The Origin and Mechanisms of CP Violation In the Two-Higgs Doublet Model and Masses of the Exotic Scalars

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Abstract

I rebuild a conventional two-Higgs doublet model by relaxing the spontaneous CP violation and considering approximate global U(1) family symmetries. So that the domain-wall problem does not explicitly arise at the weak scale, but CP violation still solely originates from a single CP-phase in the vacuum after spontaneous symmetry breaking. With this phase four types of CP-violating mechanism are induced in the model. In particular, by a new type of the mechanism, both the indirect- and direct- CP violation (i.e. $\epsilon$ and $\epsilon'/\epsilon$) in kaon decay and the neutron electric dipole moment can be consistently accommodated. The masses of the exotic scalars are weakly constrained in the model and searching for these particles is worthwhile in the presently accessible energy range. Substantial CP violation may occur in the heavy quark and lepton sectors and probing their effects provides a challenge at B-factory and colliders.

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One of the simplest extentions of the standard $SU(2)_L \times U(1)_Y$ model \cite{1} is the conventional two-Higgs doublet model (2HDM). Recently, I investigated \cite{2} one of the simplest cases in the 2HDM, that is, CP is broken spontaneously \cite{3} and the neutral currents conserve all flavors \cite{4} at tree level. Consequently such a simple 2HDM, in which Cabibbo-Kobayashi-Maskawa (CKM) matrix \cite{6} is known to be real \cite{7}, possesses a new type of CP-violating mechanism, by which the CP-violating parameters $\epsilon$ \cite{8} and $\epsilon'/\epsilon$ \cite{9} in kaon decay and the neutral electric dipole moment \cite{10} can be consistently accommodated. In particular, the constraint on the masses of the exotic scalars is very weak in that model.

These observations are, on the one hand, of interest and fascinating in physical phenomena, on the other hand they bring forth the necessity to further justify and improve that model. As it is known that there are two essential issues which need to be clarified. They are the domain-wall problem \cite{11} and radiative stability \cite{12}.

Naturally, a simple and interesting way is to further relax the conditions of the spontaneous CP violation (SCPV) and the neutral flavor conservation (NFC) in such a consideration that the improved model remains possessing its initial attractive features in the physical phenomena, and having the basic requirement that CP violation solely originates from the vacuum, namely if vacuum has no CP violation then the theory becomes CP invariant (but CP is not necessary to be broken spontaneously). Indeed, the answer is positive. I shall describe in detail this consideration below.

It is clear that in the limit that CKM matrix is unity, the conventional 2HDM with NFC at tree level generates global U(1) family symmetries and the stability becomes manifest. In the realistic case, it is known that CKM matrix deviates only slightly from unity. This implies that at the electroweak scale any successful models can only possess approximate global U(1) family symmetries (AGUFS).

In general, without imposing any additional conditions, the AGUFS should also play a role on the neutral currents. Therefore it is natural to assume, instead of demanding NFC, a partial conservation of neutral flavor (PCNF). Clearly, the radiative stability now becomes manifest due to the existence of the small terms of the flavor-changing neutral
scalar interactions (FCNSI) at tree level. The smallness of the CKM mixings and FCNSI can be regarded as being naturally in the sense of ’t Hooft’s criterion [13]. It should be mentioned that the motivation that imposing approximate global symmetries on the Yukawa couplings is not novel and has been in fact discussed in the literature. The examples are the approximate flavor symmetries [14] and the approximate discrete symmetries [15]. Nevertheless, different considerations will result in different parameterizations and structures of the Yukawa coupling matrices, consequently the resulting physical phenomena should be also distinguished.

Motivated from the general condition [2] of NFC at tree level, I shall present in this paper an alternative parameterization for the Yukawa coupling matrices by considering the AGUFS and PCNF. It will be seen that this type of parameterization provides a more conventional and general structure for the Yukawa coupling matrices, which will be found to be very useful in classifying various CP violations and analysing possible new physical phenomena. In particular, the suppressions of the FCNSI become more manifest, as a consequence, the constraint on the masses of the exotic scalars can become very weak in this model, so that the masses of these exotic scalars can be just around the present experimental bound.

