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# Cp violation in the MSSM and role of the KM phase

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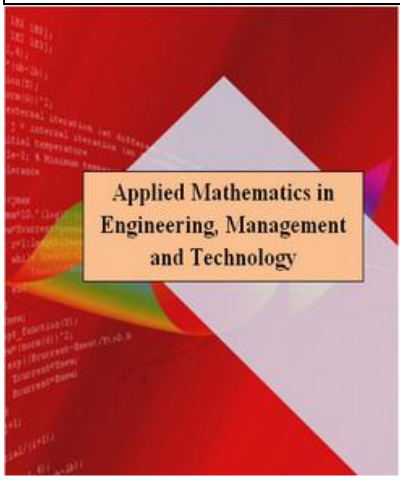


**Abstract**

Another aspect of SM extensions, namely non-standard CP violation, is also essential for baryogenesis scenarios.SM extensions, in particular, an extended non-gauge sector; that is to say, a richer set of Yukawa and Higgs-boson self-interactions than in the SM. It is these interactions that break, in general, CP invariance. Thus, in SM extensions additional sources of CP violation besides the KM phase are usually present. In the minimal super symmetric extension (MSSM) of the Standard Model CP violating phases can appear, apart from the complex Yukawa interactions of the quarks yielding a non-zero KM phase $\delta_{KM}$.

## 1. **introduction**

The standard model of particle physics combined with the SCM has, it seems, all the ingredients for generating a baryon asymmetry. First we recall the salient features of the SM at temperatures T $\simeq$ 0 which apply to present-day physics. The observed particle spectrum tells us that the electroweak gauge symmetry $SU(2)_L \times U(1)_Y$, for which there is solid empirical evidence, cannot be a symmetry of the ground state. In the SM this spontaneous symmetry breaking is accomplished by a $SU(2)_L$ doublet of scalar fields $\Phi(x)$, the Higgs field, that is assumed to have a non-zero ground state expectation value $\langle 0|\Phi|0\rangle = 246Gev$. This classical field selects a direction in the internal $SU(2)_L \times U(1)_Y$ space and hence breaks the electroweak symmetry, leaving intact the gauge symmetry of electromagnetism. The W and Z bosons, quarks, and leptons acquire their masses by coupling to this field (which may be viewed as a Lorentz-invariant ether).
C and CP are violated by the charged weak quark interactions

$$\mathcal{L}_{cc} = -\frac{g_w}{\sqrt{2}}\overline{U}_L\gamma^\mu V_{KM} D_L W_\mu^+ + h.c. \quad (1)$$

Here $U_L = (u_L\,,\ c_L, t_L\ )^T$, $D_L = (d_L\,,\ s_L, b_L\ )^T$, denote the left-handed quark fields ($q_L = \frac{(1-\gamma_5)q}{2}$), $W_\mu^+$is the W boson field, $g_w$ is the weak gauge coupling, and $V_{KM}$ is the Cabibbo-Kobayashi-Maskawa mixing matrix. CP is violated if the KM phase angle $\delta_{KM} \neq 0, \pm\pi$. By this "mechanism" the CP effects observed so far in the K and B meson systems can be explained.

There is also baryon number violation in the SM, but this is a subtle, nonperturbative effect which is completely negligible for particle reactions in the laboratories at present-day collision energies, but very significant for the physics of the early universe.

$\mathcal{L}_{SM}$ is invariant under the following two sets of global phase transformations of the quark and lepton fields q = u, ..., t; ℓ = e, ..., ντ :

$$q(x) \to e^{iw/3} q(x) \qquad \ell(x) \to \ell(x) \quad (2)$$

$$\ell(x) \to e^{i\lambda}\ell(x) \qquad q(x) \to q(x) \quad (3)$$

Applying Noether's theorem we obtain the associated symmetry currents $J_\mu^B$ μ and $J_\mu^L$, which are conserved at the Born level:

$$\partial^\mu J_\mu^B = \partial^\mu \sum_q \frac{1}{3} \bar{q}\gamma_\mu q = 0 \quad (4)$$

$$\partial^\mu J_\mu^L = \partial^\mu \sum_\ell \bar{\ell}\gamma_\mu \ell = 0 \quad (5)$$

