CP violation in rare lepton-number-violating $W$ decays at the LHC

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ABSTRACT: Some models of leptogenesis involve a quasi-degenerate pair of heavy neutrinos $N_{1,2}$ whose masses can be small, $O$(GeV). Such neutrinos can contribute to the rare lepton-number-violating (LNV) decay $W^\pm \rightarrow \ell_1^\pm \ell_2^\pm (q\bar{q})^\mp$. If both $N_1$ and $N_2$ contribute, there can be a CP-violating rate difference between the LNV decay of a $W^-$ and its CP-conjugate decay. In this paper, we examine the prospects for measuring such a CP asymmetry $A_{\text{CP}}$ at the LHC. We assume a value for the heavy-light neutrino mixing parameter $|B_{eN}|^2 = 10^{-5}$, which is allowed by the present experimental constraints, and consider $5 \text{ GeV} \leq M_N \leq 80 \text{ GeV}$. We consider three versions of the LHC — HL-LHC, HE-LHC, FCC-hh — and show that small values of the CP asymmetry can be measured at $3\sigma$, in the range $1\% \lesssim A_{\text{CP}} \lesssim 15\%$.

KEYWORDS: Beyond Standard Model, CP violation, Neutrino Physics

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1 Introduction

The standard model (SM) has been extremely successful in explaining most of the data taken to date. Still, there are questions that remain unanswered. For example, in the SM, neutrinos are predicted to be massless. However, we now know that neutrinos do have masses, albeit very small. What is the origin of these neutrino masses? Furthermore, are neutrinos Dirac or Majorana particles? If the latter, lepton-number-violating (LNV) processes, such as neutrinoless double-beta ($0\nu\beta\beta$) decay, may be observable.

The most common method of generating neutrino masses uses the seesaw mechanism [1–3], in which three right-handed (sterile) neutrinos $N_i$ are introduced. The diagonalization of the mass matrix leads to three ultralight neutrinos ($m_\nu \lesssim 1$ eV) and three heavy neutrinos, all of which are Majorana.

Another question is: what is the explanation for the baryon asymmetry of the universe? All we know is that out-of-equilibrium processes involving baryon-number violation and CP violation are required [4]. One idea that has been proposed to explain the baryon asymmetry is leptogenesis. Here the idea is that CP-violating LNV processes can produce an excess of leptons over antileptons. This lepton asymmetry is converted into a baryon asymmetry through sphaleron processes [5, 6].

A great deal of work has been done trying to combine these two ideas. One scenario that often arises is the appearance of a pair of heavy neutrinos, $N_1$ and $N_2$, whose masses are nearly degenerate. With this quasi-degenerate pair, leptogenesis can be produced through CP-violating decays of the heavy neutrinos [7, 8], or via neutrino oscillations [9, 10].

One particularly intriguing aspect of this scenario is that the nearly-degenerate neutrinos can have masses as small as $O$(GeV) [11]. The possibility that there can be CP-violating LNV processes involving these light sterile neutrinos has led some authors to examine ways to see such effects in the decays of mesons [12–19] and $\tau$ leptons [20, 21]. Note that these...
studies all use as motivation the neutrino minimal standard model (νMSM) [22–25], which combines the seesaw mechanism and leptogenesis, and even provides a candidate for dark matter. However, it is argued in ref. [26] (see also refs. [27, 28]) that the size of CP violation in the νMSM, while large enough to explain the baryon asymmetry of the universe, is too small to lead to a measurable effect at low energies. Still, CP-violating effects in other models may not be so small, which is the motivation for our work.

The idea of refs. [12–21] is as follows. The seesaw mechanism yields heavy-light neutrino mixing, which generates a $W$-$\ell$-$N$ coupling. This leads to decays such as $B^\pm \to D^0 \ell_1^\mp \bar{\ell}_2^\pm \pi^\pm$ via $B^\pm \to D^0 W^{\mp \pm} (\to \ell_1^\mp N)$, with $N \to \ell_2^\pm W^{\mp \pm} (\to \pi^\pm)$ [17]. CP violation occurs because there are two heavy neutrinos, $N = N_1$ or $N_2$, and these are nearly degenerate in mass. The interference of the two amplitudes leads to a difference in the rates of process and anti-process, which is a signal of CP violation.

The key point here is that the underlying LNV process is a $W$ decay. In the above meson and $\tau$ decays, the $W$ is virtual, but similar effects can be searched for in the decays of real $W$s at the LHC. To be specific, the $0\nu\beta\beta$-like process is $W^- \to \ell_1^- \bar{\ell}_2^- (q\bar{q})^+$. This decay has been studied extensively, both theoretically [29–37] and experimentally [38–41, 43, 44], as a signal of LNV. In the present paper, we push this further and study CP violation in this decay.

We consider both the decay $W^- \to \ell_1^- \bar{\ell}_2^- (q\bar{q})^+$ and its CP-conjugate. In order to generate a CP-violating rate difference between the two processes, the two interfering amplitudes mediated by the nearly-degenerate $N_1$ and $N_2$ must have different CP-odd and CP-even phases. The CP-odd phase difference is due simply to different couplings of the two heavy neutrinos. As for the CP-even phase difference, this can be generated through propagator effects or heavy neutrino oscillations. (These mirror the two different ways of producing CP-violating LNV processes for leptogenesis.) We take both into account in our study of these decays at the LHC. We will show that, if the new-physics parameters are such that $W^- \to \ell_1^- \bar{\ell}_2^- (q\bar{q})^+$ is observable, a CP-violating rate asymmetry $A_{CP}$ may be as well.

In section 2, we consider the decay $W^- \to \ell_1^- \bar{\ell}_2^- (q\bar{q})^+$. We work out the individual amplitudes $M_i^{--}$, the square of the total amplitude, $|M_{1}^{--} + M_{2}^{--}|^2$, and the CP asymmetry $A_{CP}$. The experimental prospects for measuring $A_{CP}$ are examined in section 3. We compute the expected number of events at the LHC and the corresponding minimal value of $|A_{CP}|$ measurable. We include the production of $W^\pm$ in $pp$ collisions, and take into account the lifetime of the $N_i$ and experimental efficiency. A summary & discussion are presented in section 4.

