Heavy neutrinos and the $pp \to lljj$ CMS data

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We show that the excess in the $pp \to eejj$ CMS data can be naturally interpreted within the Minimal Left Right Symmetric model (MLRSM), keeping $g_L = g_R$, if CP phases and non-degenerate masses of heavy neutrinos are taken into account. As an additional benefit, a natural interpretation of the reported ratio (14:1) of the opposite-sign (OS) $pp \to t^\pm t^\mp jj$ to the same-sign (SS) $pp \to t^\pm t^\mp jj$ lepton signals is possible. Finally, a suppression of muon pairs with respect to electron pairs in the $pp \to lljj$ data is obtained, in accordance with experimental data. If the excess in the CMS data survives in the future, it would be a first clear hint towards presence of heavy neutrinos in right-handed charged currents with specific CP phases, mixing angles and masses, which will have far reaching consequences for particle physics directions.

I. INTRODUCTION

LHC is a perfect laboratory to test Beyond Standard Model (BSM) scenarios. Recently, the CMS Collaboration announced an interesting excess in data, see point B on Fig. 1. This point is related to the process $pp \to eejj$ collected at $\sqrt{s} = 8 \text{ TeV}$ LHC corresponding to an integrated luminosity 19.7 fb$^{-1}$. Several analyses showed that this excess can be interpreted as a signal of charged gauge boson $W_2^\pm$ with mass about 2.2 TeV in the Left-Right symmetric model \cite{6,8}. It is possible when gauge couplings connected with left and right $SU(2)$ groups are not equal to each other. For a case $g_L = g_R$ see point A on Fig. 1 the measured cross section is suppressed by a factor of $\gamma_{CMS} = 0.23$ when compared with scenario in which $g_L = 0$. Moreover, the number of events with same-sign (SS) leptons to the number of events with opposite-sign (OS) leptons is

$$r = \frac{N_{SS}}{N_{OS}} = \frac{1}{14},$$

and, finally, no excess in $\mu\mu$ channel has been reported \cite{1,9}.

Theoretical analyses of the left-right symmetric models speeded up considerable in recent years \cite{10,30}, after the LHC has started its operation. It is not surprising as this collider is operating at highest available so far energies, which means that new states of matter or new interactions can be probed more effectively. For instance, the left-right symmetric models offer an elegant, dynamical explanation for suppression of right handed currents at low energies, and it might be that finally LHC can see them directly in experimental data analyses \cite{11}.

Discovery of right-handed currents and new elementary states of matter in form of a charged heavy gauge boson and heavy neutrinos would be of paramount importance for our understanding of Physics in microscale. It would also impact Physics in macroscale. For instance, details of leptogenesis depend on CP phases of decaying particles, or Big Bang Nucleosynthesis and the dark matter problem raise questions about the matter content of the Universe \cite{32,33}.

It is natural that experimental data analysis employs simplifications of theoretical models which quite often, thinking in terms of BSM, are much more complicated than worthy Standard Model theory. However, in this way conclusions can be distorted or even some interesting and natural scenarios can be overlooked. We think that our discussion here is a good example showing that including some additional theoretical issues into analysis can finally pay back in terms of better understanding of experimental results.

In this paper, we show how including details of heavy neutrinos mass spectrum, their CP phases and non-trivial

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\caption{CMS data for production of the first generation leptons with two jets in $pp$ collision with $\sqrt{s} = 8 \text{ TeV}$ \cite{1}. Blue solid line shows CMS estimation of the cross section in the MLRSM model with $g_L = g_R$, diagonal heavy neutrino couplings and $M_N = M_{W_2}/2$. Explaining excess in the data around $M_{W_2} \sim 2.2 \text{ TeV}$ (point B) requires refinement of those assumptions (point A).}
\end{figure}

\footnote{For more suitable arguments in favor of left-right symmetry, see e.g. \cite{31}.}
mixing matrix can change a picture, leading to natural interpretation of the data within MLRSM, with $g_L = g_R$ and relatively light $W^\pm_2$ charged gauge boson mass. In other words, we can get down from point A to point B on Fig. 1 while holding $g_L = g_R$. We can also accommodate value of $r$ in (1) and explain a shortage of muon pairs.