(The currents are to be normal-ordered.) Thus the associated charge operators

$$\hat{B} = \int d^3x J_0^B(x) \quad (6)$$

$$\hat{L} = \int d^3x J_0^L(x) \quad (7)$$

are time-independent. At the level of quantum fluctuations beyond the Born approximation these symmetries are, however, explicitly broken because eqs. (4),
(5) no longer hold. This is seen as follows. Decompose the vector current

$$\bar{f}\gamma_\mu f = \bar{f}_L\gamma_\mu f_L + \bar{f}_R\gamma_\mu f_R \quad (8)$$

where $f = q, \ell$, into its left- and right-handed pieces. Because of the clash between gauge and chiral symmetry at the quantum level the gauge-invariant chiral currents are not conserved: in the quantum theory the current-divergencies suffer from the Adler-Bell-Jackiw anomaly[1], [2]. For a gauge theory based on a gauge group G, which is a simple Lie group of dimension $d_G$ , the anomaly equations for the L- and R-chiral currents $\bar{f}_L\gamma_\mu f_L$and $\bar{f}_R\gamma_\mu f_R$read

$$\partial^\mu \bar{f}_L\gamma_\mu f_L = -c_L \frac{g^2}{32\pi^2} F_{\mu\nu}^a \tilde{F}^{a\mu\nu} \quad (9)$$

$$\partial^\mu \bar{f}_R\gamma_\mu f_R = +c_R \frac{g^2}{32\pi^2} F_{\mu\nu}^a \tilde{F}^{a\mu\nu} \quad (10)$$

where $F_{\mu\nu}^a$ is the (non)abelian field strength tensor ($a = 1, \dots, d_G$) and $\tilde{F}^{a\mu\nu} = \epsilon^{\mu\nu\alpha\beta} \frac{F_{\alpha\beta}^a}{2}$ is the dual tensor[1], g denotes the gauge coupling, and the constants $c_L$, $c_R$ depend on the representation which the $f_L$ and $f_R$ form. Let us apply (8)- (10) to the above baryon and lepton number currents of the SM where the gauge group is $SU(3)_c \times SU(2)_L \times U(1)_Y$. Because gluons couple to right-handed and lefthanded quark currents with the same strength, we have $c_L^{QCD} = c_R^{QCD}$.Therefore $J_\mu^B$ has no QCD anomaly. However, the weak gauge bosons $W_\mu^a, a = 1,2,3$, couple only to left-handed quarks and leptons, while the weak hypercharge boson couples to $f_L$ and $f_R$.with different strength. Hence $c_R^W = 0$, $c_L^Y \neq c_R^Y$.
Putting everything together one obtains

$$\partial^\mu J_\mu^B = \partial^\mu J_\mu^L = \frac{n_F}{32\pi^2} (-g_w^2 W_{\mu\nu}^a \tilde{W}^{a\mu\nu} + g'^2 B_{\mu\nu} \tilde{B}^{\mu\nu}) \quad (11)$$

where $W_{\mu\nu}^a$ and $B_{\mu\nu}$ denote the $SU(2)_L$and $U(1)_Y$field strength tensors, respectively, $g'$ is the $U(1)_Y$ gauge coupling, and $n_F = 3$ is the number of generations.

Eq. (11) implies that $\partial^\mu (J_\mu^B - J_\mu^L) = 0$. Thus the difference of the baryonic and leptonic charge operators $\hat{B} - \hat{L}$ remains time-independent also at the quantum level and therefore the quantum number B - L is conserved in the SM.

