majorana contribution  $J_{ij\alpha\beta}^M$  is indeed highly suppressed as shown in figure 2 in the next section.

function changes at  $m_4 = m_6$ , corresponding to the “singularities” in the absolute value of  $I'_D$ , as can be seen on the figure, and the loop function becomes approximately flat for regimes where  $m_4 \gg m_6$ , i.e., for a strongly hierarchical heavy spectrum. Although the loop function could be larger for heavier sterile neutrinos, these regimes are theoretically and experimentally constrained, in particular by the perturbative unitarity bound and by constraints arising from cLFV observables.

A full analysis requires to carry a numerical diagonalisation of the  $7 \times 7$  neutrino mass matrix of eq. (2.4). In order to account for neutrino oscillation data, the Inverse Seesaw mechanism parameters must fulfill the following condition

$$|\mu| \ll |d| \ll |n|, \tag{4.1}$$

where  $\mu$ ,  $d$  and  $n$  are the elements of the sub-matrices in eq. (2.4). In addition, since the light neutrino mass matrix can be approximately given by

$$m_\nu^{\text{light}} \simeq d (n^{-1})^T \mu n^{-1} d^T, \tag{4.2}$$

the lepton number violating parameter  $\mu$  would be given by  $\mu \sim m_\nu^{\text{light}} n^2/d^2$ . The ratio  $d/n$  is related with the non-unitarity of the  $\tilde{U}_{\text{PMNS}}$  matrix and is thus experimentally constrained. To satisfy the latter constraints, the ratio should obey  $d/n \lesssim 0.1$  [14]. Accordingly, we take the following intervals of the different entries of the neutrino mass matrix,

$$1 \text{ GeV} \leq n_i \leq 10^7 \text{ GeV}, \quad 10^{-3} \leq \frac{|d_{ij}|}{\max[n_i]} \leq 10^{-1}, \tag{4.3}$$

$$\frac{\min[n_i^2]}{\max[|d_{ij}|^2]} \leq \frac{\mu_{11}}{10^{-10} \text{ GeV}} \leq \frac{\max[n_i^2]}{\min[|d_{ij}|^2]}, \quad 10^{-6} \leq \left| \frac{\mu_{12}}{\mu_{11}} \right| \leq 10^{-3}, \quad 10^{-1} \leq \left| \frac{\mu_{22}}{\mu_{11}} \right| \leq 10. \tag{4.4}$$

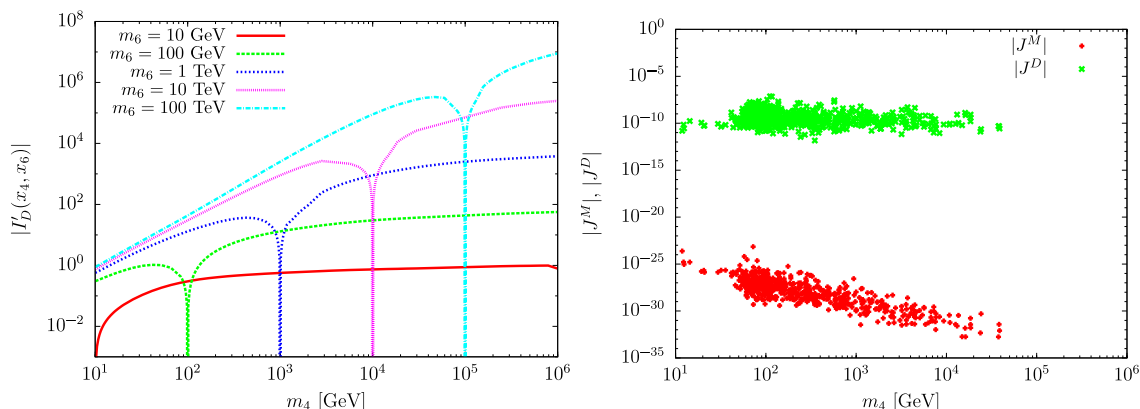
All the parameters in the mass matrix are randomly taken in the above ranges, and the different CP phases are also randomly varied in the  $[0, 2\pi]$  interval.