We shall consider production of $W_2$ which further decays to a charged lepton $l_i$ ($i = 1, 2, 3$) and an on-shell heavy neutrino $N_a$ ($a = 1, 2, 3$). The latter further decays mainly via 3-body process $N_a \to l_i j j$ leading to two jets and two charged leptons in the final state:

$$pp \to W^\pm_2 \to l_i^+ l_i^- N_a \to l_i^+ l_i^- j j, \quad \text{(OS)}$$

$$pp \to W^\pm_2 \to l_3^+ N_a \to l_3^+ j j. \quad \text{(SS)} \quad \text{(2)}$$

We have left considering distributions of kinematical variables of leptons and jets as well as cuts issue in the discussed processes for future studies.

II. HEAVY NEUTRINO INTERACTIONS AND THEIR CP PARITIES

We shall consider the MLRSM model in which $g_L = g_R$, $v_L = 0$ and $v_R = 0$ (what results in no $W_L - W_R$ mixing). The scale of breaking $SU(2)_R$ is set to $v_R = 4.77$ TeV, such that the mass of $W_2$ is about 2.2 TeV (see Fig. 1). Moreover, to simplify our considerations, let us assume that the scalar potential parameters are chosen such that all scalar particles beside the lightest Higgs boson have masses of order $v_R$. We leave discussion of their influence on $pp \to l l j j$ for future studies.

Neutrino mass matrix is chosen to be of the form

$$M_\nu = \begin{pmatrix} 0 & M_D \cr M_D^T & M_{R} \end{pmatrix}. \quad \text{(3)}$$

Typically, Dirac masses $M_D$ are much smaller than Majorana masses $M_R$ i.e. $M_D \ll M_R$, e.g. $M_D \sim 10^{-3}$ GeV, $M_R \sim 10^3$ GeV. Hence light neutrinos $\nu_{1,2,3}$ obtain masses $M_{\nu_{1,2,3}}$ of the order of 1 eV via type I see-saw mechanism. Unitary matrix $U$ which enters Takagi decomposition $M_\nu = U^T \text{diag}(M_{\nu_1}, M_{\nu_2}, M_{\nu_3}) U$ is of the following form:

$$U \approx \begin{pmatrix} 1 & 0 \\
0 & K_R \end{pmatrix}, \quad \text{(4)}$$

where, $K_R$ is an unitary $3 \times 3$ matrix defined by $M_R = K_R^T \text{diag}(M_{N_1}, M_{N_2}, M_{N_3}) K_R$, $M_{N_a} > 0$. For simplicity, we assume no light-heavy neutrino mixings (they are negligible or very small [23]). Such choice of $U$ means that $W^\pm_2$ does not couple to light neutrinos $\nu_a$, and heavy neutrinos $N_a$ do not couple to $W^\pm_2$. Exact neutrino mixing matrix $U$ can also be considered, which include non-zero off-diagonal light-heavy matrix elements in [4] [19]. $K_R$ matrix enters directly heavy neutrinos - $W_2$ interactions, which can be cast in the following form [25]:

$$\mathcal{L} \supset \frac{g_R}{\sqrt{2}} \bar{N}_a \gamma^\mu P_R(K_R)_{aj} l_j W^+_{2j} + \text{h.c.} \quad \text{(5)}$$

In general elements of the $K_R$ matrix can be complex. In a CP-conserving case, CP parities of heavy neutrinos are purely imaginary [36] [37] and, in fact, they can be connected with elements of the $K_R$ matrix. In practice, if CP parities of all three heavy neutrinos are the same, $\eta_{CP}(N_1) = \eta_{CP}(N_1) = \eta_{CP}(N_3) = +i$, then elements of the $K_R$ matrix can all be made real. If, for instance, $\eta_{CP}(N_1) = \eta_{CP}(N_1) = -\eta_{CP}(N_3) = +i$ then $K_{R,3}$ element is complex. Choosing different scenarios have far-reaching consequences in phenomenological studies. Let us consider processes where heavy neutrinos propagate as virtual states, then they contributions to the amplitudes must be summed over. In general, constructive or destructive interferences between heavy neutrinos can appear. For instance, in the neutrinoless double beta decay $(\beta\beta)_{0v}$ process, or its inverse collider version process $e^- e^- \to W^- W^-$, amplitudes include squared matrix elements $(K_R)_{aa}^2$. If all heavy neutrinos have the same CP parities, then elements of the $K_R$ matrix can be made real, and all heavy neutrinos contribute constructively into the amplitudes, otherwise destructive interferences can appear. Such scenarios have been considered in full details in phenomenological analyses in [58]. It has been shown there that cancellations among contributions to the amplitude from heavy neutrinos with opposite CP parities can appear. In this way, low energy $(\beta\beta)_{0v}$ constraints can be avoided and for instance the collider signal $e^- e^- \to W^- W^-$ can be substantial. We will see in the next Section that CP phases of heavy neutrinos play a crucial role also in a case of SS and OS $pp \to l l j j$ signals.

