Axial Currents from CKM Matrix CP Violation and Electroweak Baryogenesis

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The first principle derivation of kinetic transport equations suggests that a CP-violating mass term during the electroweak phase transition can induce axial vector currents. Since the important terms are of first order in gradients there is a possibility to construct new rephasing invariants that are proportional to the CP phase in the Cabibbo-Kobayashi-Maskawa matrix and to circumvent the upper bound of CP-violating contributions in the Standard Model, the Jarlskog invariant. Qualitative arguments are given that these new contributions still fail to explain electroweak baryogenesis in extensions of the Standard Model with a strong first order phase transition.

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I. INTRODUCTION

Following the seminal work [1] about electroweak baryogenesis (EWB) many models have been proposed in the last years, that intend to explain the baryon asymmetry of the universe (BAU) by sphaleron processes that couple to an axial quark current during a first order electroweak phase transition (EWPT). The main reason that this topic attracted such an attention is, that the related elementary particle physics is accessible to experiments these days.

However all models depend on extensions of the Standard Model (SM) since the SM fails on the following grounds:

• Lack of CP violation

Since the only source of CP violation in the Standard Model is the Cabibbo-Kobayashi-Maskawa (CKM) matrix (apart from the neutrino mass matrix, which provides an even tinier source of CP violation) one has to face that it is too weak to account for the observed magnitude of BAU.

• First order phase transition

Sakharov [2] pointed out that baryogenesis necessarily requires non-equilibrium physics. The expansion of the universe is too slow at the electroweak scale and one needs bubble nucleation during a first order EWPT. The phase diagram of the Standard Model is studied in detail [3, 4], and it is well known that there is no first order phase transition in the Standard Model for the experimentally allowed Higgs mass.

• Sphaleron bound

To avoid washout after the phase transition, the vev of the broken Higgs field has to meet the criterion $|\langle \Phi \rangle| \gtrsim T_c$, i.e. a strong first order phase transition. This results in the Shaposhnikov bound on the Higgs mass [5, 6].

In the following we will address the first point - the lack of sufficient CP violation. The strong first order phase transition is assumed to occur at about $T_c \simeq 100$ GeV and is parametrised by the velocity of the phase boundary (wall velocity) $v_w$ and its thickness $l_w$. It may be induced by adding massive scalars and gauge fields to the SM.

A first attempt to account for the BAU within the SM was given by Farrar and Shaposhnikov [7]. Their method was based on reflection coefficients in the thin wall regime and the need of different diagonalization matrices in the broken and the unbroken phases. However, it has been argued, that the fermion damping by gluons annihilates the coherent modes too fast to account for an axial quark current in the wall [8, 9].

On the other hand one important effect has been neglected by the assumption that the wall is infinitely thin. The use of a continuous wall profile in the WKB approach [10] leads to a dependence of the dispersion relation on the CP-violating phase of complex mass terms. This effect is in strong contrast to the coherent generation of axial fermion currents as discussed in [7], since in the former case the wall produces a CP sensitive mass of the fermions and the damping is required to convert this slight mass change into a displacement in the fermionic distribution functions.

A first principle derivation of this dispersion relation and the associated transport equation in the Schwinger-Keldysh formalism was given in [11]. This formalism is shortly reviewed in the next two sections. Content of the present publication is to generalise this method to several flavors and to the Standard Model type case, in which the
As additional simplification, one can treat all appearing self-energies as being in equilibrium – which is more crude than a strict linear response approximation, valid close to thermal equilibrium and whose implementation would imply additional integral terms – such that by using the Kubo-Martin-Schwinger condition and the thermal fermionic distribution function \( f_F(k \cdot u) = [\exp (k \cdot u/T) + 1]^{-1} \) \( (w^i = \gamma_w(1,0,0,v_w)) \) denotes the plasma vector in the wall frame) the collision terms in the first equation can be transformed into

\[
\frac{1}{2} \Sigma^S S^> - \frac{1}{2} \Sigma^S S^< = \frac{1}{2} \Sigma_A (S^S - f_F(k \cdot u) A).
\]

We will not solve the full transport equations, but only look for the appearing CP-violating source terms. The explicit form of the collision term will not be discussed and is generally denoted by \( \text{Coll.} \), even though it can contain CP-violating contributions \([13]\). We expect to obtain the usual classical transport equation to order \( h^0 \) and CP-violating effects to order \( h^1 \), which appear at the first order in gradients.