To prevent the domain-wall problem from arising explicitly in the model, I come to the following observation that

In the gauge theories of spontaneous symmetry breaking (SSB), CP violation can be required solely originating from the vacuum after SSB, even if CP symmetry is not good prior to the symmetry breaking. In other words, there exists a kind of explicit CP violation which can be attributed to the one in the vacuum after SSB, namely if vacuum conserves CP then such an explicit CP violation disappears simultaneously and the theory becomes CP invariant.

In general, the demanded condition for such a statement is: CP nonconservation occurs only at one place of the interactions in the Higgs potential. In particular, this condition may be simply realized by an universal rule that in a renormalizable lagrangian all the interactions with dimension-four conserve CP and only interactions with dimension-two
possess CP nonconservation. It may also be naturally implemented through imposing some symmetries. An interesting example is the Weinberg 3HDM \cite{16} in which one can always choose a basis so that there is only one place allowing to violate CP because of discrete symmetries. To distinguish from the SCPV, such a CP violation may be referred as a Vacuum CP Violation (VCPV).

As an interesting case, let me directly write down in the 2HDM the most general Higgs potential subject to gauge invariance and the universal rule stated above.

\[
V(\phi) = \lambda_1(\phi_1^\dagger \phi_1 - \frac{1}{2}v_1^2)^2 + \lambda_2(\phi_2^\dagger \phi_2 - \frac{1}{2}v_2^2)^2 \\
+ \lambda_3(\phi_1^\dagger \phi_1 - \frac{1}{2}v_1^2)(\phi_2^\dagger \phi_2 - \frac{1}{2}v_2^2) + \lambda_4[(\phi_1^\dagger \phi_1)(\phi_2^\dagger \phi_2) - (\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_1)] \\
+ \frac{1}{2}\lambda_5(\phi_1^\dagger \phi_2 + \phi_2^\dagger \phi_1 - Re(\hat{v}_1^* \hat{v}_2))^2 + \lambda_6(\phi_1^\dagger \phi_2 - \frac{1}{2}(\hat{v}_1^* \hat{v}_2))(\phi_2^\dagger \phi_1 - \frac{1}{2}(\hat{v}_2^* \hat{v}_1)) \\
+ [\lambda_7(\phi_1^\dagger \phi_1 - \frac{1}{2}v_1^2) + \lambda_8(\phi_2^\dagger \phi_2 - \frac{1}{2}v_2^2)][(\phi_2^\dagger \phi_1 + \phi_1^\dagger \phi_2 - Re(\hat{v}_1^* \hat{v}_2)]
\]

where the \(\lambda_i\) (\(i = 1, \cdots, 8\)) are all real parameters. If all the \(\lambda_i\) are non-negative the minimum of the potential then occurs at \(< \phi_i^0 > = \hat{v}_i/\sqrt{2} \equiv v_i e^{i\delta_i}/\sqrt{2}\) with choosing \(\delta_1 = \delta\) and \(\delta_2 = 0\). It is clear that in the above potential CP nonconservation can only occur through the vacuum, namely \(Im(\hat{v}_1^* \hat{v}_2) \neq 0\). Obviously, such a CP violation appears as an explicit one in the potential provided \(\lambda_6 \neq 0\), therefore the domain-wall problem does not explicitly arise. Nevertheless, a further quantitative calculation should be of interest. A similar potential with \(\lambda_7 = \lambda_8 = 0\) due to discrete symmetry was also considered in \cite{17} for strong CP, where discrete symmetry is softly broken by dimension-two terms, its relevant CP violation was referred as a soft CP violation (SOCPV) which is one of the special cases of VCPV and makes sense only for that case. CP-violating effects in our case are no longer soft due to the hard violating terms of discrete symmetry in the Yukawa interactions (in fact no discrete symmetry is imposed). Therefore it is better to refer it as VCPV.