2 $W^- \to \ell_1^- \bar{\ell}_2^- (f\bar{f})^+$

As described in the Introduction, the seesaw mechanism produces three ultralight neutrinos, $\nu_j$ ($j = 1, 2, 3$), and three heavy neutrinos, $N_i$ ($i = 1, 2, 3$). The flavour eigenstates $\nu_\ell$ are expressed in terms of the mass eigenstates as follows:

$$\nu_\ell = \sum_{j=1}^{3} B_{ij} \nu_j + \sum_{i=1}^{3} B_{iN_i} N_i.$$ (2.1)
Here the parameters $B_{iN_i}$ describe the heavy-light neutrino mixing. These parameters are small, but nonzero. Because of this, there are $W$-$\ell$-$N_i$ couplings. We are particularly interested in the couplings that involve the nearly-degenerate heavy neutrinos $N_1$ and $N_2$. They are

$$\mathcal{L} \supset \frac{g}{\sqrt{2}} \bar{\ell}_i \gamma^\mu P_L (B_{iN_1} N_1 + B_{iN_2} N_2) W_\mu + h.c.$$  \hfill (2.2)

These couplings generate the $W$ decay $W^- \to \ell_1^- \bar{N}_i$. Using the fact that the $N_i$ is Majorana ($N_i = \bar{N}_i$), the $\bar{N}_i$ can subsequently decay (as an $N_i$) to $\ell_2^+ W^{*-} (\to f' \bar{f})$, where $f' \bar{f} = q' \bar{q}$ or $\ell_3^+ \nu_{\ell_3}$.

This leads to the (apparently) LNV $W$ decay $W^- \to \ell_1^- \ell_2^- (f' \bar{f})^+$. But if $f' \bar{f} = \ell_3^+ \nu_{\ell_3}$, there is a complication. The $\bar{N}_i$ can also decay as an $\ell_1^- \ell_2^- W^{*-} (\to \ell_2^- \nu_{\ell_2})$. This is a lepton-number-conserving (LNC) decay. But since neither the $\bar{\nu}_{\ell_2}$ nor the $\nu_{\ell_3}$ is detected, this final state is indistinguishable from the one above. That is, there are effectively both LNV and LNC contributions to the same decay. Since we wish to focus on pure LNV decays, hereafter we consider only $f' \bar{f} = q' \bar{q}$.

Thus, we have the rare LNV $W$ decay $W^- \to \ell_1^- \ell_2^- (q' \bar{q})^+$. This is the same decay that appears with a virtual $W$ in the decays of mesons and $\tau$ leptons, studied in refs. [12–19] and [20, 21], respectively. In those papers, it was pointed out that the interference between the $N_1$ and $N_2$ contributions can lead to a CP-violating rate difference between process and anti-process. But if this effect is present in these processes, it should also be seen in rare LNV decays of a real $W$. In the present paper we study the prospects for measuring CP violation in such decays at the LHC.

### 2.1 Preamble

It is useful to make some preliminary remarks. For the decay $W^- \to F$, where $F$ is the final state, the signal of CP violation will be a nonzero value of

$$A_{CP} = \frac{BR(W^- \to F) - BR(W^+ \to \bar{F})}{BR(W^- \to F) + BR(W^+ \to \bar{F})}.$$  \hfill (2.3)

Suppose this decay has two contributing amplitudes, $A$ and $B$. The total amplitude is then

$$A_{tot} = A + B = |A| e^{i\phi_A} e^{i\delta_A} + |B| e^{i\phi_B} e^{i\delta_B},$$  \hfill (2.4)

where $\phi_{A,B}$ and $\delta_{A,B}$ are CP-odd and CP-even phases, respectively. With this,

$$A_{CP} = \frac{2|A||B| \sin(\phi_A - \phi_B) \sin(\delta_A - \delta_B)}{|A|^2 + |B|^2 + 2|A||B| \cos(\phi_A - \phi_B) \cos(\delta_A - \delta_B)}.$$  \hfill (2.5)

The point is that, in order to produce a nonzero $A_{CP}$, the two interfering amplitudes must have different CP-odd and CP-even phases. In $W^- \to \ell_1^- \ell_2^- (q' \bar{q})^+$, the two amplitudes are $W^- \to \ell_1^- \bar{N}_{1,2}$, with $\bar{N}_{1,2}$ each subsequently decaying to $\ell_2^- (q' \bar{q})^+$. The two CP-odd phases are $\arg[B_{\ell_1 N_1} B_{\ell_2 N_1}]$ and $\arg[B_{\ell_1 N_2} B_{\ell_2 N_2}]$, which can clearly be different.
For the CP-even phases, things are a bit more complicated. The usual way such phases are generated is via gluon exchange (which is why they are often referred to as “strong phases”). However, since this decay involves the $W^\pm, \ell_1^\pm$ and $N_i$, which are all colourless, this is not possible. Instead, the CP-even phases can be generated in one of two ways. First, the propagator for the $N_i$ is proportional to
\[
\frac{1}{(p_N^2 - M_N^2) + i M_N \Gamma_N} = \frac{1}{\sqrt{(p_N^2 - M_N^2)^2 + M_N^2 \Gamma_N^2}} e^{i \eta_i},
\]
with
\[
\tan \eta_i = \frac{M_N \Gamma_N}{(p_N^2 - M_N^2)}. \tag{2.6}
\]
Thus, $\eta_i$ is the CP-even phase associated with the propagator. Since $N_1$ and $N_2$ do not have the same mass — they are nearly, but not exactly, degenerate — if one of the $N_i$ is on shell ($p_N^2 = M_N^2$), the other is not. This creates a nonzero CP-even phase difference: the on-shell $N_i$ has $\eta_i = -\pi/2$, while the CP-even phase of the other $N_i$ obeys $|\eta_i| < \pi/2$. This leads to what is known as resonant CP violation.1

A second way of generating a CP-even phase difference is through oscillations of the heavy neutrinos. As we will see below, the time evolution of a heavy $N_i$ mass eigenstate involves $e^{-i E_i t}$ (in addition to the exponential decay factor). Since $N_1$ and $N_2$ do not have the same mass, their energies are different, leading to different $e^{-i E_i t}$ factors. This is another type of CP-even phase difference, and can also lead to CP violation.