III. CROSS SECTIONS

We shall show that interference effects, CP phases of heavy neutrinos and their mass splittings are relevant for the prediction of $pp \to l l j j$ cross section. To expose interference effects in a clear way, the following three different setups will be discussed: (A) neutrinos have degenerate masses, (B) one neutrino is lighter than $W_2$, (C) two neutrinos are lighter than $W_2$, and, (D) finally, there is only small mass splitting among neutrinos. The numerical analysis has been done with the help of MADGRAPH5 (v2.2.2) [39] and with our implementation of the MLRSM in FEYNRULES (v2.0.31) [40] [41].

To simplify notation we shall denote cross-sections for the process $pp \to l_i^+ l_i^- j j$ by $\sigma_{l_i l_i j j}$ etc. For reference points it is assumed, as in CMS [41] analysis, that $M_N = M_{W_2}/2$ with diagonal and real $K_R$ mixing matrix in (5), which for $\sqrt{s} = 8$ TeV and $M_{W_2} = 2.2$ TeV, gives:

$$\sigma(pp \to W^\pm_2) = \begin{cases} 71.16 \text{ fb} \\
21.09 \text{ fb} \end{cases} \quad \text{(6)}$$

what agrees with recent estimations on $pp \to W_2 \to jj$ cross section [42]. For chosen value of $v_R$ and diagonal
matrix $K_R$ relevant branching ratios are:

$$
\text{BR}(W^+_2 \to e^+ N_a) = 0.058, \\
\text{BR}(N_a \to e^+ jj) = 0.35 
$$

when all heavy neutrinos have the same mass $M_N = M_{W_2}/2$, and

$$
\text{BR}(W^+_3 \to e^+ N_1) = 0.066 
$$

when only $M_{N_1} = M_{W_2}/2$ while $M_{N_{2,3}}$ are heavier than $W_2$.

## A. Degenerate masses of heavy neutrinos

First, let us examine the following mass pattern in which all heavy neutrinos are degenerate and lighter than $W_2$:

$$
M_N := M_{N_a} = M_{W_2}/2.
$$

In this setup, and also for small mass differences between heavy neutrinos, the narrow width approximation (NWA) will not work because of the interference effects. Let us take $K_R$ in the following form (which is in fact a product of real, orthogonal transformation and diagonal phase matrix)

$$
K_R = \left( \begin{array}{ccc} 
\cos \theta_{12} & \sin \theta_{12} & 0 \\
-e^{i\phi_2} \sin \theta_{12} & e^{i\phi_2} \cos \theta_{12} & 0 \\
0 & 0 & 1 \end{array} \right).
$$

This is a simplified version of a complete unitary rotation matrix $[13]$. In this way, we assume mixings between two lepton flavors only. Phase $\phi_2$ is connected with CP parity of heavy neutrinos $N_{1,2}$, CP-conserving case is realized when $\phi_2 = 0, \pm \pi/2, \pm \pi$. All phases which do not fulfill the above relations break CP symmetry. In general, in the MLRSM with the mass matrix of the form $[3]$ we have six CP phases; if $v_L \neq 0$ ($M_L \neq 0$) then there are 18 CP phases $[14]$. Using this simple form of the matrix $K_R$ we are already able to discuss all relevant effects connected with mixings and CP phases in the considered process.

First, in the case of degenerate neutrinos, $\sigma_{i i}^{\pm \mp}$ with $i = j$ does not depend on mixing angles at all, and is zero for $i \neq j$:

$$
\sigma_{i i}^{\pm \mp} = \frac{g^4}{4} \left| \sum_a F_{i i}^{\pm \mp} (s, M_{W_2}^2, M_{N_a}^2) (K_R^i)^a (K_R^i)_{ai} \right|^2
$$

$$
= \delta_{i j} \frac{g^4}{4} F_{i i}^{\pm \mp} (s, M_{W_2}^2, M_{N_a}^2) = \delta_{i j} \tilde{\sigma}_{SF}^{\pm \mp},
$$

where the second equality comes from the unitarity of the $K_R$, while the third defines $\tilde{\sigma}_{SF}^{\pm \mp}$. $F^{\pm \mp}$ is a function of center of mass energy $\sqrt{s}$ and masses $M_{W_2}$ and $M_N$ (leptons and constituents of jets are treated as massless).