Let me now present a detailed description for the model with VCPV and PCNF by starting with the general Yukawa interactions

\[
L_Y = \bar{q}_L \Gamma_D^a D_R \phi_a + \bar{q}_L \Gamma_U^a U_R \phi_a + \bar{q}_L \Gamma_E^a E_R \phi_a + H.C.
\]
where \(q^i, l^i\) and \(\phi_\alpha\) are \(SU(2)_L\) doublet quarks, leptons and Higgs bosons, while \(U_R^i, D_R^i\) and \(E_R^i\) are \(SU(2)_L\) singlets. \(\Gamma^a_F\) (\(F = U, D, E\)) are the arbitrary real Yukawa coupling matrices. 

\(i = 1, \ldots, n_F\) is a family label and \(a = 1, \ldots, n_H\) is a Higgs doublet label.

I now come to make another essential step, namely to parameterize \(\Gamma^a_F\) in such a conventional way that the global \(U(1)\) family symmetry violations in the charged currents and the neutral currents can be easily distinguished and characterized by the different sets of parameters. Motivated from the general condition \(2\) of NFC at tree level, it is evident to make the following parameterization for \(\Gamma^a_F\)

\[
\Gamma^a_F = O^F_L \sum_{i,j=1}^{n_F} \{ \omega_i (g^F_a \delta_{ij} + \zeta^F \sqrt{g^F} S^F_a \sqrt{g^F_i}) \omega_j \} (O^F_R)^T \tag{3}
\]

\[
g^F_i = | \sum_a g^F_a \hat{v}_a | / (\sum_a | \hat{v}_a |^2)^{1/2}
\]

with \(\{ \omega_i, i = 1, \ldots, n_F \}\) the set of diagonalized projection matrices \((\omega_i)_{jj'} = \delta_{ji} \delta_{j'i}\). \(\hat{v}_a (a = 1, \ldots, n_H)\) are the VEV’s which will develop from the Higgs bosons after SSB. \(g^F_a\) are the arbitrary real Yukawa coupling constants. \(S^F_a = 0\) for \(a = n_H\) and \(S^F_a (a \neq n_H)\) are the arbitrary off-diagonal real matrices. \(g^F_i\) are introduced so that a comparison between the diagonal and off-diagonal matrix elements becomes available. \(\zeta^F\) is a conventional parameter introduced to scale the off-diagonal matrix elements with \((S^F_i)^{12} \equiv 1\) and others \((S^F_a)_{ij}\) being expected to be of order unity (of course, some elements of \(S^F_a\) may be off by a factor of 2 or more). \(O^F_{L,R}\) are the arbitrary orthogonal matrices. Generally, one can choose, by a redefinition of the fermions, a basis so that \(O^F_L = O^F_R \equiv O^F\) and \(O^U = 1\) (or \(O^D = 1\)). The AGUFS and PCNF then imply that

\[
(O^F)^2_{ij} \ll 1 \ , \quad i \neq j \ ; \quad \zeta^2_F \ll 1 \tag{4}
\]

where \(O^F\) describe the AGUFS in the charged currents and \(\zeta_F\) mainly characterizes the PCNF. Obviously, if taking \(\zeta_F = 0\), it turns to the case of NFC at tree level. When \(\zeta_F = 0\) and \(O^U = O^D = 1\), the lagrangian then possesses global \(U(1)\) family symmetries, i.e., \((U, D)_i \rightarrow e^{\alpha_i} (U, D)_i\).