### III. MODEL WITH CP-VIOLATING COMPLEX MASS

To start with \([11]\) we add a pseudoscalar imaginary mass term to the normal Dirac operator. The inverse propagator in a convenient coordinate system, in the wall frame in which a particle moves perpendicular to the wall \( \vec{k}_w = 0 \) and in the case of a stationary wall, reads

\[
(S_0^{-1} - \Sigma R) \rightarrow \vec{k}_w^0 \gamma_0 + k^3 \gamma_3 + m_R(X_3) + i m_I(X_3) \gamma_5,
\]

where \( \vec{k}_w^0 := \text{sign}[k^0] (k^2 - \vec{k}_w^0)^{1/2} \). The wall velocity will enter in the boundary conditions of \( S^< \).

Although \( m_R \) and \( m_I \) are Lorentz scalars, they can appear in our solutions only in certain combinations. Since a chiral transformation will change the complex mass by a constant phase, but should not have any physical relevance, only the following terms are possible up to second order in gradients (prime means differentiation with respect to \( X_3 \), \( m_R + i m_I = m e^{i \theta}, m^2 = m_R^2 + m_I^2 \))

\[
(m^2, \theta m^2, (m^2)^{1/2}, \theta m^2, (m^2)^{1/2}, \theta' m^2).
\]

The first CP-violating effect will therefore be at least of first order in gradients.

Since the inverse propagator commutes with the spin projector \( P_s = \frac{1}{2} (1 + s^7 \gamma_0 \gamma_5), \) spin is conserved and the spin diagonal entries can be written in the block-diagonal form in spin,

\[
S^{<s,s} = P_s S^< P_s = P_s (s_0 \gamma_0 + s_1^2 \gamma_2 + s_2 \gamma_5 + s_3^3 \gamma_3),
\]

\[
A^{<s,s} = P_s AP_s = P_s (a_0 \gamma_7 + a_1^2 \gamma_1 + a_2 \gamma_5 + a_3^3 \gamma_3).
\]

A consistent iterative solution of equations \([3]\) and \([4]\) yields the following equations for \( s_0 \) and \( a_0 \) (for details
\[ \hat{C}s_0^* = 0, \quad \hat{C}a_0^* = 0, \quad \hat{K}s_0^* = \text{Coll.}, \quad \hat{K}a_0^* = 0, \]

with the constraint and kinetic operators
\[ \hat{C} = k_0^2 - \tilde{k}^2 - m^2 - s \frac{m^2 \theta'}{k_0}, \]
\[ \hat{K} = \hat{C} \partial \omega - (m^2) \frac{\theta}{k_0}, \]

where
\[ s \frac{m^2 \theta'}{k_0} \text{ and } \hat{K} \partial \omega \]

The correctly normalized solutions for the spectral functions \( a_0^* \) are
\[ a_0^* = \pi i \text{ sign}(k_0)(\partial \omega \hat{C}) \delta(\hat{C}). \]

In this form we can immediately see, that the axial current will not contain any CP violation if the wall velocity vanishes, since this would permit the solution
\[ \frac{\delta \omega}{\omega} = \frac{s \frac{m^2 \theta'}{k_0}}{2k_0k^*_0}. \]

**IV. WHY THE STANDARD MODEL (NAÍVELY) FAILS**

In the Standard Model the Jarlskog determinant \[ 16 \] is believed to be an upper bound on CP violating effects. The basis for this relation is the following reasoning: Suppose the SM Lagrangian contains two non-hermitian mass matrices for the quarks (due to the coupling to the Higgs field, denoted by \( \tilde{m}_u \) and \( \tilde{m}_d \)) while the coupling of the left-handed quarks to the W bosons is still proportional to unity in flavour space.

Using four unitary flavour matrices for the left/right handed up/down quarks (\( U^L_u, U^R_u, U^L_d, U^R_d \)) the Lagrangian can be written in terms of the mass eigenstates (\( m_u = U^L_u \tilde{m}_u U^R_u, m_d = U^L_d \tilde{m}_d U^R_d \)). The unitary matrices for the right handed quarks have no physical significance, while the product of the left handed up/down matrices lead to the CKM matrix in the coupling term between left handed quarks and W bosons (\( C = U^L_d U^L_u \)).