From now on we will not write down arguments of the $F$ functions. On the other hand, for same-sign signature i.e. $l^+ l^+$ or $l^- l^-$ the mixing matrix $K_R$ does not cancel from the cross section formula:

$$
\sigma_{i i}^{\pm \mp} = \frac{g^4}{4} \left[ F_{i i}^{\pm \mp} + (-1)^{\delta_{ij}} F_{j j}^{\pm \mp} \right]^2 \sum_a (K_R^i)^a (K_R^a)_{ai}.
$$

As a consequence, cross section for $l_i^\pm l_j^\mp$ with $i = j$ is correlated with that for which $i \neq j$ i.e.

$$
\sigma_{i i}^{\pm \mp} = \begin{cases} \tilde{\sigma}_{SF}^{\pm \mp} (1 - \sin^2 2\theta_{12} \sin^2 \phi_2) & \text{for } i = j, \\
\tilde{\sigma}_{DF}^{\pm \mp} \sin^2 2\theta_{12} \sin^2 \phi_2 & \text{for } i \neq j, \end{cases}
$$

where $\tilde{\sigma}_{SF}^{\pm \mp}$ and $\tilde{\sigma}_{DF}^{\pm \mp}$ correspond to maximal values of cross sections for same-flavour (SF) and different flavour (DF) cases. For the numerical results see Fig. 2. The dif-

![FIG. 2. Cross section $pp \to lljj$ for the production of two SS light leptons with two jets $jj$ for $\phi_2 = 0$, $\pi/4$ and $\pi/2$: $e^+ e^+$ (red), $\mu^+ \mu^+$ (red, same as for $e^+ e^+$) and $e^+ e^-$ (green). Plots for $\sigma^{++}$ are of the same shape but with $\tilde{\sigma}_{SF}^{++}$ changed by $\tilde{\sigma}_{SF}^{--}$ in (14). Solid lines show formulas (14), while green and red dots are numerical results obtained in MADGRAPH5. The blue dotted line shows corresponding cross section for the OS process $pp \to W_2^+ \to e^+ e^- jj$, which is independent of $\theta_{12}$ — see (12).](image-url)
where \( c = (\hat{\sigma}_{SF}^+ + \hat{\sigma}_{SF}^-)/\hat{\sigma}_{SF}^+ \approx 1. \) As a consequence same-sign different-flavour cross section is \( \sigma_{e\mu}^\pm \approx \hat{\sigma}_{DF}^\pm (1 - r/c). \) Moreover the total cross section for \( pp \to eejj \) is then
\[
\sigma_{ee}^{(tot)} = \hat{\sigma}_{SF}^+[1 + r - r\sin^2 2\theta_{12}\sin^2 \phi_2].
\]
One can check that the total cross section \( \sigma_{ee}^{(tot)} \) is suppressed by a factor
\[
\gamma = \frac{1 + r}{1 + c} \approx 0.54
\]
with respect to \( \theta_{12} = \phi_2 = 0 \) case \( (\sigma_{ee}^{(tot,0)}). \) Our numerical calculations yield \( \sigma_{ee}^{(tot,0)} = 3.41 \text{ fb}. \) Hence when \( \theta_{12} \) and \( \phi_2 \) are chosen such that \( [15] \) is satisfied then
\[
\sigma_{ee}^{(tot)} = \gamma \sigma_{ee}^{(tot,0)} = 1.84 \text{ fb}\]
what is about 81% of the excess reported by the CMS (point B on Fig. [1]). Moreover, in consequence of [14] total cross section for production of two muons with two jets is the same as for electrons: \( \sigma_{ee}^{(tot)} = \sigma_{e\mu}^{(tot)}. \) Hence the discussed scenario would also result in excess in \( \sigma(pp \to \mu\mu jj) \). That is in contradiction with the CMS data related to \( pp \to \mu\mu jj \) [1].