Below we derive the amplitudes for $W^- \rightarrow \ell_1^- \bar{N}_i$, with each $\bar{N}_i$ subsequently decaying to $\ell_2^- (q \bar{q})^+$, including both types of CP-even phases.

### 2.2 Decay amplitudes $\mathcal{M}_{i}^{--}$

Consider the diagram of figure 1, with $N_i = N_1$. If this were the only contribution, its amplitude could be written simply as the product of two amplitudes, one for $W^- \rightarrow \ell_1^- \bar{N}_1$, the other for $\bar{N}_1 \rightarrow \ell_2^- (q \bar{q})^+$. However, because there are contributions from $N_1$ and $N_2$, and because $N_1$ and $N_2$ cannot be on shell simultaneously, we must include the heavy neutrino propagator.

Furthermore, although the neutrino is produced as $\bar{N}_i$, it actually decays as $N_i$, leading to the fermion-number-violating and LNV process $W^- \rightarrow \ell_1^- \ell_2^- (q \bar{q})^+$. This implies that (i) conjugate fields will be involved in the amplitudes, and (ii) the amplitudes will be proportional to the neutrino mass.

We can now construct the amplitudes $\mathcal{M}_{i}^{--} \equiv \mathcal{M}(W^- \rightarrow \ell_1^- \bar{N}_1, N_1 \rightarrow \ell_2^- W^{*+} (\rightarrow (q \bar{q})^+))$. Writing $\mathcal{M}_{i}^{--} = \mathcal{M}_{i}^{\mu\nu} \epsilon_\mu \epsilon_\nu$, where $\epsilon_\mu$ is the polarization of the initial $W^-$ and

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1 Note that it is important that the $N_i$ be nearly degenerate. From eq. (2.5) we see that $A_{CP}$ is sizeable only when the two contributing amplitudes are of similar size ($|A| \sim |B|$). But if the masses of $N_1$ and $N_2$ were very different, the size of their contributions to the decay would also be very different, leading to a small $A_{CP}$. 

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Thus, we can consider only the cases. As a result, when the total amplitude is squared, \( W \) and the total amplitude is for this second diagram is adds the diagrams, while if \( p \ell \) is due to the quantum-mechanical evolution of the \( N_i \) state; its energy \( E_i \) is evaluated in the rest frame of the decaying \( W \). In the second line, we have taken the transpose of the third term, writing the current in terms of conjugate fields, \( \psi^c = C \bar{\psi}^T \). And in the third line, we have replaced \( N_i \) \( N'_i \) by the neutrino propagator.

Another contribution to this process comes from a diagram like that of figure 1, but with \( \ell_1 \leftrightarrow \ell_2 \). The amplitude for this diagram is the same as that above, but with (i) \( p_N \to p'_N \) and (ii) \( \bar{\ell}_1 \gamma^{\mu} P_R e_2 \gamma^{\nu} P_R e_1 = -\bar{\ell}_1 \gamma^{\nu} \gamma^{\mu} P_R e_2 \). Now, if \( \ell_1 \neq \ell_2 \), one simply adds the diagrams, while if \( \ell_1 = \ell_2 \), there is an additional minus sign. Thus, the amplitude for this second diagram is

\[
\mathcal{M}_i^{\mu\nu} = \pm \frac{g^2}{\sqrt{2}} B_{\ell_1 N_i} B_{\ell_2 N_i} M_i e^{-\Gamma_i t/2} e^{-iE_i t} \frac{p_N - M_i^2 + i\Gamma_i M_i}{p_N - M_i^2 + i\Gamma_i M_i} \mathcal{L}^{\mu\nu},
\]

where \( L^{\mu\nu} = \bar{\ell}_1 \gamma^{\mu} \gamma^{\nu} P_R e_2 \). In the first line, the first term is the amplitude for \( W^- \to \ell_1^- \bar{N}_i \), the second term is the time dependence of the \( N_i \) state, and the third term is the amplitude for \( N_i \to \ell_2^- W^{*+} \). The \( e^{-iE_i t} \) factor is due to the quantum-mechanical evolution of the \( N_i \) state; its energy \( E_i \) is evaluated in the rest frame of the decaying \( W \). In the second line, we have taken the transpose of the third term, writing the current in terms of conjugate fields, \( \psi^c = C \bar{\psi}^T \). And in the third line, we have replaced \( N_i \) \( N'_i \) by the neutrino propagator.

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\[
\mathcal{M}_i^{\mu\nu} = \pm \frac{g^2}{\sqrt{2}} B_{\ell_1 N_i} B_{\ell_2 N_i} M_i e^{-\Gamma_i t/2} e^{-iE_i t} \frac{p_N - M_i^2 + i\Gamma_i M_i}{p_N - M_i^2 + i\Gamma_i M_i} \mathcal{L}^{\mu\nu},
\]

and the total amplitude is \( \mathcal{M}_i^{\mu\nu} + \mathcal{M}_i^{\mu\nu} \). Now, the dominant contributions to these amplitudes come from (almost) on-shell \( N_i \)s. This means that, while both diagrams lead to \( W^- \to \ell_1^- \ell_2^- (q\bar{q})^+ \), the final-state particles do not have the same momenta in the two cases. As a result, when the total amplitude is squared, \( \mathcal{M}_i^{\mu\nu} \) and \( \mathcal{M}_i^{\mu\nu} \) will not interfere. Thus, we can consider only the \( \mathcal{M}_i^{\mu\nu} \); the contribution from the \( \mathcal{M}_i^{\mu\nu} \) will be identical.
We can now compute $\mathcal{M}^{\mu\nu} = \mathcal{M}_1^{\mu\nu} + \mathcal{M}_2^{\mu\nu}$. Writing
\begin{equation}
B_{\ell_1N_1}B_{\ell_2N_1} = B_1 e^{i\phi_1}, \quad B_{\ell_1N_2}B_{\ell_2N_2} = B_2 e^{i\phi_2}, \tag{2.9}
\end{equation}
we have
\begin{equation}
\mathcal{M}^{\mu\nu} = \frac{g^2}{2} \left( \frac{M_1 B_1 e^{i\phi_1} e^{-\Gamma_1 t/2} e^{-iE_1 t}}{p_N^2 - M_1^2 + i\Gamma_1 M_1} + \frac{M_2 B_2 e^{i\phi_2} e^{-\Gamma_2 t/2} e^{-iE_2 t}}{p_N^2 - M_2^2 + i\Gamma_2 M_2} \right) L^{\mu\nu} \tag{2.10}
\end{equation}