These Lagrangians are not in one-to-one correspondence. If we started with mass matrices, that needed the same left handed but different right handed transformation matrices, we would end up with the same CKM matrix, and the same diagonal mass matrices. If we express now our measurable quantities by the primary nondiagonal mass matrices, only combinations are allowed that do not include the right handed transformation matrices after diagonalization.

In the SM the combinations of lowest dimension, that fulfill these requirements are the matrices \( \tilde{m}_u \tilde{m}_u^\dagger \) and \( \tilde{m}_d \tilde{m}_d^\dagger \), and it turns out that the first CP sensitive contribution is the Jarlskog determinant
\[ 3 \left( \text{det}[\tilde{m}_u \tilde{m}_d^\dagger, \tilde{m}_u \tilde{m}_u^\dagger] \right) \]
\[ = \text{Tr}(Cm_u^L C^\dagger m_d^L C^\dagger m_d^L) \]
\[ \approx -2J m_u^L m_d^L m_s^2 m_s^2 \]
(8)

has dimension 12, and is suppressed by the 12th power of the W boson mass, or in a thermal system at least by the 12th power of the temperature. On these grounds the first physical effect would be of order \[ 17 \]
\[ \left( \frac{\alpha_{\text{EM}}}{2M_W^2} \right)^7 J m_u^L m_d^L m_s^2 m_s^2 \approx 10^{-22}, \]
(9)

where \( J \) denotes a specific combination of the angles of the CKM matrix \[ 16, 17 \]. For example, in the Kobayashi-Maskawa parametrization \[ 18 \]
\[ V^\text{CKM} = \left( \begin{array}{ccc} c_1 & -s_1 & -s_1s_3 \\ s_1c_2 & c_1c_2c_3 - s_2s_3c_\delta & c_1c_2c_3 + s_2s_3c_\delta \\ s_1s_2 & c_1s_2c_3 - c_2s_3c_\delta & c_1s_2c_3 - c_2s_3c_\delta \end{array} \right), \]

it is given by
\[ J = s_1^2 s_2 s_3 c_1 c_2 \sin(\delta) = (3.0 \pm 0.3) \times 10^{-5}. \]
(10)

where \( s_i \equiv \sin(\theta_i) \) and \( c_i \equiv \cos(\theta_i) \) \( (i = 1, 2, 3) \).

The calculation of the last section can as well be performed with several flavours. A numerical solution of the system shows, that the constraint equation \[ 15 \] contains a term
\[ \Im \left( \text{Tr}[m_d^L m_u^\dagger] \right) \]
(11)
as a generalization of \[ 3 \]. This term is only of dimension 3 and provides the possibility to circumvent the upper bound \[ 9 \] if one includes contributions that can produce these terms, e.g. corrections due to the thermal self-energies. This inclusion is needed since the derivatives of
the mass matrices are proportional to the mass matrices themselves. Therefore the generalization (11) can have no contribution on tree level. Even more stringent is the prejudice (17) that, in an expansion of the self-energy in masses, the most important contribution will be of the form (8) such that the bound (13) still seems to hold. However, this argument is based on the assumption, that the convergence of an expansion in the mass parameters is fast, and this turns out not to be the case.

V. SELF-ENERGIES IN THE STANDARD MODEL

The hermitian part of the thermal self-energy of the quarks in the Standard Model reads (19)

\[ \Sigma_R = \frac{1}{\kappa} (K_L P_L + K_R P_R) + \frac{1}{\kappa} (U_L P_L + U_R P_R) + M P_L + M^\dagger P_R, \]

with \( K_L, K_R, U_L, U_R \) hermitian 3 \( \times \) 3 matrices, \( M \) an arbitrary 3 \( \times \) 3 matrix, all depending on \( X_3 \), the external energy \( \omega = u \cdot k \), and the external momentum \( \kappa = \sqrt{\omega^2 - k^2} \) in the restframe of the plasma, \( P_L, P_R \) the left/right-handedness projection operators and \( u^\mu \) again the plasma vector.