### B. \( MN_1 < MW_2 < MN_{2,3} \)

In this case only \( N_1 \) can be on-shell. We choose \( MN_1 = 1.1 \text{ TeV} = MW_2/2; \) the remaining two neutrinos are much heavier, \( MN_{2,3} = 10 \text{ TeV}. \)

Here one can use narrow width approximation (NWA) to estimate cross-section for \( pp \to l_1l_2jj \) going through on-shell \( W_2 \), which decays to \( l_i \) and on-shell \( N_1 \) and the latter decays to \( l_jjj \):
\[
\sigma_{l_1l_j} = \sigma(pp \to W_2)\text{BR}(W_2 \to l_iN_1) \\
\times\text{BR}(N_1 \to l_jjj).
\]

Since quarks and leptons masses are much smaller than the \( N_1 \) mass, 3-body decay of \( N_1 \) mediated by off-shell \( W_2 \) can be treated analogously to well-known muon decay in the Fermi theory. One can check that
\[
\Gamma(N_1 \to l^-_i q_\alpha \bar{q}_\beta) = \frac{9g^2}{2048\pi^3} |(K_R)_{i\alpha}|^2 |(U_{CKM}^R)_{\alpha\beta}|^2 \\
\times MN_{a1}^2 F(x_a),
\]
where \( x_a = M_{N_a}^2/M_{W_2}^2 \) while the function
\[
F(x) = \frac{12}{x} \left[ 1 - \frac{x}{2} - \frac{x^2}{6} + \frac{1 - x}{x} \ln(1 - x) \right]
\]
encumbers full tree-level contribution from \( W_2 \) propagator [15]. The presence of such a factor makes \( N_a \) decay width really sensitive to the ratio \( x_a = M_{N_a}^2/M_{W_2}^2, \) e.g. for fixed \( M_{W_2} \) it can be enhanced by a factor of \(~27\) when \( MN_a \approx MW_2 \) with respect to the scenario \( MN_a \approx MW_2/2. \)

Summing over all possible final states and taking into account the unitarity of \( K_R \) and \( U_{CKM}^R \) one obtains the total decay width of \( N_a \)
\[
\Gamma(N_a) = \sum_{i,\alpha\beta} \left[ \Gamma(N_a \to l^-_i q_\alpha \bar{q}_\beta) + \Gamma(N_a \to l^+_i \bar{q}_\alpha q_\beta) \right]
\]
\[
= \frac{9g^2}{1024\pi^3} MN_a F(x_a)
\]
Hence the BR under consideration is
\[
\text{BR}(N_a \to l^-_i q_\alpha \bar{q}_\beta) = \frac{1}{6} |(K_R)_{i\alpha}|^2 |(U_{CKM}^R)_{\alpha\beta}|^2, \quad (23)
\]
\[
\text{BR}(N_a \to l^+_i \bar{q}_\alpha q_\beta) = \text{BR}(N_a \to l^-_i q_\alpha \bar{q}_\beta).
\]

Using assumed masses, we have scanned over \( \theta_{12} \in (0, \pi/2) \) to verify dependence of \( \sigma_{l_1l_j} \) on that angle. The CP phase \( \phi_2 \) was set to \( \pi/2 \) (CP-conserving case), i.e.
\[
K_R = \begin{pmatrix}
\cos \theta_{12} & \sin \theta_{12} & 0 \\
-\sin \theta_{12} & i\cos \theta_{12} & 0 \\
0 & 0 & 1
\end{pmatrix}. \quad (24)
\]

![FIG. 3. Cross section \( \sigma \) for the production of two opposite-sign light leptons \( l_i = e, \mu \) with two jets \( jj \) in the process \( pp \to W^\pm_2 \to l^+_i l^-_j jj \) with \( \sqrt{s} = 8 \text{ TeV} \) for \( MN_1 = MW_2/2, MN_{2,3} > MW_2 \). The dashed lines display contributions from intermediate channels \( W^\pm_2 \to e^\pm N_1 \) and \( W^\pm_2 \to \mu^\pm N_1. \) Solid lines correspond to sum over possible channels.](image-url)

The obtained dependences are shown on the Figs. 3, 4 and 5. On these plots, we present contributions to the total cross section \( \sigma_{l_1l_j} \) from subprocesses with different charges and flavors of leptons in the final state. The scale on the vertical axes is the same for all these plots to clearly show relative values of individual cross sections. The total cross section itself is shown on Fig. 4.

