Baryogenesis from mixed particle decays

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Abstract

We consider the $CP$ violating asymmetries produced in the decay of heavy particles, studying the effects of heavy particle mixing for arbitrary mass splittings. A considerable enhancement of the asymmetries is achieved when the masses of the mixed states are comparable, and the enhancement is maximal for mass splittings of the order of the widths of the decaying particles. We apply the results to the particular case of heavy scalar neutrino decays relevant for leptogenesis scenarios.
The violation of \( CP \) is one of the crucial ingredients identified by Sakharov [1] as necessary conditions for a dynamical generation of the baryon asymmetry of the Universe. Since the known \( CP \) violation in the standard model is probably too small to be helpful in generating the observed baryon asymmetry, the existence of new \( CP \) violating interactions in most of the extensions of the standard model is particularly welcome from this point of view. In the usual scenarios of baryon (or lepton) number generation by the out of equilibrium \( B \) (or \( L \)) violating decays of heavy particles, the \( CP \) violation arises in general from the interference of the tree level and one loop decay amplitudes, which allow the phases in the complex couplings involved to show up in the partial decay rate asymmetries. The one loop contribution which is usually taken into account is the so called ‘vertex’ one, in which two light particles produced in the decay of a heavy one exchange another heavy particle to produce the required final state (Fig. 1b). However, many of the scenarios studied involve (and even require) more than one ‘flavour’ of heavy particles, allowing then for a further possibility to produce \( CP \) violation.

This new possibility arises from the so called ‘wave–function’ contribution [3, 4, 5, 6, 7], in which a loop of light particles just mixes the initial state \( \Phi_a \) with another different heavy state \( \Phi_b \), which then decays to the final state as shown in Fig. 1c. This wave contribution turns out to be comparable to the vertex one when the heavy states have large mass splittings, and may be significantly enhanced for nearly degenerate states.

The asymmetry in a global quantum number \( N \) (for instance \( B \) or \( L \)) produced in the decay of a pair made of particle \( \Phi_a \) and its antiparticle \( \bar{\Phi}_a \) is given by

\[
\epsilon_a = \sum_f \epsilon_{fa},
\]

where the quantity of interest to us is the partial rate asymmetry per decay into final

\(^1\)The existence of just one heavy triplet scalar in the minimal SU(5) theory was actually considered a problem as regards the baryon number generation, since the \( CP \) violation appeared only at three loops [3].
state $f$ (of global charge $N_f$), given by

$$\epsilon_{fa} = N_f \left[ \text{BR}(\Phi_a \to f) - \text{BR}(\bar{\Phi}_a \to \bar{f}) \right],$$

(2)

where $\text{BR}(\Phi_a \to f)$ is just the branching ratio for the decay of particle $\Phi_a$ into the final state $f$. The wave function contribution to this quantity behaves, in the limit of large mass splittings, as $\epsilon_{fa}^w \propto (M_a^2 - M_b^2)^{-1}$, due to the propagator of the intermediate state $\Phi_b$, and hence is expected to be enhanced in the limit $M_b \to M_a$. A general approach to study this quantity for arbitrary mass splittings has been considered in ref. 8. It is our purpose here to extend the formalism developed in ref. 8 and apply the results to the study of leptogenesis scenarios, for which a computation of the $CP$ violation for large mass splittings was recently obtained 7.

To be specific, we will consider the case in which the heavy decaying particles are scalars, and ignore the vertex $CP$ violation effects, which can be studied separately. The wave function mixing will have the effect of inducing an absorptive part in the heavy particle propagators, which will be responsible for the generation of the asymmetry. The effect of the one–loop self–energy diagrams in the propagators will be to modify the squared scalar mass matrix as follows

$$M_{(0)2}^{(0)2} \delta_{ab} \to H_{ab}^2 = M_{ab}^2 - i\Gamma_{ab}^2,$$

(3)

where the renormalised mass matrix $M^2$ includes the dispersive part of the loops while the matrix $\Gamma^2$ arises from the absorptive part alone. The matrices $M^2$ and $\Gamma^2$ are hermitian, but $H^2$ is not. Hence, $H^2$ will be diagonalised in general by a non–unitary transformation matrix $V$ 8], i.e.