In general all these coefficients can contain CP-violating contributions, but we will focus on the mass term, since it leads to the generalisation of terms of the form (19). The mass part of the thermal self-energy of the down quarks in the mass eigenbasis has the form

\[ M_d = h_1 m_d + \alpha_w C \frac{m_u^2}{m_W^2} h_2 C^\dagger m_d + \alpha_w \int \frac{m_u^2}{m_W^2} h_3 C^\dagger m_d + \alpha_w \int \frac{m_u^2}{m_W^2} h_3 C^\dagger m_d + \alpha_w \int \frac{m_u^2}{m_W^2} h_3 C^\dagger m_d + O(\alpha_w^3), \]

where \( h_1 \) and \( h_4 \) depend only on \( m_u^2 \), while \( h_2 \) and \( h_3 \) depend on \( m_c^2 \). The integral is performed over the energies and momenta of the particles in the loop. The terms including the CKM matrices result only from the loops of the charged Higgs bosons and are displayed in fig. (1). Since the derivatives of the mass matrices are proportional to the mass matrices themselves, in the combination 3 (\( \text{Tr}[M_d \partial M_d^\dagger] \)) only the derivatives of the \( h \) functions will contribute. Furthermore the first CP sensitive term has to include at least four CKM matrices, and using the relation (17)

\[ \text{Tr}(C X_1 C^\dagger X_2 C X_3 C^\dagger X_4) = -2 J \sum_{ij} \varepsilon_{ijkl} X_1^i X_2^j \varepsilon_{jmn} X_2^m X_4^n \]

for diagonal matrices \( X \) with the entries \( X^i \) and \( J \) as in (9) we find the following contributions (in the following prime denotes differentiation with respect to the Higgs vev)

\[ \text{Tr}(M_d M_d^\dagger) = \frac{\alpha^3_w}{m_W^4} \int \text{Tr}(C^\dagger m_d^2 h_4 C m_u^2 h_3 C^\dagger m_d^2 h_3 h_2) + \frac{\alpha^3_w}{m_W^4} \int \text{Tr}(C^\dagger m_d^2 C m_u^2 h_3 C^\dagger m_d^2 h_3 h_2) + \frac{\alpha^3_w}{m_W^4} \int \text{Tr}(C^\dagger m_d^2 C m_u^2 h_3 C^\dagger m_d^2 h_3 h_2) + \frac{\alpha^3_w}{m_W^4} \int \text{Tr}(C^\dagger m_d^2 C m_u^2 h_3 C^\dagger m_d^2 h_3 h_2) + O(\alpha_w^4) \]

We do not attempt to calculate the two loop contribution, but give qualitative arguments how the enhancement of CP-violating terms appearing in \( h_2 \) result from the one loop calculation.

The thermal propagators for the up quarks \( S(p) \) and the Higgs bosons \( D(p) \) in the Feynman gauge are given by (see (10) for details of the calculation)

\[
S(p) = \frac{(\phi + m_u)}{p^2 - m_u^2 + i\epsilon} + i\Gamma_F(p),
\]

\[
D(p) = \frac{1}{p^2 - m_h^2 + i\epsilon} - i\Gamma_B(p),
\]

with the thermal parts

\[
\Gamma_F = 2\pi\delta(p^2 - m_u^2)f_F(p \cdot u),
\]

\[
\Gamma_B = 2\pi\delta(p^2 - m_h^2)f_B(p \cdot u),
\]

and the fermionic and bosonic distribution functions

\[
f_F(p \cdot u) = \frac{1}{\exp(p \cdot u/T) + 1},
\]

\[
f_B(p \cdot u) = \frac{1}{\exp(p \cdot u/T) - 1}.
\]
The $T = 0$ contributions undergo renormalization and are absorbed into the bare parameters of the Lagrangian. The remaining hermitian terms lead to the following form of $h_2$

$$h_2(\omega, \kappa) = \int \frac{d^4p}{(2\pi)^3} \left( \frac{\Gamma_B(p)}{(p + k)^2 - m_u^2} - \frac{\Gamma_F(p + k)}{p^2 - m_h^2} \right),$$

and after three elementary integrations to

$$h_2(\omega, \kappa) = \frac{1}{\kappa} \int_0^\infty \frac{d|p|}{2\pi} \left( \frac{1}{\epsilon_u} \left| \mathcal{L}_2(\epsilon_u, |p|) f_B(\epsilon_u) \right| - \frac{|p|^2}{\epsilon_u} \mathcal{L}_1(\epsilon_u, |p|) f_F(\epsilon_u) \right),$$