2.3 $|\mathcal{M}_{\text{tot}}^{-}\rangle^2$

The complete amplitude is $\mathcal{M}_{\text{tot}}^{-} = \mathcal{M}_{\mu\nu}^{\mu\nu} \epsilon^{\mu\nu} = (g^2/2) A_{-}(t) L^{\mu\nu} \epsilon_{\mu\nu}$, where $A_{-}(t)$ is the piece in parentheses in eq. (2.10). The next step is to compute $|\mathcal{M}_{\text{tot}}^{-}\rangle^2$.

From the point of view of studying CP violation in the decay $W^- \rightarrow \ell_1^- \ell_2^- (q' \bar{q})^+$, the most important term in $\mathcal{M}_{\text{tot}}^{-}$ is $A_{-}(t)$. It is instructive to compare this quantity with eq. (2.4) above. In the first term of $A_{-}(t)$, we can identify the CP-odd phase $\phi_1$ and the CP-even phase associated with neutrino oscillations ($-E_1 t$). There is also a (different) CP-even phase $\eta_1$ associated with the propagator [see eq. (2.6)]. The phases of the second term can be similarly identified.

Consider now $|A_{-}(t)|^2$. We have
\begin{equation}
|A_{-}(t)|^2 = \frac{M_1^2 B_1^2 e^{-\Gamma_1 t}}{(p_N^2 - M_1^2)^2 + \Gamma_1^2 M_1^2} + \frac{M_2^2 B_2^2 e^{-\Gamma_2 t}}{(p_N^2 - M_2^2)^2 + \Gamma_2^2 M_2^2} + 2 \text{Re} \left( \frac{M_1 M_2 B_1 B_2 e^{-i\phi_1} e^{-\Gamma_{\text{avg}} t} e^{-i\Delta E t}}{(p_N^2 - M_1^2 + i\Gamma_1 M_1)(p_N^2 - M_2^2 - i\Gamma_2 M_2)} \right), \tag{2.11}
\end{equation}
where
\begin{equation}
\Gamma_{\text{avg}} = \frac{1}{2}(\Gamma_1 + \Gamma_2), \quad \Delta E = E_1 - E_2 = \frac{M_1^2 - M_2^2}{2M_W}, \quad \delta \phi = \phi_2 - \phi_1. \tag{2.12}
\end{equation}

There are two simplifications that can be made. First, in order to compute the rate for the decay, it will be necessary to integrate over the phase space of the final-state particles. Due to energy-momentum conservation, this will involve an integral over $p_N$. Since the $N_i$ can go on shell, we can use the narrow-width approximation to replace
\begin{equation}
\frac{1}{(p_N^2 - M_i^2)^2 + \Gamma_i^2 M_i^2} \rightarrow \frac{\pi}{\Gamma_i M_i} \delta(p_N^2 - M_i^2). \tag{2.13}
\end{equation}

Second, although it is important to take neutrino oscillations into account in considerations of CP violation, we do not focus on actually measuring such oscillations. (This is examined in refs. [45–47].) That is, we can integrate over time: $\int_0^\infty dt |A_{-}(t)|^2 = |A_{-}^2|$. Note that, in integrating to $\infty$, we assume that the $N_i$ are heavy enough that their lifetimes are sufficiently small that most $N_i$’s decay in the detector. We will quantify this in the next section.
Now consider the interference term. Using the narrow-width approximation, the product of propagators can be written

\[
\frac{1}{(p_N^2 - M_1^2 + i\Gamma_1 M_1)(p_N^2 - M_2^2 - i\Gamma_2 M_2)} = \frac{\Gamma_1 M_1 \pi \delta(p_N^2 - M_1^2)}{(\Delta M^2)^2 + \Gamma_1^2 M_1^2} + \frac{\Gamma_2 M_2 \pi \delta(p_N^2 - M_2^2)}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2}
\]

\[
- \frac{i \Delta M^2 \pi \delta(p_N^2 - M_2^2)}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} + \frac{\Gamma_2 M_2 \pi \delta(p_N^2 - M_2^2)}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2},
\]

(2.14)

where \(\Delta M^2 \equiv M_1^2 - M_2^2\). Note that the imaginary part is proportional to \(\Delta M^2 = (M_1 - M_2)(M_1 + M_2) \equiv \Delta M(M_1 + M_2)\). Referring to eq. (2.6), we see that the CP-even phase difference \(\eta_1 - \eta_2\) is proportional to \(\Delta M\).