$$(V H^2 V^{-1})_{ab} = \omega_{a}^2 \delta_{ab}.$$ 

(4)

This matrix $V$ will then transform the initial ‘flavour’ states $|\Phi_a\rangle$ into the ‘propagation’
eigenstates $|\Phi'_c\rangle$, i.e.

$$|\Phi'_c\rangle = V_{ac}^{-1}|\Phi_a\rangle.$$  \hfill (5)

Similarly, for the antiparticle states $|\bar{\Phi}_a\rangle$, one will have

$$|\bar{\Phi}'_c\rangle = V_{ca}|\bar{\Phi}_a\rangle.$$ \hfill (6)

These propagation eigenstates are the ones that will evolve as

$$|\Phi'_c(t)\rangle = e^{-i\omega_c t}|\Phi'_c(0)\rangle.$$ \hfill (7)

Considering then the transition amplitude from the state $|\Phi_a\rangle$ to a final state $|f\rangle$, we have

$$T_{fa} = \langle f|H_{int}|\Phi_a\rangle,$$ \hfill (8)

where $H_{int}$ describes the interactions of $\Phi_a$ with the final state particles. From the superposition principle and using Eqs. (3,4), one has

$$T_{fa}(t) = \sum_{b,c} T_{fb} V_{bc}^{-1} V_{ca} e^{-i\omega_c t}$$

$$\bar{T}_{fa}(t) = \sum_{b,c} T_{fb}^* V_{cb} V_{ac}^{-1} e^{-i\omega_c t}.$$ \hfill (9)

The differential partial decay rate asymmetries arising from particle mixing will be proportional to the quantities

$$\Delta_{fa}(t) \equiv |T_{fa}(t)|^2 - |\bar{T}_{fa}(t)|^2.$$ \hfill (10)

It is interesting to notice that, in the limit of degenerate propagation eigenstates, i.e. $\omega_c = \omega$, these asymmetries vanish, as can be seen from eqs. (9,10).

To continue we will concentrate in the case of mixing between just two particles, for which the matrix $V$ can be parameterised in terms of two complex mixing angles, $\theta$ and $\phi$.\footnote{The appearance of $V^{-1}$ in eq. (4) ensures that the kinetic term remains canonical, but the fact that $V^{-1} \neq V^\dagger$ implies that the propagation eigenstates are not orthonormal.}
\( \phi \), as follows

\[
V = \begin{pmatrix}
\cos \theta & -\sin \theta e^{i\phi} \\
\sin \theta e^{-i\phi} & \cos \theta
\end{pmatrix}.
\]  

(11)

Replacing this in Eq. (11) it is easy to obtain

\[
e^{2i\phi} = \frac{H_{12}^2}{H_{21}^2}; \quad (\tan 2\theta)^2 = \frac{4H_{12}^2H_{21}^2}{(H_{11}^2 - H_{22}^2)^2},
\]

(12)

where we recall that \( H_{21}^2 = M_{12}^{2*} - i\Gamma_{12}^{2*} \). The eigenvalues of \( H^2 \) are then

\[
\omega_{1,2}^2 = \frac{1}{2} \left\{ H_{11}^2 + H_{22}^2 \pm \sqrt{(H_{11}^2 - H_{22}^2)^2 + 4H_{12}^2H_{21}^2} \right\}.
\]

(13)

After an explicit computation we then get

\[
\Delta_{f_a}(t) = 2\text{Re} \left\{ T_{f_1}T_{f_2}^* [U_{a_1}U_{2a}^* - U_{a_2}U_{a_1}^*] \right\} + |T_{f_b}|^2 \left\{ |U_{ba}|^2 - |U_{ab}|^2 \right\},
\]

(14)

where we have defined

\[
U_{ab} \equiv V_{ac}^{-1}W_cV_{cb},
\]

(15)

with

\[
W_c \equiv e^{-i\omega_c t}.
\]

(16)

In the case in which the initial state under consideration is an eigenstate of the (renormalised) mass matrix, i.e. for \( M_{a b}^2 = M_{a}^2 \delta_{a b} \), these expressions simplify considerably, and we have

\[
\Delta_{f_1}(t) = 4\text{Im} \left\{ T_{f_1}T_{f_2}^* \Gamma_{12}^2 \right\} \text{Re} \left\{ (\omega_1^2 - \omega_2^2)(W_2^* - W_1^*)(\cos^2 \theta W_1 + \sin^2 \theta W_2) \right\},
\]

(17)

and a similar result holds for \( \Delta_{f_2} \) with the substitution \( W_2 \leftrightarrow W_1 \).