## 2. **EW Phase Transition in SM Extensions**

Of course, whether or not the SM Higgs field or some other mechanism provides the correct description of EW symmetry breaking remains to be clarified. In fact, this is the most important unsolved problem of present-day particle physics. Future collider experiments hope to resolve this issue. On the theoretical side, a number of extensions and alternatives to the SM Higgs mechanism have been discussed for quite some time. One may distinguish between models which postulate elementary Higgs fields (i.e., the associated spin-zero particles have pointlike couplings up to some high energy scale E $\gg$ 100 GeV) which trigger the breakdown of $SU(2)_L \times U(1)_Y$, and others which assume that it is caused by the Bose condensation of (new) heavy fermion-antifermion pairs. The dynamics of the symmetry breaking sector of these models can change the order of the EW phase transition, as compared with the SM. Let's briefly discuss results for some models that belong to the first class. The presently most popular extensions of the SM are supersymmetric (SUSY) extensions, in particular the minimal supersymmetric standard model (MSSM), the Higgs sector of which contains two Higgs doublets. Although the requirement of SUSY breaking to be soft does not allow for independent quartic couplings in the Higgs potential $V(\Phi_1, \Phi_2)$, the number of parameters of the scalar sector of this model is larger than that of the SM and a first order transition can be arranged. Investigations of $V_{eff}$ at $T \neq 0$ show that there is a region in the MSSM parameter space which allows for a sufficiently strong first order EW phase transition (see, for instance, the reviews [3], [4] and references therein). The condition for this is that the mass of the scalar partner $\tilde{t}_R$ of the right-handed top quark $t_R$ must be sufficiently light and the mass of $\tilde{t}_L$ must be sufficiently heavy. An upper bound on the mass of the lightest neutral Higgs boson $H_1$ of the model obtains from the requirement that the mass of $\tilde{t}_L$ should not be unnaturally large. In summary, the MSSM predicts a sufficiently strong 1st order EW phase transition if

$$m_{H_1} \leq 105 - 115\, Gev\,, \quad m_{\tilde{t}_R} \leq 170 \qquad (12)$$

In the next-to-minimal SUSY model which contains an additional gauge singlet Higgs field a strong first order transition can be arranged quite easily.

Non-supersymmetric SM extensions may be, in general, less motivated than SUSY models, but several of these models are, nevertheless, worth to be studied as they predict interesting phenomena. For illustrative purposes we mention here only the class of 2 Higgs doublet models (2HDM) where the field content of the SM is extended by an additional Higgs doublet, leading to a physical particle spectrum which includes 3 neutral and one charged Higgs particle. The general, renormalizable and $SU(2)_L \times U(1)_Y$ invariant Higgs potential $V(\Phi_1, \Phi_2)$contains a large number of unknown parameters. Therefore, it is not surprising that in these models, too, the requirement of a strong 1st order EW transition can be arranged quite easily as studies of the finite-temperature effective potential show no tight upper bound on the mass of the lightest Higgs boson obtains[5].

## 3. **Higgs sector CPV**

An interesting possibility is CP violation (CPV) by an extended Higgs sector which can occur already in the 2-Higgs doublet extensions of the SM. Consider the class of 2HDM which are constructed such that flavour-changing neutral (pseudo)scalar currents are absent at tree level. The appropriate $SU(2)_L \times U(1)_Y$invariant treelevel Higgs potential $SU(2)_L \times U(1)_Y$ of these models may be represented in the following way:

$$V_{tree}(\Phi_1, \Phi_2) = \lambda_1\left(2\Phi_1^\dagger\Phi_1 - v_1^2\right)^2 + \lambda_2\left(2\Phi_2^\dagger\Phi_2 - v_2^2\right)^2 + \lambda_3\left[\left(2\Phi_1^\dagger\Phi_1 - v_1^2\right) + \left(2\Phi_2^\dagger\Phi_2 - v_2^2\right)\right] + \lambda_4\left[\left((\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) - (\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1)\right)\right] + \lambda_5\left[2Re(\Phi_1^\dagger\Phi_2) - v_1 v_2 \cos\zeta\right]^2 + \lambda_6\left[2Im(\Phi_1^\dagger\Phi_2) - v_1 v_2 \sin\xi\right]^2$$
(13)

where $\lambda_i$, $v_1$‘$v_2$and $\xi$ are real parameters and the parameterization of $V_{tree}$ is chosen such that the Higgs fields have non-zero VEVs in the state of minimal energy.