Putting all the pieces together, we obtain

\[
|A_-|^2 = \frac{\pi M_1 B_2^2}{\Gamma_1^2} \delta(p_N^2 - M_1^2) + \frac{\pi M_2 B_2^2}{\Gamma_2^2} \delta(p_N^2 - M_2^2)
\]

\[
+ \frac{2 M_1 M_2 B_1 B_2}{\Gamma_{\text{avg}}^2 + (\Delta E)^2} \left( \frac{\Gamma_1 M_1 \pi \delta(p_N^2 - M_1^2)}{(\Delta M^2)^2 + \Gamma_1^2 M_1^2} + \frac{\Gamma_2 M_2 \pi \delta(p_N^2 - M_2^2)}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} \right) \left( \cos(\delta\phi)\Gamma_{\text{avg}} - \Delta E \sin(\delta\phi) \right)
\]

\[
+ \frac{2 M_1 M_2 B_1 B_2}{\Gamma_{\text{avg}}^2 + (\Delta E)^2} \left( \frac{\Delta M^2 \pi \delta(p_N^2 - M_2^2)}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} + \frac{\Delta M^2 \pi \delta(p_N^2 - M_1^2)}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} \right) \left( \cos(\delta\phi)\Delta E + \sin(\delta\phi)\Gamma_{\text{avg}} \right).
\]

(2.15)

2.4 CP violation

The time-integrated square of the amplitude for \(W^- \rightarrow \ell_1^- \ell_2^- (q'q)^+\) is therefore \(|M_-|^2 = (g^2/2)|\tilde{A}_-|^2 |L^{\mu\nu} \epsilon_{\mu j_3} j_2^\nu|^2\). The CP asymmetry is defined as [see eq. (2.3)]

\[
A_{CP} = \frac{\int d\rho \left( |\tilde{A}_-|^2 - |\tilde{A}_+|^2 \right)}{\int d\rho \left( |\tilde{A}_-|^2 + |\tilde{A}_+|^2 \right)} = \frac{\int d\rho \left( |A_-|^2 - |A_+|^2 \right)|L^{\mu\nu} \epsilon_{\mu j_3} j_2^\nu|^2}{\int d\rho \left( |A_-|^2 + |A_+|^2 \right)|L^{\mu\nu} \epsilon_{\mu j_3} j_2^\nu|^2},
\]

(2.16)

where \(|A_+|^2\) is obtained from \(|A_-|^2\) [eq. (2.15)] by changing the sign of the CP-odd phase, and \(\int d\rho\) indicates integration over the phase space.

For the phase-space integration, the only pieces that depend on the integration variables are the delta function \(\delta(p_N^2 - M_i^2)\) in eq. (2.15) and \(|L^{\mu\nu} \epsilon_{\mu j_3} j_2^\nu|^2\). The phase-space integrals are therefore

\[
\mathcal{I}(M_i) = \int d\rho \pi \delta(p_N^2 - M_i^2) |L^{\mu\nu} \epsilon_{\mu j_3} j_2^\nu|^2.
\]

(2.17)

In ref. [14], it was shown that, since \(M_1 \approx M_2\), \(\mathcal{I}(M_1) \approx \mathcal{I}(M_2)\). Thus, to a very good approximation, these terms cancel in eq. (2.16), so that

\[
A_{CP} = \frac{|\tilde{A}_-|^2 - |\tilde{A}_+|^2}{|\tilde{A}_-|^2 + |\tilde{A}_+|^2},
\]

(2.18)

where \(\tilde{A}_- (\tilde{A}_+)^*\) is the same as \(A_- (A_+)^*\), but with the \(\pi \delta(p_N^2 - M_i^2)\) factors removed.
In the numerator we have
\begin{equation}
|\tilde{A}_-|^2 - |\tilde{A}_{++}|^2 = - \frac{2M_1M_2B_1B_2}{\Gamma_2^{\text{avg}} + (\Delta E)^2} \left( \frac{\Gamma_1M_1}{(\Delta M^2)^2 + \Gamma_1^2 M_1^2} + \frac{\Gamma_2M_2}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} \right) (2\Delta E \sin(\delta \phi)) \tag{2.19}
\end{equation}
\begin{equation}
+ \frac{2M_1M_2B_1B_2}{\Gamma_2^{\text{avg}} + (\Delta E)^2} \left( \frac{\Delta M^2}{(\Delta M^2)^2 + \Gamma_1^2 M_1^2} + \frac{\Delta M^2}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} \right) (2\sin(\delta \phi)\Gamma_{\text{avg}}). \tag{2.20}
\end{equation}

In eq. (2.5), we see that $A_{\text{CP}}$ is proportional to $\sin(\phi_A - \phi_B) \sin(\delta_A - \delta_B)$, i.e., a nonzero $A_{\text{CP}}$ requires that the two interfering amplitudes have different CP-odd and CP-even phases. This is also true in the present case. Above, both terms are proportional to $\sin(\delta \phi)$ ($\delta \phi$ is the CP-odd phase difference). In the first term, the CP-even phase arises due to neutrino oscillations: $\sin(\delta_A - \delta_B)$ is proportional to $\Delta E$. And in the second term, the CP-even phase difference comes from the propagators [see eq. (2.14)]: it is proportional to $\Delta M$. In the denominator,
\begin{equation}
|\tilde{A}_-|^2 + |\tilde{A}_{++}|^2 = \frac{2M_1B_1^2}{\Gamma_1^2} + \frac{2M_2B_2^2}{\Gamma_2^2}, \tag{2.21}
\end{equation}
\begin{equation}
+ \frac{2M_1M_2B_1B_2}{\Gamma_2^{\text{avg}} + (\Delta E)^2} \left( \frac{\Gamma_1M_1}{(\Delta M^2)^2 + \Gamma_1^2 M_1^2} + \frac{\Gamma_2M_2}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} \right) (2\cos(\delta \phi)\Gamma_{\text{avg}})
\end{equation}
\begin{equation}
- \frac{2M_1M_2B_1B_2}{\Gamma_2^{\text{avg}} + (\Delta E)^2} \left( \frac{\Delta M^2}{(\Delta M^2)^2 + \Gamma_1^2 M_1^2} + \frac{\Delta M^2}{(\Delta M^2)^2 + \Gamma_2^2 M_2^2} \right) (2\cos(\delta \phi)\Delta E). \tag{2.22}
\end{equation}

We now make the (reasonable) approximations that $\Gamma_1 \simeq \Gamma_2 \equiv \Gamma$ and $M_1 \simeq M_2 \equiv M_N$ (but $\Delta M \neq 0$ and is $\ll M_N$). With the assumption that $B_1 = B_2$, $A_{\text{CP}}$ takes a simple form:
\begin{equation}
A_{\text{CP}} = \frac{2(2y - x) \sin \delta \phi}{(1 + x^2)(1 + 4y^2) + 2(1 - 2xy) \cos \delta \phi}, \tag{2.23}
\end{equation}
where
\begin{equation}
x \equiv \frac{\Delta E}{\Gamma}, \quad y \equiv \frac{\Delta M}{\Gamma}. \tag{2.24}
\end{equation}

Once again comparing to eq. (2.5), we see that $x$ and $y$ each play the role of the CP-even phase-difference term $\sin(\delta_A - \delta_B)$. Now, $x$ and $y$ reflect CP-even phases arising from neutrino oscillations and the neutrino propagator, respectively. However, they are not, in fact, independent. From eq. (2.12), we have
\begin{equation}
\Delta E = \frac{M_N^2 - M_W^2}{2M_W} = \frac{2\Delta M M_N}{2M_W} \implies x = y \frac{M_N}{M_W}. \tag{2.25}
\end{equation}
Thus, $y$ is always present; $x$ is generally subdominant, except for large values of $M_N$.