We display on the right panel of figure 2 the phase factors  $|J^D|$  and  $|J^M|$  computed using the data points complying with all the constraints discussed in section 2.2. One can observe from the figure that the Dirac contribution,  $|J^D|$ , is dominant as expected from our previous discussion; this justifies having the dominant contribution arising from the diagrams displayed in figure 1 (corresponding to the contribution of generic Dirac neutrinos in the loops). The maximum value of the factor  $|J^D|$  is approximately given by

$$|J_{\text{max}}^D| \sim 10^{-5} \times \left( \frac{\text{GeV}}{m_4} \right). \tag{4.5}$$

Combining the above quantities, one can compute the electron EDM, and the results are shown in figure 3 as a function of  $m_i$  (left) and  $\tilde{\eta}$  (right), where  $\tilde{\eta} = 1 - |\text{Det}(\tilde{U}_{\text{PMNS}})|$ . As can be seen, the maximum value of the predicted electron EDM is  $|d_e^{\text{max}}|/e \sim 5 \times 10^{-31}$  cm, lying two orders of magnitude below the current experimental bound,  $|d_e|/e \leq 8.7 \times 10^{-29}$  cm, and thus marginally short of the future sensitivity,  $|d_e|/e \sim 10^{-30}$  cm.

Contrary to the previous study [47], where an ad-hoc “3 + 2 toy” construction was used (two sterile fermions added to the SM field content), the present framework (the



**Figure 2.** On the left, loop function  $|I'_D(x_4, x_6)|$  for several fixed values of  $m_6$ ; on the right, CP-odd factors  $|J^M|$  and  $|J^D|$  corresponding to points in parameter space satisfying all the constraints discussed in section 2.

(2,2) ISS model) leads to a much more constrained scenario: firstly, the seesaw condition of eq. (4.1) strongly constrains the different couplings; secondly, the very nature of the heavy spectrum (pseudo-Dirac pairs) reduces the set of contributing diagrams. Finally, we stress that experimental constraints, as is the case of  $\mu \rightarrow e\gamma$ , are very severe in the mass regime where the most important contributions to the EDMs are expected to arise (above the electroweak scale). We do not dismiss the possibility that the (2,2) ISS model could eventually account for larger values of the electron EDM, but this would require an important amount of fine-tuning between the relevant parameters — and in the present study we chose not to explore such fine-tuned scenarios.

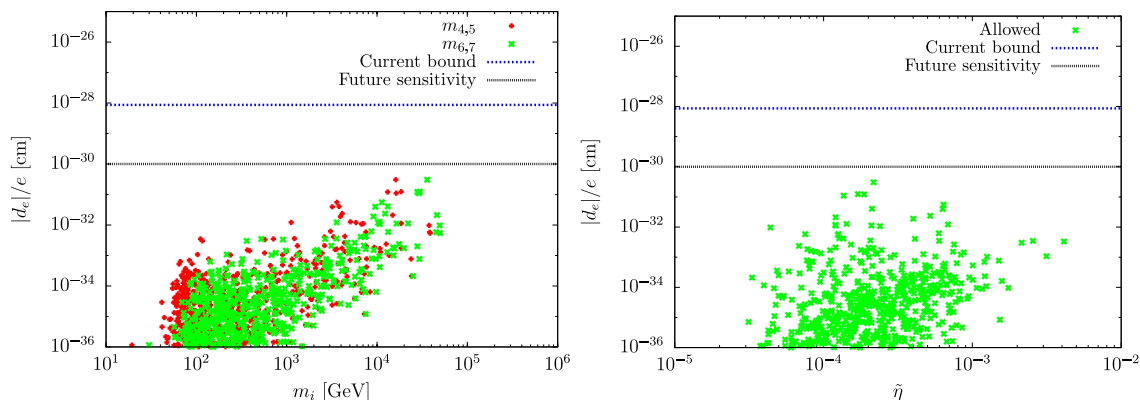
We have also computed the charged lepton anomalous magnetic moment which is given by

$$\Delta a_\ell = -\frac{4\sqrt{2}G_F m_\ell^2}{(4\pi)^2} \sum_{i=4}^7 |U_{\ell i}|^2 G_\gamma \left( \frac{m_i^2}{m_W^2} \right), \quad (4.6)$$