Performing a CP transformation,

$$\Phi_{1,2}(x,t) \overset{CP}{\rightarrow} e^{i\alpha_{1,2}}\Phi_{1,2}^\dagger(-x,t) \qquad (14)$$

we see that $H_V = \int d^3x V_{tree}(\Phi_1, \Phi_2)$ is CP-no invariant if $\xi \neq 0$. Notice that it is unnatural to assume $\xi = 0$. Even if this was so at tree level, the non-zero KM phase $\delta_{KM}$, which is needed to explain the observed CPV in K and B meson decays, would induce a non-zero $\xi$ through radiative corrections.
From eq. (13) we read off that at zero temperature the neutral components of the Higgs doublet fields have, in the electric charge conserving ground state, the expectation values

$$\langle 0|\phi_1^0|0\rangle = v_1 \frac{e^{\xi_1}}{\sqrt{2}}, \quad \langle 0|\phi_2^0|0\rangle = v_2 \frac{e^{\xi_2}}{\sqrt{2}} \tag{15}$$

where $v = \sqrt{v_1 + v_2} = 246\ Gev$, and $\xi = \xi_2 - \xi_1$ is the physical CPV phase.
The spectrum of physical Higgs boson states of the two-doublet models consists of a charged Higgs boson and its antiparticle, $H^{\pm}$, and three neutral states. As far as CPV is concerned, $H^{\pm}$ carries the KM phase. This particle affects the (CPV) phenomenology of flavor-changing $|\Delta F| = 2$ neutral meson mixing and $|\Delta F| = 1$ weak decays of mesons and baryons. ( Experimental data on $b \to s + \gamma$ imply that this particle must be quite heavy, $m_{H^+} > 210\ Gev$.)
Let's briefly discuss some implications of Higgs sector CPV for present-day physics. If $\xi$ were zero, the set of neutral Higgs boson states would consist of two scalar (CP=1) and one pseudo scalar (CP= –1) state. If $\xi \neq 0$ these states mix. As a consequence the 3 mass eigenstates, $|\varphi_{1,2,3}\rangle$, no longer have a definite CP parity. That is, they couple both to scalar and to pseudo scalar quark and lepton currents.
In terms of Weyl fields the corresponding Lagrangian reads

$$\mathcal{L}_{\varphi} = -\sum_{\psi} c_{\psi} \frac{m_{\psi}}{v} \bar{\psi}_L \psi_R \varphi + h.c. \tag{16}$$

