On the spontaneous CP breaking in the Higgs sector of the Minimal Supersymmetric Standard Model *

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Abstract

We revise a recently proposed mechanism for spontaneous CP breaking at finite temperature in the Higgs sector of the Minimal Supersymmetric Standard Model, based on the contribution of squarks, charginos and neutralinos to the one-loop effective potential. We have included plasma effects for all bosons and added the contribution of neutral scalar and charged Higgses. While the former have little effect, the latter provides very strong extra constraints on the parameter space and change drastically the previous results. We find that CP can be spontaneously broken at the critical temperature of the electroweak phase transition without any fine-tuning in the parameter space.

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1. The issue of spontaneous CP breaking (SCPB) at zero temperature has been recently proposed in supersymmetric theories as an appealing alternative to the Peccei-Quinn mechanism to solve the strong CP problem, and also as a natural explanation for the smallness of CP violating phases contributing to the neutron electron dipole moment (NEDM). In the minimal supersymmetric standard model (MSSM), SCPB at zero temperature can be triggered by radiative corrections. Unfortunately, as established on general grounds by Georgi and Pais, it requires the existence of a Higgs boson with a mass of a few GeV, which has been ruled out at LEP.

On the other hand, CP violation at finite temperature is one of the key ingredients to generate baryon asymmetry at the electroweak phase transition. It has been recently realized that temperature effects can trigger SCPB in the MSSM at the critical temperature of the electroweak phase transition, and baryon asymmetry with the help of tiny CP violating phases giving rise to a NEDM well below its present experimental limit. The analysis of Refs. was based on the contribution of squarks, charginos and neutralinos to the one-loop effective potential at finite temperature.

In this paper we have analyzed SCPB in the MSSM at finite temperature, taking into account plasma effects for all bosons and including the contribution of the Higgs sector to the effective potential. We have found that plasma effects have little influence on the SCPB, but the contribution from neutral scalar and charged Higgses changes drastically the results found in Ref. giving rise to a much stronger constraint on the parameter space. Nevertheless we find that SCPB can be accomplished without any fine-tuning of the parameters.

2. The most general scalar potential for the two-Higgs doublet model that is renormalizable and $SU(2) \times U(1)$ invariant is given by

$$V = m_1^2|H_1|^2 + m_2^2|H_2|^2 - (m_3^2H_1H_2 + h.c.) + \lambda_1|H_1|^4 + \lambda_2|H_2|^4 + \lambda_3|H_1|^2|H_2|^2 + \lambda_4 H_1^2 H_2^2 + \lambda_5 (H_1 H_2)^2 + \lambda_6 |H_1|^2 H_1 H_2 + \lambda_7 |H_2|^2 H_2 H_1 + h.c.),$$

where $m_3^2$, $\lambda_5$, $\lambda_6$ and $\lambda_7$ are taken real, so that there is no explicit CP violation at tree level. However, SCPB can be triggered if the vacuum expectation values of the neutral components of the Higgs fields get a relative phase $\delta \neq n\pi$, with integer $n$. (We can take $\langle H_1^0 \rangle = v_1$, $\langle H_2^0 \rangle = v_2 e^{i\delta}$ without any loss of generality.) This requires,

$$\lambda_5 > 0,$$

$$-1 < \cos \delta = \frac{m_3^2 - \lambda_6 v_1^2 + \lambda_7 v_2^2}{4\lambda_5 v_1 v_2} < 1.$$  

In the MSSM these conditions are not satisfied (at tree level) because in this case supersymmetry gives

$$\lambda_1 = \lambda_2 = \frac{1}{8}(g^2 + g'^2),$$

$$\lambda_3 = \frac{1}{4}(g^2 - g'^2),$$

$$\lambda_4 = -\frac{1}{2}g^2,$$

$$\lambda_5 = \lambda_6 = \lambda_7 = 0,$$
where \( g \) and \( g' \) are the \( SU(2)_L \) and \( U(1)_Y \) gauge couplings. Of course, after supersymmetry breaking, one–loop effects change these couplings and there are finite contributions (proportional to the soft parameters) that can produce SCPB. At finite temperature, there will be additional corrections coming from the thermal modes running in the loops.