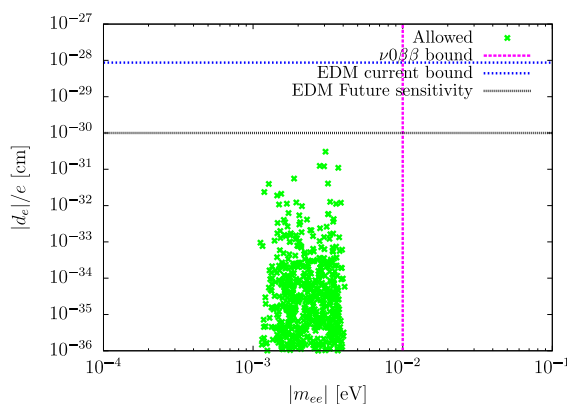
where the loop function  $G_\gamma(x)$  is given by

$$G_\gamma(x) = \frac{x - 6x^2 + 3x^3 + 2x^4 - 6x^3 \log x}{4(1-x)^4}. \quad (4.7)$$

For the muon anomalous magnetic moment, the new contribution is negative and cannot explain the discrepancy of  $3.5\sigma$  deviation between the SM prediction and the experimental value given by  $\Delta a_\mu = 2.88 \times 10^{-9}$  [78]. For the electron anomalous magnetic moment, we can naïvely expect from  $\Delta a_\mu$  — assuming a common New Physics origin to both discrepancies - that the deviation between the theoretical prediction of the SM and the corresponding experimental value will be of the order of  $\Delta a_e \sim 6.7 \times 10^{-14}$ , by scaling the value for the muon by  $m_e^2/m_\mu^2$ . However, the current difference between experimental observation and the SM prediction is  $\Delta a_e = 8.2 \times 10^{-13}$  [79]. Hence, this observable has the potential to probe and constrain New Physics contributions to  $\Delta a_e$  of order  $10^{-13}$ . Nevertheless, our analysis has shown that the present (2,2) ISS framework leads to contributions to the



**Figure 3.** Electron EDM as a function of  $m_i$  (left) and  $\tilde{\eta}$  (right), all the points displayed comply with the constraints discussed in section 2.2.



**Figure 4.** Electron EDM vs the effective neutrino mass  $|m_{ee}|$ . All the points displayed comply with the constraints discussed in section 2.2. The violet line denotes the conservative limit  $|m_{ee}| \leq 10^{-2}$  eV for  $0\nu\beta\beta$  decay.

electron anomalous magnetic moment of the order of  $|\Delta a_e| \sim 10^{-17}$ , which are too small compared to the present value.

Finally, we have considered the contribution of the present model to the effective mass  $m_{ee}$  (for the neutrinoless double beta decay). In the present scenario (characterised by a normal hierarchy of the light neutrino spectrum) we have found that the effective neutrino mass lies in the range  $1 \lesssim |m_{ee}| / (10^{-3} \text{ eV}) \lesssim 4$ , below the conservative experimental limit  $|m_{ee}| \leq 10^{-2}$  eV (see section 2.2), as can be seen in figure 4.

## 5 Discussion

### 5.1 Adding more sterile fermions

In the above analysis, we considered the Inverse Seesaw mechanism for neutrino mass generation with the minimal number of RH and sterile states. If more sterile fermions are added, one could expect an enhancement of the electron EDM induced by the extra neutrinos in the loop. In order to have an estimation of the possible increase we consider

the case in which the SM is extended by  $N$  right-handed neutrinos and  $N$  sterile fermions.<sup>6</sup> Assuming that all the mixing matrix elements  $U_{\alpha i}$  are of the same order, i.e.  $U_{\alpha i} \simeq \mathcal{O}(U_{\text{av}})$ , and that the loop function  $I'_D$  can be approximated by a constant (cf. left panel of figure 2, in the regime of a strong hierarchy between the pseudo-Dirac pairs), one can rewrite the EDM expression of eq. (3.4) for an  $(N, N)$  ISS realisation as

$$|d_e^{(N,N)}| \sim \frac{g_2^4 e m_e}{2(4\pi)^2 m_W^2} \left| 4N(N-1) \text{Im}(U_{\text{av}}^4) I'_D \right|. \quad (5.1)$$

On the other hand, the new contributions are subject to the constraints discussed on section 2.2, among them  $\mu \rightarrow e\gamma$ . Under the above assumptions, the branching ratio for  $\mu \rightarrow e\gamma$  increases, being approximately given by

$$\text{Br}(\mu \rightarrow e\gamma) \sim \frac{\sqrt{2} G_F^2 m_\mu^5}{\Gamma_\mu} (2N)^2 |U_{\text{av}}|^4 \leq 5.7 \times 10^{-13}. \quad (5.2)$$

Combining eqs. (5.1) and (5.2), the current experimental limit on the above branching ratio leads to an upper bound for the electron EDM obtained in or the  $(N, N)$  Inverse Seesaw realisation given by

$$|d_e^{(N,N)}| \lesssim (5.7 \times 10^{-13}) \frac{e m_e m_W^2 \Gamma_\mu}{4\sqrt{2}\pi^2 m_\mu^5} \left(1 - \frac{1}{N}\right) \sin\theta |I'_D|, \quad (5.3)$$

where  $\sin\theta$  is the phase defined by  $\text{Im}(U_{\text{av}}^4) \equiv |U_{\text{av}}|^4 \sin\theta$ .