The sum over the Higgs fields i = 1, 2, 3 is implicit, $\psi$ denotes a quark or lepton field, $m_{\psi}$ is the mass of the associated particle, and the dimensionless reduced Yukawa couplings $c_w = a_w + ib_w$ depend on the parameters of the Higgs potential and on the type of model.
The Yukawa interaction leads to CPV in flavour-diagonal reactions for quarks and for leptons $\psi$. The induced CP effects are proportional to some power $(m_{\psi})^p$. For example, consider the reaction $\bar{\psi}\psi \to \psi\bar{\psi}$. The exchange of a $\varphi$ boson at tree level induces an effective CPV interaction of the form $(\bar{\psi}\psi)(\bar{\psi}i\gamma_5\psi)$ with a coupling strength proportional to $\frac{m_{\psi}^2}{m_{\varphi}^2}$. The search for non-zero electric dipole moments (EDM) of the electron and the neutron has traditionally been a sensitive experimental method to trace non-SM CP violation. If a light $\varphi$ boson exists and the CPV phase ξ is of order 1 the Yukawa interaction can induce electron and neutron EDMs of the same order of magnitude as their present experimental upper bounds.
We assume that the parameters of the 2HDM are such that the transition is strongly first order. Moreover, in order to simplify the discussion we assume that the passage from the symmetric to the broken phase occurs in one step, at some temperature $T_c$. Somewhat below $T_c$ bubbles filled with Higgs fields start to nucleate and expand. That is, the Higgs VEVs are space and time dependent. Let's consider, for simplicity, only one of the bubbles and assume its expansion to be spherically symmetric. When the bubble has grown to some finite size we can use the following one-dimensional description. Consider the rest frame of the bubble wall. The wall is taken to be planar and the expansion of the bubble is taken along the z axis. The wall, i.e., the phase boundary has some finite thickness $l_{wall}$, extending from $z = 0$ to $z = z_0$. The symmetric phase lies to the right of this boundary, $z > z_0$ while the broken phase lies to the left, z < 0. Thus the neutral Higgs fields have VEVs whose magnitudes and phases vary with z:

$$\langle 0|\phi_1^0|0\rangle_T = \frac{\rho_1(z)}{\sqrt{2}} e^{i\theta(z)}, \quad \langle 0|\phi_2^0|0\rangle_T = \frac{\rho_2(z)}{\sqrt{2}} e^{iw(z)} \tag{17}$$

In the symmetric phase, $z \gg z_0$, both VEVs vanish, whereas in the broken phase the VEVs should be close to their zero temperature values:

$$\rho_i(z) \cong v_i\ , \quad \theta(z) \cong \xi_1\ , \qquad \omega(z) \cong \xi_2 \tag{18}$$

if $z \gg z_0$.The variation of the moduli and phases with z can be determined by solving the field equations of motion that involve the finite-temperature effective
potential of the model.
As to the couplings of the Higgs fields to fermions, we assume here and in the
following subsection, for definiteness, that all quarks and leptons couple to $\Phi_1$ only. Then the Yukawa coupling of a quark or lepton field $\psi = q, \ell$ to the neutral Higgs field is given by

$$\begin{aligned}\mathcal{L}_1 &= -h_\psi \bar{\psi}_L \psi_R \phi_1^0 + h.c.\\ &= -m_\psi(z)\bar{\psi}_L\psi_R - m_\psi^*(z)\bar{\psi}_R\psi_L + \cdots\ ,\end{aligned} \tag{19}$$

Where

$$m_\psi(z) = h_\psi \frac{\rho_1(z)}{\sqrt{2}} e^{i\theta(z)} \tag{20}$$

is a complex-valued mass and the ellipses in (19) indicate the coupling of the quantum field, i.e., the coupling of a neutral Higgs particle to $\psi$. Thus the interaction of a fermion field $\psi(x)$ with the CP-violating Higgs bubble, treated as an external, classical background field, is summarized by the Lagrangian

$$\mathcal{L}_\psi = \bar{\psi}_L i\gamma^\mu \partial_\mu \psi_L + \bar{\psi}_R i\gamma^\mu \partial_\mu \psi_R - m_\psi\ \ (x)\bar{\psi}_L\psi_R - m_\psi^*\ \ (x)\bar{\psi}_R\psi_L \tag{21}$$

## 4. **CP Violation in the MSSM**

The requirement of gauge invariance and hermiticity of the Lagrangian (21) allows for the following new sources of CP violation:
i) A complex mass parameter $\mu_c = \mu \exp(i\varphi_\mu)$, $\mu$ real, describing the mixing of the two Higgs chiral super fields in the superpotential.
ii) A complex squared mass parameter $m_{12}^2$ describing the mixing of the two Higgs doublets and contributes to the Higgs potential

$$V(\Phi_1, \Phi_2) \supset \mu_c \Phi_1^\dagger \cdot \Phi_2^\dagger + h.c\ , \tag{22}$$