For the simplest 2HDM, the physical basis after SSB is defined through

\[
f_L = (O^F_L V^F_L)^\dagger F_L
\]
and \( f_R = (O_R^f P_f V_R^f)^T F_R \) with \( V_{L,R}^f \) being unitary matrices and introduced to diagonalize the mass matrices

\[
(V_L^f)^\dagger (\sum_i m_i^o \omega_i + \zeta F c_\beta \sum_{i,j} \sqrt{m_i^o} \omega_i S_i^F \omega_j \sqrt{m_j^o} e^{i\sigma_f (\delta - \delta_f)}) V_R^f = \sum_i m_i^o \omega_i
\]

with \( m_i \) the masses of the physical states \( f_i = u, d, e \). Where the following definitions are introduced

\[
(c_\beta g_1^F e^{i\sigma_f \delta} + s_\beta g_2^F) v \equiv \sqrt{2} m_i^o e^{i\sigma_f \delta_f},
\]

with \( v^2 = v_1^2 + v_2^2 = (\sqrt{2} G_F)^{-1} \), \( g^F v = \sqrt{2} m_i^o \) and \( c_\beta \equiv \cos \beta = v_1/v \) and \( s_\beta \equiv \sin \beta = v_2/v \).

Where \( P_{ij}^f = e^{i\sigma_f \delta_i \delta_j} \), with \( \sigma_f = + \) for \( f = d, e \), and \( \sigma_f = - \) for \( f = u \).

As a convention, writing \( V_{L,R}^f \equiv 1 + \zeta_f T_{L,R}^f \), thus the scalar interactions of the fermions can be written in the physical basis into

\[
L_Y \equiv L_Y^o + L_Y' (\zeta_f)
\]

where \( L_Y^o \) possesses no FCNSI and has a simple structure

\[
L_Y^o = (2 \sqrt{2} G_F)^{1/2} \sum_{i,j} \{ \xi_{d,j} \bar{u}_L^i V_{ij}^o m_{d,j} d_R^i H^+ - \xi_{u,j} \bar{d}_L^i V_{ij}^o m_{u,j} u_R^i H^- + \xi_{e,j} \bar{v}_L^i \delta_{ij} m_{e,j} e_R^j H^+ + H.C. \} + (\sqrt{2} G_F)^{1/2} \sum_{k} \sum_{i=1}^{3} \eta_{u,i}^{(k)} m_{u,i} \bar{u}_L^i u_R^i + \eta_{d,i}^{(k)} m_{d,i} \bar{d}_L^i d_R^i + \eta_{e,i}^{(k)} m_{e,i} \bar{e}_L^i e_R^i + H.C. \} H_k^0
\]

and

\[
\xi_{f_i} = \frac{\sin \delta f_i}{\delta_\beta^2 \sin \delta} e^{i\sigma_f (\delta - \delta_f)} \tan \beta - \cot \beta ; \quad \eta_{f_i}^{(k)} = O_{2k}^H + (O_{1k}^H + i \sigma_f O_{3k}^H) \xi_{f_i}
\]

where \( V^o = (O_L^f)^T O_L^D \) is real. \( H_k^0 = (h, H, A) \) are the three physical neutral scalars and \( O_{kl}^H \) is the \( 3 \times 3 \) orthogonal mixing matrix among these three scalars, \( H \) plays the role of the Higgs boson in the standard model.

\( L_Y' \) contains FCNSI, its complete expression and physical phenomena will be presented in a longer paper [20]. I mention here only their main features. Firstly, it is in favor of having
\[
\tan \beta > 1 \quad \text{and} \quad \sin \delta f_i / \sin \delta \lesssim 1 \quad \text{in order to have the FCNSI be suppressed manifestly.}
\]

To the first order in \( \zeta_F \) the couplings of the FCNSI are of order \( O(\zeta_F s_\beta^{-1} \sqrt{2} G_F m_{f_i} m_{f_j})^{1/2} (S_{ij}^F) \). When \( \zeta_D \lesssim 10^{-3} m_{H^0} (GeV) / (1 GeV) \) the masses of the exotic scalars are unconstrained from the \( K^0 - \bar{K}^0 \) and \( B^0 - \bar{B}^0 \) mixings. For a typical value \( \zeta_D \sim V_{cb} \sim 0.04 \), the masses of the exotic scalars are allowed to be just around the present experimental bound \( m_H \sim 48 GeV \).