When compared with the (2,2) ISS case, the  $(N, N)$  ISS contribution to the electron EDM is given by

$$\left| \frac{d_e^{(N,N)}}{d_e^{(2,2)}} \right| \lesssim 2 \left(1 - \frac{1}{N}\right). \quad (5.4)$$

For example, if we consider the  $N = 3$  case, a factor 4/3 enhancement is expected compared to the results found for the (2,2) ISS realisation. In the limit of  $N \gg 1$ , one obtains, at most, a factor 2 enhancement, which would nevertheless lead to contributions  $\sim 10^{-30}$  cm, thus within the optimistic sensitivity of the future experiment by ACME Collaboration.

## 5.2 Resonant leptogenesis

The baryon asymmetry of the Universe is another observation pointing towards New Physics scenarios. The current center value of the baryon asymmetry, as determined by the PLANCK Collaboration is given by [80]

$$\frac{n_b - n_{\bar{b}}}{s} = 8.59 \times 10^{-11}, \quad (5.5)$$

---

<sup>6</sup>As shown in [14], the next-to-minimal ISS realisation involving 2 RH neutrinos and 3 sterile fermions leads to the addition of a mass state at the LNV scale,  $\mathcal{O}(\mu)$ , to the spectrum obtained in the (2,2) ISS. The latter state is too light to provide an enhancement to the charged lepton EDMs. This statement can be generalised for a number of sterile states larger than the number of RH neutrinos: one finds a spectrum composed by pseudo-Dirac pairs and lighter (mostly) sterile states with masses  $\mathcal{O}(\mu)$ .

where  $n_b - n_{\bar{b}}$  is the asymmetry between the number density of baryons and anti-baryons, and  $s$  is the entropy density. In Inverse Seesaw realisations, it might be possible to generate the BAU via thermal leptogenesis. A viable leptogenesis is possible in the presence of out-of-equilibrium processes violating CP and lepton number. In the following, we discuss whether or not the regimes leading to a large electron EDM can also offer the interesting possibility of generating the BAU via leptogenesis.

Despite being a low-scale seesaw realisation, a successful baryogenesis could be achieved in the ISS via resonant leptogenesis [81] since the spectrum contains pseudo-Dirac pairs of nearly degenerate states. Even if lepton number violation and CP violation are present — in particular a significant amount of CP violation potentially leading to a large electron EDM —, it turns out that it is not possible to have a successful leptogenesis as one cannot satisfy the out-of-equilibrium condition. The decay width of the lightest heavy sterile neutrino  $\nu_4$  (whose decays are assumed to be responsible for the generation of a lepton asymmetry) obtained for its dominant channel  $\nu_4 \rightarrow \ell_{\alpha}^{\pm} W^{\mp}$  is given by

$$\Gamma_{\nu_4} = \frac{g_2^2 m_4^3}{16\pi m_W^2} \sum_{\alpha} |U_{\alpha 4}|^2. \tag{5.6}$$

The out-of-equilibrium condition corresponds to having a decay width smaller than the expansion rate of the Universe, which translates into the following inequality,

$$\Gamma_{\nu_4} < H(T)|_{T=m_4}, \tag{5.7}$$

where  $H$  is the Hubble parameter given by  $H \approx 1.66\sqrt{g_*}T^2/m_{\text{Pl}}$  for the radiation-dominated epoch, with  $g_*$  the effective degrees of freedom of relativistic particles,  $T$  the temperature of the universe and the Planck mass given by  $m_{\text{Pl}} = 1.22 \times 10^{19}$  GeV. From eqs. (5.6) and (5.7), one can obtain the required order of magnitude for the mixing angle  $|U_{\alpha 4}|$ ,

$$\sum_{\alpha} |U_{\alpha 4}|^2 \sim 10^{-15} \left( \frac{1 \text{ TeV}}{m_4} \right) \left( \frac{g_*}{100} \right)^{1/2}. \tag{5.8}$$

Thus very small mixing angles are needed to satisfy the out-of-equilibrium condition. This means that the amount of baryon asymmetry generated by resonant leptogenesis becomes quite small if large mixings (maximally  $|U_{\alpha i}| \sim 10^{-3}$ ) are assumed, as required to induce a large electron EDM (see section 4). If additional sterile states are added, it might be possible to have a successful BAU consistent with a large electron EDM since one of the pseudo-Dirac pairs of the sterile states is responsible for generating the BAU, while the other pairs can account for large contributions to the electron EDM.