iii) Complex Majorana masses $M_i$in the gaugino mass terms($\epsilon \equiv i\sigma_2$),

$$-\sum_i M_i \frac{(\lambda_i^T \epsilon \lambda_i)}{2} + h.c.\ , \tag{23}$$

Where i = 1, 2, 3 refers to the $U(1)_Y$ ,$SU(2)_L$ gauginos, and gluinos, respectively. A standard assumption is that the $M_i$ have a common phase.
iv) Complex trilinear scalar couplings of the scalar quarks and scalar leptons, respectively, to the Higgs doublets $\Phi_1, \Phi_2$. These couplings form complex 3×3 matrices $A_\psi$ in generation space. Motivated by supergravity models it is often assumed that the matrices $A_\psi$ are proportional to the Yukawa coupling matrices $h_\psi$:

$$A_\psi = A h_\psi\ , \psi = u, d, \ell, \tag{24}$$

where A is a complex mass parameter.
Thus the parameter set $\mu_c, m_{12}^2, M_i$ , and A involves 4 complex phases. Exploiting

two (softly broken) global $U(1)$ symmetries of the MSSM Lagrangian, two of these phases can be removed by re-phasing of the fields. A common choice, we we shall also use, is a phase convention for the fields such that the gaugino masses $M_i$ and the mass parameter $m_{12}^2$ are real. Then the observable CP phases in the MSSM (besides the KM phase) are $\varphi_\mu = \arg(\mu_c)$ and $\varphi_\mu = \arg(\mu_c)$. The experimental upper bounds on the electric dipole moments $d_e, d_n$ of the electron and the neutron put, however, rather tight constraints on these CP phases, in particular on $\varphi_\mu$. Even if there are correlations between these phases such that there are cancellations among the contributions to $d_e$ and to $d_n$, finds(see also [6],[7]) that $\varphi_\mu$ is constrained by the data to be smaller than $|\varphi_\mu| \leq 0.03$. A way out of this constraint would be heavy first and second generation sleptons and squarks with masses of order 1 TeV.

In the MSSM the tree-level Higgs potential $V_{tree}$ is CP-invariant. Supersymmetry does not allow for independent quartic couplings in $Y$. They are proportional to linear combinations of the $SU(2)_L$ and $U(1)_Y$ gauge couplings squared. At one-loop order the interactions of the Higgs fields $\Phi_1, \Phi_2$ with charginos, neutralinos, (s)tops, etc. generate quartic Higgs self-interactions of the form

$$V_{eff} \supset \lambda_1(\Phi_1^\dagger\Phi_2)^2 + \lambda_2(\Phi_1^\dagger\Phi_2)(\Phi_1^\dagger\Phi_1) + \lambda_3(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_2) + h.c., \qquad (25)$$

In the effective potential. The CP phases $\varphi_\mu$ and $\varphi_A$ induce complex $\lambda_{1,2,3}$. Thus, explicit CP violation in the Higgs sector occurs at the quantum level which leads to
Yukawa interactions of the neutral Higgs bosons being of the form (16).

In the context of baryogenesis a potentially more interesting possibility is spontaneous CP violation at high temperatures $T \leq T_{EW}$.This kind of CP violation could not be traced any more in the laboratory! Ref. [8] pointed out that, irrespective of whether or not $\varphi_\mu$ and $\varphi_A$ are sizeable, the MSSM effective potential receives, at high temperatures $T \leq T_{EW}$, quite large one-loop corrections of the form (25). As a consequence, the neutral Higgs fields can develop complex VEVs of the form (17) with a large CP-odd classical field. This would signify spontaneous CPV at finite temperatures, even if $\varphi_\mu$ and $\varphi_A$ would be very small or even zero. However, ref [9] finds that experimental constraints on the parameters of the MSSM and the requirement of the phase transition to be strongly first order preclude this possibility in the case of the MSSM.