We will then compute in detail the integrated rate asymmetry in this case. The branching ratios entering in eq. (18) are just

\[
\text{BR}(\Phi_a \rightarrow f) = \int d\Omega_a \int_0^\infty dt \ |T_{f_a}(t)|^2,
\]

(18)
with $d\Omega_a$ the phase space element of particle $\Phi_a$. We have then

$$
\epsilon_{fa}^w = N_f \Omega_a \int_0^\infty dt \Delta_{fa}(t),
$$

where $\Omega_a = M_a/16\pi$ in our case of two body scalar decay.

Integrating over time we find

$$
\epsilon_{fa}^w = 2N_f \Omega_a \text{Im} \left\{ T_{f1} T_{f2}^* \Gamma_{12}^2 \right\} F_a,
$$

with

$$
F_1 = \frac{1}{|\omega_1^2 - \omega_2^2|^2} \left\{ \text{Re} \{\omega_1^2 - \omega_2^2\} \left[ \frac{1}{\gamma_2} - \frac{1}{\gamma_1} \right] - (M_1^2 - M_2^2) \left[ \frac{1}{\gamma_1} + \frac{1}{\gamma_2} - \frac{\gamma_1 + \gamma_2}{|\omega_1^2 - \omega_2^2|^2} \right] \right.
\left. + 2 \text{Im} \{\omega_1^2 - \omega_2^2\} \left( \frac{m_1 - m_2}{|\omega_1^2 - \omega_2^2|^2} \right) \right\},
$$

where we defined $\omega_a \equiv m_a - i\gamma_a/2$. The result for $F_2$ is similar with the replacements $\gamma_1 \leftrightarrow \gamma_2$ and $m_1 \leftrightarrow m_2$ in the expression for $F_1$.

In the case $M_2^2 \gg M_1^2 \gg |\Gamma_{ab}^2|$, one has

$$
F_a \to \frac{2}{\gamma_a (M_2^2 - M_1^2)},
$$

where in this limit $\gamma_a \simeq \Gamma_{aa}^2/M_a$ becomes just the total width of particle $\Phi_a$. Hence, the results obtained in the literature in the limit of large mass splittings can be recovered.

For decreasing mass splittings, the function $F_1$ reaches a maximum, which for $|\Gamma_{12}^2| \ll |\Gamma_{22}^2 - \Gamma_{11}^2|$ takes place for

$$
M_2^2 - M_1^2 \simeq \Gamma_{11}^2 + \Gamma_{22}^2.
$$

The value of $F_1$ at the maximum is

$$
F_1 \simeq M_1/|\Gamma_{11}^2 (\Gamma_{11}^2 + \Gamma_{22}^2)|.
$$

On the other hand, for $|\Gamma_{11}^2 - \Gamma_{22}^2| \gg |M_1^2 - M_2^2|, |\Gamma_{12}^2|$, i.e. in the limit of small mass splittings and small mixing, we get

$$
F_1 \simeq \frac{(M_2^2 - M_1^2) 2M_1}{|\Gamma_{11}^2 - \Gamma_{22}^2|^2 \Gamma_{11}^2} \left. \left\{ 1 - \frac{4\Gamma_{11}^2 \Gamma_{22}^2}{(\Gamma_{11}^2 + \Gamma_{22}^2)^2} \right\} \right. - \frac{1}{M_1 (\Gamma_{11}^2 + \Gamma_{22}^2)}. \tag{24}
$$
This result coincides with the one in ref. [8] only in the limit \( \Gamma_{11}^2 \ll \Gamma_{22}^2 \) (or \( \Gamma_{11}^2 \gg \Gamma_{22}^2 \)) and if we neglect the last term, which although small is non-vanishing and survives in the limit \( M_2^2 \rightarrow M_1^2 \) (in which the propagation eigenstates are not degenerate due to \( \Gamma_{11}^2 \neq \Gamma_{22}^2 \)).