## 6 Summary

We have considered the contribution of sterile neutrinos to the electric dipole moment of charged leptons in the (2,2) realisation of the Inverse Seesaw mechanism. We have shown that the two pairs of (heavy) pseudo-Dirac mass eigenstates can give significant contributions to the electron EDM, close to the future experimental sensitivity of ACME,

if their masses are above the electroweak scale. We have further investigated whether or not a successful leptogenesis can be accommodated in this framework, compatible with a large electron EDM. Although such a possibility is precluded for the minimal (2,2) ISS realisation, we do not dismiss its viability in an ISS framework with an extended spectrum ( $N, N$ ) ISS, where  $N > 2$ .

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## A Loop functions

The analytic expression of the loop function  $I_D(x_i, x_j)$  has been derived by FeynCalc in the limit of  $x_{\alpha,\beta} \ll 1$  and  $x_{i,j} \gg 1$ . The loop function represents the contributions of the two diagrams displayed in figure 1. Furthermore, each contribution of the diagrams of figure 1 is decomposed into two integral pieces. As a result, the loop function  $I_D(x_i, x_j)$  can be written as

$$I_D(x_i, x_j) = I_{D1}^L(x_i, x_j) + I_{D2}^L(x_i, x_j) + I_{D1}^R(x_i, x_j) + I_{D2}^R(x_i, x_j). \quad (\text{A.1})$$

The superscripts  $L, R$  correspond to the contribution of the left and right diagrams of figure 1, while the indices 1 and 2 refer to two different types of integrals as detailed below. Finally it is convenient to anti-symmetrize the loop function (as discussed in section 3) in terms of  $x_i$  and  $x_j$  as

$$I_D(x_i, x_j) \rightarrow \frac{1}{2} (I_D(x_i, x_j) - I_D(x_j, x_i)), \quad (\text{A.2})$$

since only the anti-symmetric part of the loop function contributes to the charged lepton EDM. The terms in the right hand side in eq. (A.1) are given by the integral in terms of the Feynman parameters  $s_A, t_B$  as

$$I_{Dn}^{L/R}(x_i, x_j) = \int_0^1 \prod_{A=1}^5 ds_A \delta\left(\sum_{A=1}^5 s_A - 1\right) \int_0^1 \prod_{B=1}^{3+\delta_{2n}} dt_B \delta\left(\sum_{B=1}^{3+\delta_{2n}} t_B - 1\right) F_n^{L/R}(x_i, x_j), \quad (\text{A.3})$$

where the integrands  $F_n^{L/R}(x_i, x_j)$  are given by

$$F_n^{L/R}(x_i, x_j) = \frac{N_n^{L/R}(x_i, x_j)}{D^{L/R}(x_i, x_j)}. \quad (\text{A.4})$$

The denominators  $D^L(x_i, x_j), D^R(x_i, x_j)$  are given by

$$D^L(x_i, x_j) = t_1(s_4 + s_5)(s_4 + s_5 - 1) - (1 - t_1)(s_1 + s_2x_i + s_3x_j), \quad (\text{A.5})$$

$$D^R(x_i, x_j) = -(1 - t_1)(s_1 + s_4 + s_5 + s_2x_i + s_3x_j), \quad (\text{A.6})$$