Furthermore, we note that $x$ and $y$ have the same sign, and that $|x| < |y|$. Thus, $|2y - x| \leq |2y|$. That is, as $|x|$ increases, $A_{\text{CP}}$ decreases. We therefore expect to see smaller CP-violating effects for larger values of $M_N$. The reason this occurs is as follows. Above, we said that $x$ and $y$ each play the role of $\sin(\delta_A - \delta_B)$. However, in this system, their contributions have the opposite sign, hence the factor $2y - x$ in eq. (2.21).
In order to get an estimate of the potential size of $A_{CP}$, we set $\delta\phi = \pi/2$. In figure 2, we show $A_{CP}$ as a function of $y$, for various values of $M_N$. We see that large values ($\geq 0.9$) of $|A_{CP}|$ can be produced for light $M_N$. The maximal values of $|A_{CP}|$ are found when $y \simeq \pm \frac{1}{2}$, with $|A_{CP}|$ decreasing for larger values of $|y|$. As expected, the size of $|A_{CP}|$ decreases as $M_N$ increases, with $|A_{CP}|_{\text{max}} < 0.6$ for larger values of $M_N$.

3 Experimental analysis

In this section, we explore the prospects for measuring $A_{CP}$ at the LHC. We consider three versions of the LHC: (i) the high-luminosity LHC (HL-LHC, $\sqrt{s} = 14$ TeV, peak $L_{\text{int}} = 3$ ab$^{-1}$), (ii) the high-energy LHC (HE-LHC, $\sqrt{s} = 27$ TeV, peak $L_{\text{int}} = 15$ ab$^{-1}$) [48], (iii) the future circular collider\footnote{The Future $e^+e^-$ Circular Collider, FCC-ee (TLEP) [50], would also be a promising place to make this measurement.} (FCC-hh, $\sqrt{s} = 100$ TeV, peak $L_{\text{int}} = 30$ ab$^{-1}$) [49]. We implement the model in FeynRules [34, 51] and use MadGraph [52] to generate events.

The CP asymmetry of eq. (2.18) involves the branching ratios of the decay $W^- \rightarrow \ell^-_1 \ell^-_2 (q'\bar{q})^+$ and its CP-conjugate decay, $W^+ \rightarrow \ell^+_1 \ell^+_2 (q\bar{q}')^-$. Another way of describing $A_{CP}$ is: given an equal number of initial $W^-$ and $W^+$ bosons,

$$A_{CP} = \frac{N_{--} - N_{++}}{N_{--} + N_{++}},$$

where $N_{--}$ and $N_{++}$ are the number of observed events of $W^- \rightarrow \ell^-_1 \ell^-_2 (q'\bar{q})^+$ and $W^+ \rightarrow \ell^+_1 \ell^+_2 (q\bar{q}')^-$, respectively.

But there is a problem: these decays are not measured directly at the LHC. Instead, one has $pp$ collisions, so that the processes are $\bar{u}_id_j \rightarrow W^- \rightarrow \ell^-_1 \ell^-_2 (q'\bar{q})^+$ and $\bar{d}_ju_i \rightarrow W^+ \rightarrow \ell^+_1 \ell^+_2 (q\bar{q}')^-$, where $u_i$ and $d_j$ represent up-type and down-type quarks, respectively.
Since protons do not contain equal amounts of $\bar{u}, d_j$ and $\bar{d}_j u_i$ pairs, the number of $W^-$ and $W^+$ bosons produced will not be the same, and this must be taken into account in the definition of the CP asymmetry.

This is done by changing eq. (3.1) to

$$A_{CP} = \frac{N_{pp}^{++}/\sigma^- - N_{pp}^{--}/\sigma^+}{N_{pp}^{-+}/\sigma^- + N_{pp}^{--}/\sigma^+} = \frac{R_W N_{pp}^{pp} - N_{pp}^{pp}}{R_W N_{pp}^{pp} + N_{pp}^{pp}},$$

where $N_{pp}^{++}$ and $N_{pp}^{--}$ are the number of observed events of $pp \to XW^- (\to \ell_1^- \ell_2^- (q\bar{q})^+)$ and $pp \to XW^+ (\to \ell_1^+ \ell_2^+ (q\bar{q})^-)$, respectively, and $R_W = \sigma^+/\sigma^-$, with

$$\sigma^+ = \sigma(pp \to W^+X), \quad \sigma^- = \sigma(pp \to W^-X).$$

Experimentally, it is found that $R_W = 1.295 \pm 0.003 \ (stat) \pm 0.010 \ (syst)$ at $\sqrt{s} = 13\,\text{TeV}$ [53]. Presumably, $R_W$ can be measured with equally good precision (if not better) at higher energies, so it is clear how to obtain a CP-violating observable from the experimental measurements.\footnote{In ref. [54], it is argued that a more promising way to search for $W^- \to \ell_1^- \ell_2^- (q\bar{q})^+$ is to use $W^-$s coming from the decay of a $t$. If this is true, then if such a decay is observed, one can measure CP violation in these decays using the above formalism. And since top quarks mainly arise through $t\bar{t}$ production, there are equal numbers of $W^-$ and $W^+$ bosons, so that an adjustment using $R_W$ is not required.}

Now, given a CP asymmetry $A_{CP}$, the number of events ($N_{\text{events}} = N_{pp}^{++} + N_{pp}^{--}$) required to show that it is nonzero at $n\sigma$ is

$$N_{\text{events}} = \frac{n^2 \epsilon}{A_{CP}^2},$$

where $\epsilon$ is the experimental efficiency. This can be turned around to answer the question: given a certain total number of events $N_{\text{events}}$, what is the smallest value of $|A_{CP}|$ that can be measured at $n\sigma$?