Another simple example is for the case \( \Gamma_{11}^2 = \Gamma_{22}^2 \), for which we have

\[
\omega_2^2 - \omega_1^2 = \sqrt{(M_2^2 - M_1^2)^2 - 4|\Gamma_{12}^2|^2} \tag{25}
\]

We see that here for \( |M_2^2 - M_1^2| = 2|\Gamma_{12}^2| \) the two propagation eigenstates become degenerate, and in fact for this mass splitting \( e_{fa}^w = 0 \), but it will be non-vanishing otherwise. In particular, in the limit \( M_2^2 \rightarrow M_1^2 \) one has \( F_1 \simeq -1/(2M_1\Gamma_{11}^2) \).

Let us also emphasise that a crucial ingredient in all this computation is the proper specification of the initial state. The asymmetry of course depends on the starting basis for \( \Phi_a \) considered, and hence on the process which produces the initial state, so that ignoring the mixing at production would lead to incorrect results. For instance, in the case in which \( M^2 \propto 1 \) (or more generally whenever the matrices \( M^2 \) and \( \Gamma^2 \) commute), it is possible to change basis with a unitary transformation to make \( M^2 \) and \( \Gamma^2 \) both diagonal. Hence, the asymmetry computed in this new basis will vanish, since \( \Gamma_{12}^2 = 0 \) now. However, these new states may not be the quantum states generated in the production process, and therefore the new basis may not be the appropriate one to compute the resulting asymmetry.

Let us now consider the particular example of lepton number generation in the out of equilibrium decay of heavy scalar neutrinos, i.e. the supersymmetric version of the Fukugita and Yanagida scenario [10]. In these type of models [11, 12, 13, 14, 15], the ‘degenerate’ situation \( \Gamma_{11}^2 = \Gamma_{22}^2 \) and \( M_2^2 = M_1^2 \) actually is present in well known cases such as \( K^0\bar{K}^0 \) or \( B^0\bar{B}^0 \) mixing, where those constraints are imposed by CPT relations, and for which the integrated CP violating asymmetries are non-vanishing [1].
decay of the electroweak singlet (s)neutrinos, with masses $M \gg \text{TeV}$, produces a lepton asymmetry. This is then partially converted into a baryon asymmetry \cite{16} by the effects of the anomalous $B + L$ violation in the standard model \cite{17}, which is in equilibrium at temperatures larger than the electroweak phase transition one ($\simeq 10^2 \ \text{GeV}$). The study of the $CP$ violation in these models, considering both the ‘vertex’ part as well as the ‘wave’ contribution in the limit of large mass splittings, was carried out recently \cite{7}. We now consider the effects of mixing for arbitrary masses using the formalism introduced above\cite{4}.

The Lagrangian for the scalar neutrinos is, in a basis in which the mass matrix is diagonal,

$$\mathcal{L} = -\lambda_{\alpha} \epsilon_{\alpha\beta} \left\{ M_{\alpha} \bar{N}_{a}^{\dagger} L_{i}^{\alpha} H^{\beta} + (\bar{h}^{\beta})^{c} P_{L} \ell_{i}^{\alpha} \bar{N}_{a} \right\} + h.c. \tag{26}$$

where $\ell_{i}^T = (\nu_i, l_i^-)$ and $H^T = (H^+, H^0)$ are the lepton and Higgs doublets ($i = e, \mu, \tau$, and $\epsilon_{\alpha\beta} = -\epsilon_{\beta\alpha}$, with $\epsilon_{12} = +1$).

Since we are interested in the possible implications of small mass splittings, we will assume that the right handed neutrino masses consist of two almost degenerate states, with the third one being much heavier and hence effectively decoupled from the mixing mechanism. In this case, the effects of the third scalar neutrino can be included independently, and the mixing effects can be studied with the two flavour formalism discussed before. It is particularly interesting that scenarios with this type of spectrum have been widely considered in the literature \cite{19}, and can naturally arise in $SO(10)$ models.