and the numerators  $N_n^L(x_i, x_j)$ ,  $N_n^R(x_i, x_j)$  are given by

$$\begin{aligned}
 N_1^L(x_i, x_j) &= \frac{2(1-t_1)^2(s_2-s_3)s_1}{(s_4+s_5)(s_4+s_5-1)} + \frac{2(-s_2+s_4+s_5-1)}{(s_4+s_5)(s_4+s_5-1)^2}x_i \\
 &\quad + \frac{(1-t_1)\left(3(s_4+s_5)^2 + (6s_3-7)(s_4+s_5) - 2s_2 - 6s_3 + 4\right)}{(s_4+s_5)(s_4+s_5-1)^2}x_i \\
 &\quad + \frac{(s_2-s_3)s_1(1-t_1)^2}{2(s_4+s_5)(s_4+s_5-1)^3}x_ix_j + \frac{(s_2-s_3)\left(-3(s_4+s_5)t_1 + 2t_1 + 1\right)}{2(s_4+s_5)(s_4+s_5-1)^2}x_ix_j, \quad (\text{A.7})
 \end{aligned}$$

$$\begin{aligned}
 N_2^L(x_i, x_j) &= -\frac{4(s_2-s_3)s_1t_1^2}{(s_4+s_5)(s_4+s_5-1)}x_ix_j - \frac{4(s_2-s_3)s_1}{(s_4+s_5)^2(s_4+s_5-1)^2} \\
 &\quad - \frac{(s_2-s_3)s_1}{(s_4+s_5)^2(s_4+s_5-1)^2}x_ix_j - \frac{(s_2-s_3)\left(-3(s_4+s_5)t_1 + 1\right)}{(s_4+s_5)^2(s_4+s_5-1)}x_ix_j \\
 &\quad + \frac{-2\left(-3t_1(s_4+s_5)^2 + (s_4+s_5)(-6t_1s_3 + 3t_1 + 2) - 2s_2\right)}{(s_4+s_5)^2(s_4+s_5-1)}x_i, \quad (\text{A.8})
 \end{aligned}$$

and

$$\begin{aligned}
 N_1^R(x_i, x_j) &= \frac{2s_1(s_2-s_3)t_1(1-t_1)}{(s_4+s_5)(s_4+s_5-1)^2} + \frac{2s_1(1-t_1)}{(s_4+s_5)(s_4+s_5-1)^2}x_i \\
 &\quad + \frac{6s_1(s_2-s_3)(1-t_1)\left(1-(s_4+s_5)t_1\right)}{(s_4+s_5)(s_4+s_5-1)^3} - \frac{(1+s_1t_1)(s_2-s_3)(1-t_1)}{2(s_4+s_5)(s_4+s_5-1)^2}x_ix_j, \quad (\text{A.9})
 \end{aligned}$$

$$\begin{aligned}
 N_2^R(x_i, x_j) &= \frac{(3-2s_1t_1)(s_2-s_3)(1-t_1)}{(s_4+s_5)(s_4+s_5-1)^2}x_ix_j \\
 &\quad + \frac{2s_1(1-t_1)\left(2t_1(s_4+s_5)(s_2-s_3) + s_4+s_5 - 2(s_2-s_3) - 1\right)}{(s_4+s_5)(s_4+s_5-1)^3}x_i \\
 &\quad - \frac{4s_1(s_2-s_3)\left(t_1(s_4+s_5)^2(1+4t_1) - 6t_1(s_4+s_5) + 1\right)}{(s_4+s_5)^2(s_4+s_5-1)^3} \\
 &\quad - \frac{10s_1(s_2-s_3)\left(-t_1(s_4+s_5) + 1\right)\left((s_4+s_5)(3-4t_1) + 2\right)}{(s_4+s_5)^2(s_4+s_5-1)^3} \\
 &\quad + \frac{6(s_4+s_5)(1-t_1)(4s_2+2s_3+3s_4+3s_5-3)}{(s_4+s_5)^2(s_4+s_5-1)^2}x_i \\
 &\quad + \frac{6\left(-(s_4+s_5)(3s_2+s_3+2s_4+2s_5-1) + 2s_2\right)}{(s_4+s_5)^2(s_4+s_5-1)^2}x_i \\
 &\quad + \frac{-3(s_4+s_5)(s_4+s_5+2s_2-1)(1-t_1) + (s_4+s_5-1)(4s_2-2s_3+s_4+s_5)}{(s_4+s_5)(s_4+s_5-1)^3}x_i \\
 &\quad + \frac{(s_2-s_3)}{(s_4+s_5)^2(s_4+s_5-1)}x_ix_j \\
 &\quad + \frac{-s_1(s_2-s_3)\left(t_1(s_4+s_5)^2(1+4t_1) - 6t_1(s_4+s_5) + 1\right)}{(s_4+s_5)^2(s_4+s_5-1)^4}x_ix_j. \quad (\text{A.10})
 \end{aligned}$$



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