We can write the one–loop effective potential, including the zero temperature terms, as

\[
\Delta V_1 = T^4 \left( \sum_b g_b V_b - \sum_f g_f V_f \right),
\]

where \( g_b \) \((g_f)\) denotes the number of degrees of freedom of bosons \( (\text{fermions}) \). In the bosonic part we sum over stops, neutral scalar and charged Higgses, and Goldstone bosons \( (b = \tilde{t}_{L,R}, h, H, H^\pm, G^\pm) \), and in the fermionic one we consider neutralinos and charginos \( (f = \tilde{\chi}^0, \tilde{\chi}^\pm) \). Working in the \( 't \) Hooft–Landau gauge and in the DR renormalization scheme, we will decompose the bosonic components as

\[
V_b = V_{nz}(y_b^2) + V_z(\bar{y}_b^2).
\]

The bosonic zero temperature part plus the finite temperature contribution coming from \( n \neq 0 \) Matsubara modes \[11\],

\[
V_{nz}(y_b^2) = Tr \left\{ \frac{y_b^4}{64\pi^2} \left[ \log \left( \frac{T^2 y_b^2}{Q^2} \right) - \frac{3}{2} \right] + \frac{1}{2\pi^2} J_b(y_b^2) \right\},
\]

where \( y_b^2 = \mathcal{M}_b^2/T^2 \) is the bosonic mass matrix rescaled with temperature, \( Q \) is the renormalization scale, and the finite temperature one–loop term is given by the usual integral expression with the \( n = 0 \) part subtracted out,

\[
J_b(y_b^2) = \int_0^\infty dx \ x^2 \log \left( 1 - e^{-\sqrt{x^2+y_b^2}} \right) - \left( -\frac{\pi^3}{6} y_b^3 \right).
\]

The \( n = 0 \) term is given by \[11\]

\[
V_z(\bar{y}_b^2) = Tr \left\{ -\frac{\bar{y}_b^3}{12\pi} \right\},
\]

and now \( \bar{y}_b^2 = (\mathcal{M}_b^2 + \Pi_b)/T^2 \), where \( \Pi_b \) is the thermal polarization matrix for bosons, contains the Debye masses \[12\].

As the fermions do not have zero Matsubara modes, \( V_f \) is just given by \[11\]

\[
V_f = Tr \left\{ \frac{y_f^4}{64\pi^2} \left[ \log \left( \frac{T^2 y_f^2}{Q^2} \right) - \frac{3}{2} \right] + \frac{1}{2\pi^2} J_f(y_f^2) \right\},
\]

with

\[
J_f(y_f^2) = \int_0^\infty dx \ x^2 \log \left( 1 + e^{-\sqrt{x^2+y_f^2}} \right),
\]

and \( y_f^2 = \mathcal{M}_f^2/T^2 \), where \( \mathcal{M}_f^2 \) is the fermionic mass matrix.

The bosonic and fermionic mass matrices are functions of the classical fields \( H_1, H_2, \tilde{H}_1^o \) and \( \tilde{H}_2^o \), and one can expand the potential (4) in powers of them to obtain the one–loop corrections to the tree level couplings of Eq. (1). Assuming as in Ref. \[8\] equal soft
supersymmetry breaking masses for left and right-handed squarks $^1 m^2_Q = m^2_{\tilde{t}} \equiv \bar{m}^2$ one can easily find the contributions from stop loops to $m^2_3$, $\lambda_5$, $\lambda_6$ and $\lambda_7$: 

\[ \Delta^{(s)}m^2_3 = -6h^2_i A_t \mu I'_i, \]  
\[ \Delta^{(s)}\lambda_5 = \frac{1}{2} h^4_i A_t^2 \mu^2 T^{-4} I''_i, \]  
\[ \Delta^{(s)}\lambda_6 = h^2_i A_t \mu \left[ \frac{3}{4}(g^2 + g'^2) I''_i + h^2_i \mu^2 T^{-4} I''_i \right], \]  
\[ \Delta^{(s)}\lambda_7 = h^2_i A_t^2 \mu \left[ (6h^2_i - \frac{3}{4}(g^2 + g'^2)) I''_i + h^2_i A_t^2 I''_i \right], \]  

where $\mu$ is the supersymmetric Higgs mixing mass, $A_t$ the trilinear soft breaking parameter corresponding to $h_t Q \cdot H_2 U$ in the superpotential, 

\[ I_b \equiv I + I_{nz} \equiv \frac{d}{dy_b} V_2(y_b^2) + \frac{d}{dy_b} V_{nz}(y_b^2), \]  
\[ I'_b \equiv I' + I'_{nz} \equiv \frac{d}{dy_b} I_2(y_b^2) + \frac{d}{dy_b} I_{nz}(y_b^2) \]  

(and a similar definition for higher derivatives), and now $y_b^2 = \bar{m}^2/T^2$, $\bar{y}_i^2 = (\bar{m}^2 + \Pi_i)/T^2$, with $\Pi_i \approx \frac{2}{3} g_s^2 T^2$.

where $g_s$ is the $SU(3)$ gauge coupling. We are neglecting in Eq. (17) contributions of order $g^2, g'^2$ and $h^2$ compared to the $g_s^2$ part. In this way the field independent mass matrix for stops with the thermal screening included is proportional to the identity matrix and the expressions (11-14) apply.