Let's now come to those CP-violating interactions of the MSSM which are of
relevance at the EW phase transition and involve $\varphi_\mu$ and $\varphi_A$ at the tree level. As discussed above, there is a small, phenomenologically acceptable range of light Higgs and light stop mass parameters which allows for a strong first order transition. The Higgs VEVs are of the form

$$\langle 0|\phi_1^0|0\rangle_T = \rho_1(z)\ , \quad \langle 0|\phi_2^0|0\rangle_T = \rho_2(z) \qquad (26)$$

where $\rho_i$ are real and for convenience, a normalization convention different from the one in (17) is used here.

These VEVs determine the interaction of the bubble wall with those MSSM particles that couple to the Higgs fields already at the classical level. Inspecting where the CP-violating phases $\varphi_\mu$ and $\varphi_A$ are located in LMSSM (we use the convention of the gaugino masses $M_i$ being real) it becomes clear that the relevant interactions of the classical Higgs background fields are those with charginos, neutralinos, and sfermions, in particular top squarks. Contrary to the case of the 2HDM discussed above the interactions of quarks and leptons with a bubble wall do not – at the classical level – violate CP invariance if (26) applies.

Inserting (26) into the respective terms of the MSSM Lagrangian we obtain the
Lagrangians describing the particle propagation in the presence of a Higgs bubble
[10], [11]. For the charged gauginos and Higgsinos in the gauge eigenstate basis
we get

$$\mathcal{L}_c = \chi_R^\dagger\sigma_\mu\partial^\mu\chi_R + \chi_L^\dagger\bar{\sigma}_\mu\partial^\mu\chi_L + \chi_R^\dagger\mathcal{M}_c\chi_L + \chi_L^\dagger\mathcal{M}_c^\dagger\chi_R \qquad (27)$$

Where $\sigma\mu = (I, \sigma i), \sigma\mu = (I, -\sigma i)$, and we have put

$$\chi_R^\dagger = (\widetilde{W}^+, \widetilde{H}_2^\dagger), \quad \chi_L = (\widetilde{W}^-, \widetilde{H}_1^-)^T, \qquad (28)$$

where $\widetilde{W}(x)$, $\widetilde{H}_{1,2}(x)$ are 2-component Weyl fields for the charged gauginos and Higgsinos, respectively. The chargino mass matrix is given by

$$\mathcal{M}_c(x) = \begin{pmatrix} M_2 & g_w\rho_2(x) \\ g_w\rho_1(x) & \mu_c \end{pmatrix}, \qquad (29)$$

Where $\mu_c$ is the complex Higgsino mass parameter defined above.

For the scalar stop fields $\tilde{t}_R(x), \tilde{t}_L(x)$ we obtain in the gauge eigenstate basis

$$\mathcal{L}_{\tilde{t}} = (\partial_\mu, \tilde{t}_L^\dagger)\partial^\mu \tilde{t}_L + (\partial_\mu, \tilde{t}_R^\dagger)\partial^\mu \tilde{t}_R - (\tilde{t}_L^\dagger\, , \tilde{t}_R^\dagger\, )\mathcal{M}_{\tilde{t}} \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix}, \qquad (30)$$

With

$$\mathcal{M}_{\tilde{t}}(x) = \begin{pmatrix} m_L^2 + h_t^2 H_2^2(x) & h_t(A_t\rho_2(x) - \mu_c^*\rho_1(x)) \\ h_t(A_t^*\rho_2(x) - \mu_c\rho_1(x)) & m_R^2 + h_t^2\rho_2^2(x) \end{pmatrix}, \qquad (31)$$

Where $m_L^2$ are SUSY breaking squared mass parameters, $h_t$ is the top-quark Yukawa coupling, and $A_t$ is the left-right stop mixing parameter.
In the mass matrices (29) and (31) the CP-violating phases combine with the spatially varying VEVs and will give rise to x-dependent CP-violating phases when the mass matrices are diagonalized, analogously to the case of the 2HDM above. This causes CP-violating particle currents.

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