Let us first note that there is no interference between different contributions to \( pp \to l^+_i l^-_j jj \), see Fig. 5, because the corresponding initial states (at the parton level)
are different. Secondly, due to their large masses $N_{2,3}$ are decoupled and effectively only contributions from Feynman diagrams containing $N_1$ are relevant. In this case NWA can be used to understand qualitative dependence on the mixing angle $\theta_{12}$. Namely, using (19) one obtains $\sigma_{ee} \sim \cos^4 \theta_{12}$, $\sigma_{\mu\mu} \sim \sin^4 \theta_{12}$ and for different-flavor signature $\sigma_{e\mu} \sim \sin^2 2\theta_{12}$, cf. Figs. 4 and 5. Thirdly, one can check that, due to decoupling of $N_{2,3}$, in this setup CP phases do not influence cross sections because the interference between diagrams with different $N_a$ is suppressed by large mass of $N_{2,3}$. It is worthwhile to note that here, as in Sec. A, the difference between maximal value of $\sigma_{ee}^{\pm\pm}$ and $\sigma_{e\mu}^{\pm\pm}$ see Figs. 4 and 5 comes from the standard factor of $(-1)$ appearing in same-flavor Feynman diagrams. Finally, our numerical analysis shows that in this scenario $\sigma_{ee}^{(tot,0)} = 3.89 \text{ fb}$ hence to address CMS excess in $\sigma_{ee}^{(tot)}$ one has to adjust $\theta_{12}$ to 0.51. At the same time $\sigma_{\mu\mu}^{(tot)} = 0.21 \text{ fb}$, see Fig. 6 so there is no excess in the $\mu\mu$jj what is in accordance with CMS data [1, 9]. However as one can check, cf. Figs. 4 and 5 sum of same-sign signature cross sections i.e. $\sigma_{ee}^+ + \sigma_{ee}^-$ is nearly equal to $\sigma_{ee}^+$ for all values of mixing angle $\theta_{12}$. As a consequence, in this setup $r \approx 1$ and one cannot address $[1]$ by adjusting $\theta_{12}$.

![FIG. 4. Cross section $\sigma$ for the production of two same-sign light leptons $l_i^+ = e^+, \mu^+$ with two jets jj in the process $pp \rightarrow W_i^+ \rightarrow l_i^+ l_i^- jj$ with $\sqrt{s} = 8 \text{ TeV}$ for $M_{N_1} = M_{W_2}/2$, $M_{N_{2,3}} > M_{W_2}$.](image)

![FIG. 5. Cross section $\sigma$ for the production of two same-sign light leptons $l_i^+ = e^+, \mu^+$ with two jets jj in the process $pp \rightarrow W_i^- \rightarrow l_i^+ l_i^- jj$ with $\sqrt{s} = 8 \text{ TeV}$ for $M_{N_1} = M_{W_2}/2$, $M_{N_{2,3}} > M_{W_2}$.](image)

C. $M_{N_1,3} < M_{W_2} < M_{N_2}$

However, it turns out that one can arrange parameters of the models such that all above-mentioned experimental constraints are fulfilled. Namely, let us now consider the following mass pattern:

$$M_{N_{1,3}} = 0.925 \text{ TeV}, \quad M_{N_2} = 10 \text{ TeV} \quad (25)$$

and mixing matrix of the form:

$$K_R = \begin{pmatrix} \cos \theta_{13} & 0 & \sin \theta_{13} \\ 0 & 1 & 0 \\ -e^{i\phi_3} \sin \theta_{13} & 0 & e^{i\phi_3} \cos \theta_{13} \end{pmatrix}. \quad (26)$$

One expects that here $\mu\mujj$ signal should be suppressed due to the large mass of $N_2$. In fact, it is confirmed by numerical computations: $\sigma_{\mu\mu}^{(tot,0)} \approx 0 \text{ fb}$ while $\sigma_{ee}^{(tot,0)} = 5\text{ fb}$.
4.21 fb. Because in this scenario $N_1$ and $N_3$ are degenerate in masses, one also gets: $\sigma^{(\text{tot})}_{\tau\tau} = \sigma^{(\text{tot})}_{e\tau}$. Let us note that here $\text{BR}(W_2^+ \rightarrow e^+ N_{1,3}) = 0.071$ due to $x_{1,3} = M^2_{N_{1,3}}/M^2_{W_2} \approx 0.18$ and $x_3 = M^2_{N_2}/M^2_{W_2} > 1$, see Appendix. That enhancement of BR with respect to $\tau$ compensates deficit in $\ell$. As previously, analysis of contributions from heavy neutrinos $N_{1,3}$ gives $\sigma_{i,j}^{\pm} = \delta_{ij} \bar{\sigma}_{SF}^i$ and:

$$\sigma_{i,j}^{\pm} = \left\{ \begin{array}{ll} \bar{\sigma}_{SF}^i (1 - \sin^2 2\theta_{13} \sin^2 \phi_3) & \text{for } i = j, \\ \bar{\sigma}_{SF}^i \sin^2 2\theta_{13} \sin^2 \phi_3 & \text{for } i \neq j, \end{array} \right. \quad (27)$$