There are two ingredients to establishing $N_{\text{events}}$. The first is the cross section for $pp \to XW^\mp$, multiplied by the branching ratio for $W^\mp \to \ell_1^\mp N_i(N_i)$, and further multiplied by the branching ratio for the decay of $N_i(N_i)$ to the final state of interest. The branching ratio for $W^\mp \to \ell_1^\mp N_i(N_i)$ depends on the value of the heavy-light mixing parameter $|B_{\ell_i N_i}|^2$. Constraints on this quantity can be obtained from experimental searches for the same $0\nu\beta\beta$-like process we consider here. A summary of these constraints can be found in ref. [55]. For 5 GeV $\lesssim M_N \lesssim 50$ GeV, $|B_{\ell N}|^2 \lesssim 10^{-5}$ ($\ell = e, \mu, \tau$), but the constraint is weaker for larger values of $M_N$. In our analysis, to be conservative, we take $|B_{\ell N}|^2 = 10^{-5}$ for all values of $M_N$.

We now use MadGraph to calculate the cross sections for $pp \to X\ell_1^\mp N$, with $N \to \ell_2^\mp (q\bar{q})^\pm$ and $pp \to X\ell_1^+ N$, with $N \to \ell_2^- (q\bar{q})^-$. The results are shown in table 1. In the table, we consider $M_N = 5$ GeV and 50 GeV. For other neutrino masses that obey $M_N \ll M_W$, such as $M_N = 1$ GeV or 10 GeV, the numbers do not differ much from those for $M_N = 5$ GeV.

We also present in table 1 the expected number of events, based on the cross section and integrated luminosity of the machine. However, that is not necessarily the final answer. The second ingredient is to look at the $N$ lifetime and see what percentage of the
Table 1. Predicted cross sections and number of events for $pp \rightarrow X \ell^-_1 \ell^-_2 \bar{q} q$ and $pp \rightarrow X \ell^+_1 \ell^+_2 \bar{q} q$. Neutrino masses $M_N = 5$ and 50 GeV are considered. Results are given for the HL-LHC ($\sqrt{s} = 14$ TeV, peak $L_{\text{int}} = 3$ ab$^{-1}$), HE-LHC ($\sqrt{s} = 27$ TeV, peak $L_{\text{int}} = 15$ ab$^{-1}$), and FCC-hh ($\sqrt{s} = 100$ TeV, peak $L_{\text{int}} = 30$ ab$^{-1}$).

| Machine | $\sigma$(fb): $\ell^-_1 \ell^-_2 jj$ | $N_{\text{events}}$ ($\times 10^{-3}$) |
|---------|-------------------------------------|----------------------------------------|
|         | $M_N = 5$ GeV | $M_N = 50$ GeV | $M_N = 5$ GeV | $M_N = 50$ GeV |
| HL-LHC  | 51.7          | 22.3          | 155.1          | 66.9          |
| HE-LHC  | 98.1          | 42.0          | 1471.5         | 630.0         |
| FCC-hh  | 323.8         | 136.7         | 9714.0         | 4101.0        |

| Machine | $\sigma$(fb): $\ell^+_1 \ell^+_2 jj$ | $N_{\text{events}}$ ($\times 10^{-3}$) |
|---------|-------------------------------------|----------------------------------------|
|         | $M_N = 5$ GeV | $M_N = 50$ GeV | $M_N = 5$ GeV | $M_N = 50$ GeV |
| HL-LHC  | 80.0          | 31.9          | 240.0          | 95.7          |
| HE-LHC  | 131.0         | 52.8          | 1965.0         | 792.0         |
| FCC-hh  | 358.2         | 147.6         | 10746.0        | 4428.0        |

For a given value of $M_N$, it is straightforward to find the neutrino lifetime, and to convert this into a distance traveled. However, the question of how many neutrinos actually decay in the detector depends on the size of the detector, and this depends on the particular experiment. As an example, we note that, in its search for $W^- \rightarrow \ell^-_1 \ell^-_2 (f' \bar{f})^+$, the CMS Collaboration considered this question [41]. They found that, for $M_N = 10$ GeV, there was essentially no reduction factor, i.e., the percentage of neutrinos decaying in the detector was close to 100%. However, for $M_N = 5$ GeV, the reduction factor was 0.1, while for $M_N = 1$ GeV, it was $10^{-3}$. Thus, the efficiency of a given experiment for observing this decay, and measuring $A_{\text{CP}}$, depends on this reduction factor.

For a given value of $M_N$, one can determine the reduction factor, and hence the total number of measurable events $N_{\text{events}}$. In order to turn this into a prediction for the smallest value of $|A_{\text{CP}}|$ that can be measured at $n\sigma$, the experimental efficiency must be included. In refs. [40, 42], the CMS Collaboration searched for heavy Majorana neutrinos at the $\sqrt{s} = 8$ TeV LHC using the final state $\ell^-_1 \ell^-_2 jj$. Including backgrounds, detector efficiency, etc., their overall efficiency was $\sim 1\%$.

Using an overall efficiency of 1%, in table 2 we present the minimum values of $A_{\text{CP}}$ measurable at $3\sigma$ at the HL-LHC, HE-LHC and FCC-hh. The results are shown for $M_N = 5$ GeV (with a reduction factor of 0.1), $M_N = 10$ GeV (with no reduction factor), and $M_N = 50$ GeV (with no reduction factor).