Secondly, the CP-violating parameter \( \epsilon \) and the mass difference between the \( K_L \) and \( K_S \) may be accommodated by FCNSI through fine-tuning the CP phases \( \delta, \delta_s, \delta_d \) and the phases \( \arg(O_{1k}^H + i O_{3k}^H) \) arising from the scalar-pseudoscalar mixings (SPM). But its contribution to \( \epsilon'/\epsilon \) is in general small. Finally, the smallness of CP violation becomes natural in the induced KM-type mechanism. This is because in the present model CP-phase in the KM-matrix \( V = (V^u_L)^\dagger V^o V^d_L \) is directly related to FCNSI. In a good approximation, to the first order of \( \zeta_F T_{L,R}^f \) the amplitudes of the imaginary part in \( V \) are of order \( O(\zeta_F c_\beta (m_{f_i} / m_{f_j})^{1/2} (S_{ij}^F)) \) with \( i < j \).

In general, there appear four types of CP-violating mechanism which are induced from the single CP-phase \( \delta \) in the vacuum. Three of them (i.e. FCNSI, KM-type and SPM) have been mentioned above. Let me now concentrate in this short paper on the new type of CP violation mechanism in \( L^\phi \) provided \( \zeta_F < 10^{-3} \) so that the effects from the FCNSI and KM-type mechanisms are negligible. Without making additional assumptions, \( m_{f_i}, V^o_{ij}, m_{H^0_k}, \xi_{f_i} \text{ (or } \delta_{f_i}) \), \( m_{H^+} \) and \( O_{kl}^H \text{ (or } \eta_{f_i}^{(k)}) \) are in general all the free parameters and will be determined only by the experiments. An important feature of the present model is that rich CP-violating phases \( (\delta_{f_i}) \) are induced from the single CP-phase \( \delta \). They are in general all different and observable. Clearly, it is unlike the Weinberg 3HDM \[16\] in which it is equivalent to the case that \( \xi_d = \xi_s = \xi_b \) and \( \xi_u = \xi_c = \xi_t \), therefore the difficulties encountered \[18,19\] in the Weinberg 3HDM do not arise in the present model.

Let us first check the CP-violating parameters \( \epsilon \) and \( \epsilon'/\epsilon \) in kaon decays. Like the Weinberg 3HDM, but unlike the KM-model, the long-distance contribution to \( \epsilon \) become important. The imaginary part of the \( K^0 - \bar{K}^0 \) mixing mass matrix to which \( \epsilon \) is proportional mainly receives the contribution from the \( \pi, \eta \) and \( \eta' \) poles \[21\]. Following the analyses in
the literature \[21,22,19\], we have in our present model
\[ \epsilon \simeq 2.27 \times 10^{-3} \Im(\xi_d \xi_c - 0.05 \xi_d^* \xi_c^*) \frac{37 \text{GeV}^2}{m_H^2} (\ln \frac{m_H^2}{m_c^2} - \frac{3}{2}) \]
which should easily accommodate the experimental value \(|\epsilon| = 2.27 \times 10^{-3}\) since \(\xi_s\) and \(\xi_c\) are all free parameters. It will be shown \[20\] that the short-distance box graphs with charged-scalar in our model can also provide a large contribution to \(\epsilon\).

The ratio \(\epsilon'/\epsilon\) in the Higgs-boson-exchange models was first correctly estimated by Donoghue and Holstein \[22\]. The evaluations can directly be applied to the present model. In fact, it was shown in the refs. \[22,19\] that the ratio does not depend on the detail of the CP-odd matrix element and is given by
\[ \epsilon'/\epsilon = 0.017 D = (0.4 - 6.0) \times 10^{-3} \]
which is comparable to the one calculated from the standard KM-model \[23\] and is also consistent with the present experimental data \[9\]. A short-distance contribution to the ratio from charged-scalar exchange at tree level is also found to be of order \(10^{-3}\) \[20\].