The functions $L_1$ and $L_2$ are defined by

$$L_{1/2}(\epsilon, |p|) = \log \left( \frac{\omega^2 - \kappa^2 \pm \Delta + 2\epsilon \omega + 2\kappa |p|}{\omega^2 - \kappa^2 \pm \Delta + 2\epsilon \omega - 2\kappa |p|} \right) + \log \left( \frac{\omega^2 - \kappa^2 \pm \Delta - 2\epsilon \omega + 2\kappa |p|}{\omega^2 - \kappa^2 \pm \Delta - 2\epsilon \omega - 2\kappa |p|} \right),$$

where $\omega$ and $\kappa$ are the energy and the momentum of the external particle in the rest frame of the plasma, $\epsilon_u = \sqrt{p^2 + m_u^2}$, $\epsilon_u = \sqrt{p^2 + m_u^2}$ and $\Delta = m_u^2 - m_h^2$.

For $|\omega^2 - k^2 - \Delta| > 2|m_u|$ or $|\omega^2 - k^2 + \Delta| > 2|m_u|$ both particles in the loop can be onshell, whilst otherwise not. This makes the function $h_2$ strongly dependent on the two loop masses. In the fig. $\ref{fig:3}$-$\ref{fig:4}$ $h_2$ is plotted as a function of the Higgs vev for a set of fixed external energies $\omega$ and momenta $k$ and three different masses $m_u$ of the quark in the loop. This strong dependence on the Higgs vev $\langle \Phi \rangle$ results in large derivatives of $h_2$ due to the wall profile. In fig. $\ref{fig:3}$ $h_2$ is plotted versus the Higgs vev for different internal quark masses. Note that the first derivative of $h_2$ is in a broad range of parameter space $\{\langle \Phi \rangle, m_u\}$ of order unity or even larger. Furthermore in the limit of vanishing external mass the sign of $h_2'$ changes in the range, where the internal quark mass agrees with the mass of the charged Higgs boson $m_h = m_W = 80$ GeV. The self-energy behaves non-perturbative in the sense that, when expanded in the mass of the internal quark, the main contributions come from higher powers of $m_u/m_W$. In fig. $\ref{fig:4}$ the derivative $h_2'$ is plotted versus the mass of the up quark in the loop. Here it is obvious that the effect is based on a resonance in the loop and can not be increased arbitrarily by increasing the mass of the quark in the loop.

It is reasonable to expect the functions $h_3$ and $h_4$ to be after integration effectively of order one as well and not proportional to unity in flavour space. This allows an estimate of the CP-violating pole dependence of the down quarks

$$\frac{\delta \omega}{\omega} \sim J m_l^4 m_s^3 m_c^2 m_l^2 \frac{\alpha_s}{m_W l_T} \sim 10^{-15},$$

(we have used the Standard Model value $l_W T = 20$ and that most of the particles carry a momentum of the order of the temperature) which is seven orders of magnitude larger than the constraint $\ref{fig:3}$, but still much too small to account for the BAU.
VI. TWO HIGGS DOUBLET MODELS

We have seen, that in spite of the enhancement of the axial current, the CP-violating source due to the CKM matrix is too weak to account for the BAU. Thus we will discuss in this section further possibilities to generate terms of the form (7) in extensions of the SM.

One attractive alternative is the extension to supersymmetric models. Analytical and lattice studies show that the additional scalars in the theory of MSSM may lead to a (two stage, see however) first order phase transition in a part of the parameter space fulfilling the sphaleron bound. However the occuring CP violation in the chargino and neutralino sector has to be maximal and the Higgs mass at the borderline to be seen experimentally to explain the observed BAU. Furthermore the current accuracy in the measurements of electric dipole moments only leaves a small window in parameter space of the MSSM and could rule out this model soon as a source of baryogenesis. An NMSSM type model leaves more freedom at the expense of further parameters.

Besides supersymmetric models, more general two Higgs doublet models are most appealing in extending the SM to explain the BAU via electroweak baryogenesis. In a certain region of parameter space according to Ref. the phase transition is of first order and the sphaleron bound is fulfilled. This derivation assumed the two Higgs vevs to be proportional to each other thus simplifying the calculation while the general case so far has not been completely studied. Baryogenesis is not compatible with this assumption, since if the quotient of the two Higgs vevs is constant, it is not possible to generate an axial fermion current via Ref. Therefore the character of the phase transition has to be examined in every specific model separately.