We will assume that sneutrinos are produced out of equilibrium by a certain unspecified mechanism (e.g. if sneutrinos are inflaton decay products \cite{13} or the inflaton itself

\footnote{There has been a recent attempt to study the asymmetry in heavy Majorana neutrino decays in the limit of small mass splittings \cite{18}, but those results are however at variance with ours. For instance, we find a dependence on $H_{12}$ through $H_{12} H_{21}$, and not through $(H_{12} + H_{21})^2$ as in ref. \cite{18}, although we expect similar results for neutrino and sneutrino decays in the supersymmetric model considered here.}
(13)], and for simplicity consider that the states produced initially correspond to one of the eigenstates \( 5 \), say \( \tilde{N}_1 \), of the mass matrix (so that \( M^2_{12} = 0 \)). Hence, the asymmetry will be given by eqs. (20, 21), where in this model a direct computation of the absorptive part of the sneutrino propagator leads to [7]

\[
\Gamma^2_{ab} = \frac{1}{8\pi} \left[ (\lambda^\dagger \lambda)_{ba} M_a M_b + (\lambda^\dagger \lambda)_{ab} s \right], \tag{27}
\]

where the square of the four momentum will just be \( s \equiv p^2 = M^2_1 \) by the on–shell condition. The first contribution to the r.h.s. of eq. (27) is given by the slepton and Higgs loop, while the second by the lepton and Higgsino one.

As final states, we need to consider two possibilities, i.e. \( f = \tilde{L}_i^\alpha \tilde{H}^\beta \) as well as \( f = \tilde{\ell}_i^\alpha \tilde{h}^c \). For the final state with sleptons, we have \( T_{fa} = -i\epsilon_{a \beta}^\alpha M_a / \sqrt{s} \), so that only the second term in the r.h.s. of eq. (27) contributes to the total asymmetry, and we have

\[
\sum_{i,\alpha,\beta} L_f \text{Im} \left\{ T_{f1} T_{f2}^* \Gamma^2_{12} \right\} = \frac{M_1 M_2}{4\pi} \text{Im} \left\{ (\lambda^\dagger \lambda)^2 \right\} = -\frac{M_1 M_2}{4\pi} \text{Im} \left\{ (\lambda^\dagger \lambda)^2 \right\}, \tag{28}
\]

where \( L_f = +1 \) is the lepton number of the final state. On the other hand, for the decay \( \tilde{N}_a \to \tilde{\ell}_i^\alpha \tilde{h}^c \), one has \( T_{fa} = -i\epsilon_{a \beta}^\alpha \lambda_{ai} \), and only the first term in the r.h.s. of eq. (27) contributes to the total asymmetry. Since now \( L_f = -1 \), we end up with the same contribution as the one coming from the slepton channel. (Notice that the asymmetry in a given final state channel results from the mixing generated by a loop involving the particles of the other final state channel.)

So, summing the contributions from both final states we get

\[
\varepsilon^w_1 = -\frac{M_1^2 M_2}{16\pi^2} \text{Im} \left\{ (\lambda^\dagger \lambda)^2 \right\} F_1, \tag{29}
\]

where \( F_1 \) is given in eq. (21). If we use that \( \Gamma^2_{11} = (\lambda^\dagger \lambda)_{11} M_1^2 / 4\pi \), and the asymptotic

\[5\]If both \( \tilde{N}_1 \) and \( \tilde{N}_2 \) are simultaneously, but incoherently, produced, one needs just to add the asymmetries from both decays.
expressions discussed previously for $M_2^2 \gg M_1^2 \gg |\Gamma_{ab}|$, one can see that in this limit

$$\epsilon_1^w = -\frac{1}{2\pi} \frac{M_1 M_2}{M_2^2 - M_1^2} \frac{\text{Im} \left\{ \left( \lambda^\dagger \lambda \right)_{21}^2 \right\}}{\left( \lambda^\dagger \lambda \right)_{11}}$$

(30)

which coincides with the expression obtained in ref. [7].