where $i, j \in \{1, 3\}$. Now the maximal values of cross sections are: $\bar{\sigma}_{SF}^1 = 2.14 \text{ fb}$, $\bar{\sigma}_{SF}^2 = 1.63 \text{ fb}$, $\bar{\sigma}_{SF}^3 = 0.48 \text{ fb}$ and $\bar{\sigma}_{DF}^{+} = 3.27 \text{ fb}$, $\bar{\sigma}_{DF}^{-} = 0.96 \text{ fb}$. Moreover,

$$\sin^2 2\theta_{13} \sin^2 \phi_3 = 1 - \frac{r}{c} \quad (28)$$

has to be satisfied in order to ensure $r = 1/14$. As previously, $c = (\bar{\sigma}_{SF}^{+} + \bar{\sigma}_{SF}^{-})/\bar{\sigma}_{SF}^0 \approx 1$ and $\gamma \approx 0.54$ what gives $\alpha_{ee}^{(\text{tot})} = \gamma \alpha_{ee}^{(\text{tot},0)} = 2.27 \text{ fb}$. It is precisely the value of $\sigma(pp \rightarrow ee\bar{\nu}j)$ reported by the CMS (point B on Fig. 1). In this way both the lepton flavor and charge independent results as well as OS (electron) dominance over SS (muon) signals can be recovered. It happens for $\theta_{13}$ and $\phi_{13}$ values which satisfies Eqs. (25).

Let us remark that naive usage of NWA would not capture dependence on CP phases $\phi_{2,3}$ at all, neither interference between diagrams with different $N_a$ correctly, what will result in wrong $\theta_{12,13}$ dependence, nor contributions from diagram with crossed lepton lines in the case of same-flavour signature. This should be kept in mind when confronting refined models with data.

D. Dependence on heavy neutrino mass splitting

$\Delta M = M_{N_2} - M_{N_1}$

Here we want to show some general dependence of the cross section on mass difference between $N_2$ and $N_1$. For simplicity it is assumed that mass of the first and third heavy neutrino are fixed to 1 TeV.

Let us note first that $\sigma$ decreases when $M_{N_2} \rightarrow M_{W_2}$, see Fig. 7. It is a consequence of decreasing branching ratio, see Eq. (31) in the Appendix. This effect is substantial; the cross section can be suppressed by a factor of 2 for considered masses.

When $M_{N_2} > M_{W_2}$ then the decay $W_2 \rightarrow lN_2$ is kinematically forbidden. It means that $N_2$ cannot be on-shell hence the contribution from such a diagram is very small because it is not enhanced by the $N_2$-resonance, cross section starts to be flat in Fig. 7.

The second effect worth mentioning, is constructive or destructive interference between diagrams with $N_1$ and $N_2$ when $M_{N_2}$ goes across $M_{N_{1,3}}$. Due to very small width of $N_a$, $\Gamma(N_a) \sim 10^{-3} \text{ GeV}$, the interference effect is visible only in the ‘very degenerate’ case i.e when mass difference $|M_{N_{1,3}} - M_{N_3}|$ between heavy neutrinos is smaller than about 0.005 GeV. Let us stress that due to these interference effects cross section $\sigma$ can be suppressed by an additional factor of 0.5 or increased by 1.5, see Fig. 7. Hence, very small mass splitting between heavy neutrinos can be a source of additional suppression/enhancement of the discussed cross section. However, as a width is very small, it might be difficult to discover such effects experimentally (energy resolution).

IV. SUMMARY AND OUTLOOK

We have revisited production of two light charged leptons and two jets in $pp$ collision in the context of the genuine MLRSM with $g_L = g_R$. Taking into account details on the neutrino mass matrix parameters, interesting conclusions can be derived. Recent CMS data showed that:

(i) there is an excess in the total $pp \rightarrow ee\bar{\nu}j$ cross section at about $M_{W_2} \approx 2.2 \text{ TeV}$;

(ii) there is a suppression of same-sign electron pairs with respect to opposite-sign pairs in $pp \rightarrow ee\bar{\nu}j$ events;

(iii) there is a suppression of muon pairs with respect to electron pair in $pp \rightarrow ll\bar{\nu}j$ events.

These facts cannot be explained within the MLRSM with $g_L = g_R$, degenerate heavy neutrino mass spectrum, and no neutrino mixings in $K_R$.