From this table, we see that, as the LHC increases in energy and integrated luminosity, smaller and smaller values of $A_{\text{CP}}$ are measurable. The most promising results are for $M_N = 10$ GeV, but in all cases reasonably small values of $A_{\text{CP}}$ can be probed.
Table 2. Minimum value of $A_{\text{CP}}$ measurable at $3\sigma$ at the HL-LHC ($\sqrt{s} = 14$ TeV, peak $L_{\text{int}} = 3$ ab$^{-1}$), HE-LHC ($\sqrt{s} = 27$ TeV, peak $L_{\text{int}} = 15$ ab$^{-1}$), and FCC-hh ($\sqrt{s} = 100$ TeV, peak $L_{\text{int}} = 30$ ab$^{-1}$). Results are given for $M_N = 5$ GeV (reduction factor = 0.1), $M_N = 10$ GeV (no reduction factor), and $M_N = 50$ GeV (no reduction factor).

| Machine     | $M_N = 5$ GeV | $M_N = 10$ GeV | $M_N = 50$ GeV |
|-------------|---------------|----------------|----------------|
| HL-LHC      | 15.0%         | 4.8%           | 7.4%           |
| HE-LHC      | 5.1%          | 1.6%           | 2.5%           |
| FCC-hh      | 2.1%          | 0.7%           | 1.0%           |

4 Summary & discussion

Two subjects whose explanation requires physics beyond the SM are neutrino masses and the baryon asymmetry of the universe. The standard method for generating tiny neutrino masses is the seesaw mechanism, in which one introduces three right-handed neutrinos $N_i$. As for the baryon asymmetry, leptogenesis is often used: CP-violating, lepton-number-violating processes produce a lepton asymmetry, and this is converted into a baryon asymmetry through sphaleron processes. Models that combine these two ideas often involve a quasi-degenerate pair of heavy neutrinos $N_1$ and $N_2$: leptogenesis arises through the CP-violating decays of these heavy neutrinos.

Here, an intriguing aspect is that the masses of $N_{1,2}$ can be small, $O$(GeV). This has led to suggestions to look for CP-violating LNV effects in decays of light mesons or $\tau$ leptons. These processes all involve the exchange of a virtual $W$. However, one can also consider CP-violating LNV decays of real $W$'s at the LHC. Indeed, searches for LNV at the LHC use the decay $W^\pm \rightarrow \ell_1^\pm \ell_2^\mp (q'\bar{q})^\mp$. In this paper, we have examined the prospects for measuring CP violation in such decays at the LHC.

The point is that the decay $W^\pm \rightarrow \ell_1^\pm \ell_2^\mp (q'\bar{q})^\mp$ arises via $W^\pm \rightarrow \ell_i^\pm N_i$, with $N_i \rightarrow \ell_j^\pm W^*\mp(\rightarrow q'\bar{q})^\mp$. Here, the $W-\ell-N_i$ couplings are generated due to the heavy-light neutrino mixing of the seesaw mechanism. CP violation occurs due to the interference of the $N_1$ and $N_2$ contributions.

A signal of CP violation would be the measurement of a nonzero difference in the rates of the decay $W^- \rightarrow \ell_1^- \ell_2^+ (q'\bar{q})^+$ and its CP-conjugate. This type of CP asymmetry requires that the two interfering amplitudes have both CP-odd and CP-even phase differences. The CP-odd phase difference is due to different $W-\ell-N_1$ and $W-\ell-N_2$ couplings. A CP-even phase difference can be generated in two ways, via propagator effects or oscillations of the heavy neutrino. Both are taken into account in our study.

Our analysis has two pieces, theory predictions and experimental prospects. On the theory side, we have computed the expression for the CP-violating rate asymmetry $A_{\text{CP}}$ [eqs. (2.21) and (2.22)]. We consider neutrino masses in the range $5$ GeV $\leq M_N \leq 80$ GeV. (The LHC has poor sensitivity to smaller masses.) For various values of the neutrino mass, we compute the potential size of $A_{\text{CP}}$. For low masses, e.g., $5$ GeV $\leq M_N \leq 20$ GeV, we find that (i) the contribution of neutrino oscillations to the CP-even phase is much suppressed.
compared to that from propagator effects, and (ii) $A_{\text{CP}}$ can be large, $\gtrsim 0.9$. For large masses, e.g., $M_N \geq 60\,\text{GeV}$, the contribution of neutrino oscillations to the CP-even phase becomes important, but has the effect of reducing the CP asymmetry, $A_{\text{CP}} \leq 0.6$.

On the experimental side, we want to determine the smallest value of $A_{\text{CP}}$ that can be measured at $3\sigma$ at the LHC. This depends on the number of observed events, and we use MadGraph to find this for three versions of the LHC: (i) the high-luminosity LHC (HL-LHC, $\sqrt{s} = 14\,\text{TeV}$), (ii) the high-energy LHC (HE-LHC, $\sqrt{s} = 27\,\text{TeV}$), (iii) the future circular collider (FCC-hh, $\sqrt{s} = 100\,\text{TeV}$). We assume an experimental efficiency of 1% [40, 42]. The one input required is the size of the heavy-light neutrino mixing parameter $|B_{\ell_1 N_1}|^2$. Taking into account the present experimental constraints, in our analysis we take $|B_{\ell_1 N_1}|^2 = 10^{-5}$.

We find that, while the minimum value of $A_{\text{CP}}$ measurable at the LHC depends on the neutrino mass $M_N$, smaller and smaller values of $A_{\text{CP}}$ can be measured as the LHC increases in energy and integrated luminosity. The most promising result is for the FCC-hh with $M_N = 10\,\text{GeV}$, where $A_{\text{CP}} = O(1\%)$ is measurable. But even for the worst case, the HL-LHC with $M_N = 5\,\text{GeV}$, a reasonably small value of $A_{\text{CP}} = O(10\%)$ can be measured.

The point to take away from all of this is the following. The simple observation of the LNV decay $W^\pm \rightarrow \ell_1^\mp \ell_2^\pm (q\bar{q})^\mp$ would itself be very exciting. But the next step would then be to try to understand the underlying new physics. If a CP asymmetry in this decay were measured, it would tell us that (at least) two amplitudes contribute to the decay, with different CP-odd and CP-even phases, and would hint at a possible connection with leptogenesis models.

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