Since these models are not subject to the stringent restrictions of supersymmetry, there are several possibilities to introduce new sources of CP violation. One feasible approach is to avoid flavour changing neutral currents (FCNC) by construction, which leads to the so called type I and type II models (for a comprehensive discussion see Ref. ). An additional source of CP violation in this context is the complex phase between the two Higgs fields. Similarly to the supersymmetric case this model can, with a reasonable choice of parameters, just marginally explain the generated BAU to be consistent with primordial nucleosynthesis.

Another possibility, even if less attractive because of minor predictivity, is to admit FCNCs at the tree level, called two Higgs doublet models type III. Due to the large parameter space these models still resist to be ruled out by experiments (for some implications on experimental bounds see ) even with quite natural choices for the new parameters and impressive experimental lower bounds on FCNC processes. The rich phenomenology can even account for deviations from the SM as for example the difference between the measurement of the $g-2$ muon factor and it’s SM prediction.

The main difference to models without FCNC and particularly the SM is, that during the electroweak phase transition the derivatives of the mass matrices are not necessarily proportional to the mass matrices themselves. This gives the possibility to construct CP odd rephasing invariants of the form on the tree level and even with just two flavours.

The Lagrangian considered for the the Yukawa couplings of the Higgs fields to the quarks is of the form

$$\mathcal{L}_Y^{(III)} = \eta_{ij}^U \hat{Q}_{i,L} \hat{\bar{u}}_j \hat{D}_{j,R} + \eta_{ij}^D \hat{Q}_{i,L} \phi_1 \hat{D}_{j,R} + \xi_{ij}^U \hat{Q}_{i,L} \tilde{\phi}_2 \bar{u}_j \hat{D}_{j,R} + h.c.,$$

where we used the standard notation: $Q_{i,L}$ denote the left-handed quark doublets, $U_{j,R}$ and $D_{j,R}$ $(i,j = 1, 2, 3)$ the up and down quark singlets, and $\phi_1, \phi_2$ are the two Higgs doublets. To fulfill the experimental bounds, it is sufficient to assume a hierarchy between the couplings...
\[ \eta_{ij}^{U,D} = \frac{m_i \delta_{ij}}{v}, \]
\[ \xi_{ij}^{U,D} = \lambda_{ij} \sqrt{m_i m_j} / v, \]

(14)

\[ |\lambda_{ij}| \lesssim 10^{-1} \]

is needed to suppress \( D^0 - \bar{D}^0 \) and \( B^0 - \bar{B}^0 \) mixing sufficiently. Note that a change in the quotient of the two Higgs \( vevs \) in the mass eigenbasis of the Higgs fields leads to a change in the Yukawa couplings \( \eta \) and \( \xi \) in the above used basis with only one Higgs \( vev \) and therefore to terms of the form \( \eta \xi^\dagger \). The effect can be for example quite large in a two-stage phase transition, as it was seen in Ref. [44, 45]. Baryogenesis in these models occurs at the second phase transition, and it is efficient provided the first phase transition is sufficiently weak, such that the baryon number violating processes are not too suppressed in the weakly broken phase.

The resulting pole shift will be of order

\[ \frac{\delta \omega}{\omega} \approx |\lambda| \frac{m^2}{4K_0 l_w}, \]

and for the top quark this is approximately

\[ \frac{\delta \omega}{\omega} \sim |\lambda| / T l_w, \]

(15) (16)

with \( T \) the temperature and \( l_w \) the wall thickness. The high degree of arbitrariness in these models opens this way a large window for electroweak baryogenesis.

VII. CONCLUSIONS

We have shown that, due to a resonance in the quark self-energies, the CP-violating pole shift induced by the CKM matrix at high temperatures can be by about seven orders of magnitude larger than the CP-violating shift naively expected from the Jarlskog invariant. However, the effect is still too small to account for the BAU via baryogenesis from the Standard Model CP violation at the electroweak phase transition.

Finally we point out that smallness of the CP violation in the Standard Model can be resolved within a certain class of two Higgs doublet models.

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