Consider now the neutron EDM, \(d_n\). The present experimental limit on \(d_n\) is \(d_n < 1.2 \times 10^{-25} \text{ e cm}\) \[10\]. Applying various well-known scenarios for the calculations of \(d_n\) to the present model, the \(d_n\) is then accommodated by choosing the parameters \(\xi_{fi}\) and \(\eta_{fi}^{(k)}\) to satisfy the following conditions
\[ D1 : \quad \Im(\xi_d \xi_c) \frac{1}{m_H^2} (\ln \frac{m_H^2}{m_c^2} - \frac{3}{4}) \lesssim 6.3 \times 10^{-2} \text{GeV}^{-2} , \]
\[ D2 : \quad \Im[\eta_{l}^{(k)}] h_{NH}(m_t, m_{H_0}) \lesssim 0.18 , \]
\[ D3 : \quad \Im(\xi_b \xi_t) h_{CH}(m_t, m_b, m_H) \lesssim 3.0 \times 10^{-2} , \]
\[ D4 : \quad \Im(\eta_{d}^{(k)} \eta_{l}^{(k)} + 0.5 \eta_{u}^{(k)} \eta_{l}^{(k)}) \lesssim 0.2 , \quad (m_t \sim m_H) \]
where the condition \(D1\) is from the quark model through charged-Higgs boson exchange, \(D2\) and \(D3\) from the Weinberg’s gluonic operator through the neutral-Higgs boson exchange \[18\] and the charged-Higgs boson exchange \[24\] respectively. \(h_{NH}\) and \(h_{CH}\) are the functions of the quark- and Higgs-mass arising from the integral of the loop. \(h_{NH} \simeq 0.05\) for \(m_t = m_H\).
\( h_C \leq 1/8 \) and \( h_C = 1/12 \) for \( m_b \ll m_t = m_H \). \( D4 \) is from the gluonic chromoelectric dipole moment (CEDM) \([23]\) induced by the Barr-Zee two-loop mechanism \([26]\).

For the present experimental bound of \( m_{H^+) \gtrsim 45 \text{GeV} \), it is not difficult to find that \( \text{Im} \xi_s \xi_c \gtrsim 10 \) and \( \text{Im} \xi_d \xi_c \sim 21 \). A natural solution for them is \( \tan \beta \gg 1 \), i.e., \( v_2 \gg v_1 \), such a hierarchy was already discussed in \([27]\).

The hierarchic structure of the VEV’s also implies that large CP violation may occur in the heavy quark and lepton sectors due to their possible large complex Yukawa couplings. One of the interesting examples is the T-odd and CP-odd triple momentum correlations in the exclusive semileptonic B-meson decay \( B \to D^*(D\pi)\tau \nu_\tau \) \([28]\). Using the analysis of the ref. \([28]\) to the present model, the CP asymmetry reads

\[
A_{VT} \approx 10^{-2} \frac{m_b m_\tau}{m_{H^+}^2} \text{Im}(\xi_b \xi_c^*) \lesssim 10^{-2} - 10^{-3} \tag{16}
\]

where the last numerical values are estimated by taking \( |\xi_b| \sim |\xi_\tau| \sim (v_2/v_1) \sim (m_t/m_b) \), \( m_{H^+} \sim m_W \) and \( \sin(\delta_\tau - \delta_b) \sim 1.0 - 0.1 \).

Before ending this paper, I would like to briefly remark that the strong CP problem may be simply evaded by using the well-known Peccei-Quinn mechanism \([29]\) realized in a heavy fermion invisible axion scheme \([30]\). Since this scheme, on the one hand, is one of the simplest schemes and has an advantage that it leaves the above model almost unchanged, and on the other hand it is also free from the axion domain-wall problems.

In a word, the above considerations indicate that the conventional 2HDM with VCPV and PCNF may provide one of the simplest and attractive models in understanding the origin and mechanisms of CP violation. It also opens a window for probing new physical phenomena at B-factory and colliders. In particular, the weak constraint on the masses of the exotic scalars makes the searching for these particles in the presently accessible energy range more worthwhile. Various interesting features arising from this model are going to be discussed in a longer paper \([20]\).

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