In figure 2 we plot the total asymmetry $\epsilon_1^w$ for arbitrary mass splittings, normalised to the vertex contribution $\epsilon_1^v$ arising from the exchange of the second state $\tilde{N}_2$ [7]

$$\epsilon_1^v = -\frac{1}{4\pi} g \left( \frac{M_2^2}{M_1^2} \right) \frac{\text{Im} \left\{ \left( \lambda^\dagger \lambda \right)_{21}^2 \right\}}{\left( \lambda^\dagger \lambda \right)_{11}}$$

(31)

where $g(x) \equiv \sqrt{x} \ln[(1 + x)/x]$.

Notice that $\epsilon_1^v$ contains actually the same combination of Yukawas appearing in the wave function contribution $\epsilon_1^w$ (i.e. the factor $\text{Im}\{(\lambda^\dagger \lambda)_{21}^2\}$), so that the $CP$ violating phase cancels in the ratio.

In fig. 2 we adopted for definiteness $\Gamma_{22}^2 = M_1^2/10$, $\Gamma_{11}^2 = |\Gamma_{12}^2| = \Gamma_{22}^2/10$, plotting the result as a function of $x \equiv M_2/M_1$. In the limit of large mass splittings, the wave contribution approaches twice the value of the vertex one, as expected [7]. For decreasing mass splittings ($x \to 1$), the enhancement in the wave contribution due to the mixing of the states is apparent, and reaches a maximum value $\simeq M_2^2/(\Gamma_{22}^2 \ln 2)$ for $M_2^2 - M_1^2 \simeq \Gamma_{22}^2$ (corresponding to $x \simeq 1.05$ in this case) as discussed in eq. (23).

The dotted line corresponds to the asymptotic expression for the wave contribution in eq. (30), and gives a reasonable approximation to the result for $M_2^2 - M_1^2 > 4\Gamma_{22}^2$. For smaller values of $\Gamma_{22}^2/M_1^2$, the enhancement in the wave contribution is larger (and can be in principle of many orders of magnitude), and the asymptotic expression is valid down to smaller values of $x$. On the other hand, for $M_2^2 - M_1^2 < \Gamma_{22}^2$, $\epsilon_1^w$ decreases significantly,
and for $M_2 \to M_1$ one has

$$\frac{\epsilon_{1w}^w}{\epsilon_{1}^1} \to -\frac{\Gamma_{11}^2}{(\Gamma_{11}^2 + \Gamma_{22}^2) \ln 2},$$  \hspace{1cm} (32)

which is tiny. The results are quite insensitive to the actual values of $\Gamma_{11}^2$ and $|\Gamma_{12}^2|$, as long as this last stays much smaller than $\Gamma_{22}^2$, so that the mixing angle $\theta$ is small. If $|\Gamma_{12}^2| \sim \Gamma_{22}^2$, the maximum enhancement is somewhat smaller but the general behaviour remains similar.

It is important to keep in mind that the vertex and wave contribution arising from the exchange of the heavier third scalar neutrino $\bar{N}_3$ may however be larger than the one coming from the exchange of the second one $\bar{N}_2$, even taking into account the possible enhancements for small mass splittings, due to the probably larger Yukawa couplings involved and the unknown size of the $CP$ violating phases appearing in both channels (for three families, there are actually three independent $CP$ violating phases entering in the asymmetry [20]).

In conclusion, we have considered in detail the integrated $CP$ violating asymmetries arising from heavy particle mixing, and studied the effects that appear when the mass splittings are of the order of the particle widths. The large enhancements which can be achieved can be helpful to explain the observed baryon asymmetry of the Universe, as we have exemplified with the study of a scenario for leptogenesis.

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Figure captions:

Figure 1: Diagrams which interfere to produce the $CP$ violation in the heavy particle decay. Fig. 1b gives the so called vertex contribution while Fig. 1c gives the wave function one.

Figure 2: Wave function contribution to the $CP$ asymmetry, normalised to the vertex one, as a function of $M_2/M_1$, assuming $M_{12}^2 = 0$ and taking $\Gamma_{22}^2 = M_1^2/10$, $\Gamma_{11}^2 = |\Gamma_{12}^2| = \Gamma_{22}^2/10$. 