However, we have shown that all the facts (i)–(iii) listed above can be reconciled with the $g_L = g_R$ MLRSM, if non-degenerate heavy neutrino mass spectrum, neutrino mixings in $K_R$ and CP phases are taken into account.
We also conclude that it is worth to undertaken more careful analyses of the neutrino sector when exclusion plots are considered, otherwise too strong limits can be inferred from a simplified scenario (in this case assuming real neutrino mixing matrix elements with degenerate heavy neutrinos). An example is specific, but conclusions which we can derive are more general as heavy neutrinos are present within many BSM models.

In our analyses we kept $M_{W_2}$ fixed at the CMS value 2.2 TeV, however, in the light of leptogenesis [46], it would be interesting to check if it is possible to reproduce CMS data with $M_{W_2}$ shifted up to about 3 TeV by relaxing $M_{W_2} - M_N$ mass relation ($M_N = M_{W_2}/2$) and exploring wide space of heavy neutrino mixing angles, phases and masses (not necessarily of the degenerate nature), similarly as we have made in this work. It will be worthwhile to study that issue when better statistics is available.

As an outlook, we would like to check more carefully contributions from the scalar sector in MLRSM and confront our scenarios which include heavy neutrino mixing parameters and CP phases with other delicate low-energy data as neutrinoless double beta decay, just to mention [11, 30, 47–50].

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### APPENDIX

We collect here some basic formulas useful in calculations and basic estimations.

- **Quark gauge interactions** [53]:

\[
\mathcal{L} \supset \frac{g_R}{\sqrt{2}} \bar{l}_a \left( U_{CKM}^R \right)^{\alpha \beta} \gamma_5 q_\beta W^-_{2\mu} + \frac{g_R}{\sqrt{2}} \bar{q}_\beta \gamma^\mu P_+ \left( U_{CKM}^R \right)_{\beta a} q_a W^+_{2\mu}. \tag{29}
\]

- **2-body decay contributions to BR were calculated in FeynRules.** We treat leptons $l_i^\pm$ as massless:

\[
\Gamma(W_2^+ \rightarrow l_i^+ N_a) = \frac{g^2_W M_{W_2}}{96\pi} |(K_R^i)_{ia}|^2 F_W(x_a), \tag{30}
\]

where $x_a = M^2_{N_a}/M^2_{W_2}$ and $F_W(x) = (2 - 3x + x^3\theta(1-x))$. $\theta$ function in the definition of $F_W$ takes care of kinematic constraints for the decays. Because even for top quark the ratio $M^2_{W_2}/M^2_{W_2}$ is of the order of $10^{-2}$ one can treat quarks in the final states as massless. Hence, taking into account $F_W(0) = 2$ and summing over colors:

\[
\Gamma(W_2^+ \rightarrow q_\alpha \bar{q}_\beta) = \frac{g^2_W M_{W_2}}{16\pi} |(U_{CKM}^R)^{\alpha \beta}|^2, \tag{31}
\]

\[
\Gamma(W_2^- \rightarrow \bar{q}_\alpha q_\beta) = \frac{g^2_W M_{W_2}}{16\pi} |(U_{CKM}^R)^{\alpha \beta}|^2. \tag{32}
\]

That yields the total width of $W_2^\pm$

\[
\Gamma(W_2^\pm) = \frac{g^2_W M_{W_2}}{96\pi} \left[ \sum_a F_W(x_a) + 18 \right] \tag{33}
\]

and branching ratio for $W_2^\pm \rightarrow l_i^+ N_a$:

\[
\text{BR}(W_2^\pm \rightarrow l_i^+ N_a) = \frac{|(K_R^i)_{ia}|^2 F_W(x_a)}{18 + \sum_c F_W(x_c)}. \tag{34}
\]

That formula gives very good estimate of branching ratio; e.g. for $x_1 = 1/4$, $x_{2,3} > 1$ and $K_R = 1$ one gets $\text{BR}(\ldots)/\text{BR}(\ldots)_{665} \approx 0.0657/0.0659 \approx 0.997$, and similarly for $x_{1,2,3} = 1/4$: $\text{BR}(\ldots)/\text{BR}(\ldots)_{665} \approx 0.0581/0.0582 \approx 0.998$.

The branching ratio for decay to quarks is:

\[
\text{BR}(W_2^\pm \rightarrow q_\alpha \bar{q}_\beta) = \frac{|(U_{CKM}^R)_{\beta a}|^2}{18 + \sum_c F_W(x_c)} \tag{35}
\]

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