The mass matrices of neutral and charged Higgses at the origin, $H^0_1 = H^0_2 = 0$, are never proportional to the identity and so the corresponding one–loop corrections involve finite differences instead of derivatives of the $I_b$ functions. In particular,

\[ \Delta^{(h)}m^2_3 = -\frac{1}{2} \left\{ (g^2 + g'^2) + 2g^2 \right\} m^2_3 \Delta I_h, \]  
\[ \Delta^{(h)}\lambda_5 = \frac{1}{48} m^4_3 T^4 \left\{ (g^2 + g'^2)^2 + 2g^4 \right\} \Delta_3 I_h, \]  
\[ \Delta^{(h)}\lambda_6 = \frac{1}{8} m^2_3 T^2 \left\{ (g^2 + g'^2)^2 + 2g^4 \right\} \Delta I'_h \]  
\[ + \left\{ (g^2 + g'^2)^2 + 2g^2 g'^2 \right\} \frac{m^2_1 - m^2_2}{3 T^2} \Delta_3 I_h \]  
\[ \Delta^{(h)}\lambda_7 = \frac{1}{8} m^2_3 T^2 \left\{ (g^2 + g'^2)^2 + 2g^4 \right\} \Delta I'_h \]  

\[ ^1\text{This is just a technical simplification such that the mass matrix at the origin, } H^0_1 = H^0_2 = 0, \text{ is proportional to the identity matrix.} \]

\[ ^2\text{The high-}T \text{ expansion of Eqs. (11-14) coincides with Eqs. (15-18) in Ref. [8], except that we find 1/512 instead of 1/256 in Eq. (16) of [8].} \]
where the contribution to $\Delta^{(h)}$ proportional to $(g^2 + g'^2)$ inside the curly brackets comes from the neutral scalar Higgs sector, and the rest, proportional to $g^2$, from the charged Higgs sector. $\Delta I_h$ and $\Delta_3 I_h$ in Eqs. (18-21) are defined by

$$
\Delta I_h \equiv \frac{I_n(y^2_1) - I_n(z^2_1)}{y^2_1 - y^2_2} + \frac{I_n(y^2_2) - I_n(z^2_2)}{y^2_2 - y^2_1},
$$

$$
\Delta_3 I_h \equiv \frac{6}{(y^2_3 - y^2_2)^2} \left[ I_n(y^2_3) + I_n(z^2_3) - 2I_n(y^2_2) + I_n(z^2_2) - 2I_n(z^2_1) - I_n(y^2_1) \right],
$$

and $y^2_\pm = m^2_\pm/T^2$, $y^2_\mp = \tilde{m}^2_\mp/T^2$ as usual, with $m^2_\pm$ the Higgs masses for zero classical fields:

$$
m^2_\pm = \frac{1}{2} \left\{ m_1^2 + m_2^2 \pm \sqrt{(m_2^2 - m_1^2)^2 + 4m_3^4} \right\},
$$

and $\tilde{m}^2_\pm$ the corresponding masses in the presence of screening:

$$
\tilde{m}^2_\pm = \frac{1}{2} \left\{ m_1^2 + m_2^2 \pm \sqrt{(m_2^2 - \tilde{m}_1^2)^2 + 4\tilde{m}_3^4} \right\},
$$

where $m_1^2$ and $m_2^2$ are functions of $m^2_3$, tan $\beta$ and the other parameters of the theory, by means of the zero-temperature (including the one-loop corrections) minimization conditions, and [13, 14]

$$
\tilde{m}_1^2 = m_1^2 + \Pi_{h_1} = m_1^2 + \frac{1}{8}(3g^2 + g'^2)T^2,
$$

$$
\tilde{m}_2^2 = m_2^2 + \Pi_{h_2} = m_2^2 + \frac{1}{8}(3g^2 + g'^2 + 6h_1^2)T^2.
$$

We can also assume that $\mu^2 = M_1^2 = M_2^2$ (where $M_2$ and $M_1$ are the soft Majorana masses for $SU(2)_L$ and $U(1)_Y$ gauginos, respectively) and then find, for charginos:

$$
\Delta^{(c)} m^2_3 = -\text{sign}(\mu)4g^2\mu^2 I'_\chi,
$$

$$
\Delta^{(c)} \lambda_5 = -\frac{1}{3}g^4 T^4 \frac{\mu^4}{I'^2 \chi},
$$

$$
\Delta^{(c)} \lambda_6 = \Delta^{(c)} \lambda_7 = \text{sign}(\mu)4g^2 \frac{\mu^2}{3 T^2} \left[ 3I'_\chi + \frac{\mu^2}{T^2} I'' \chi \right],
$$

and for neutralinos [14]

$$
\Delta^{(n)} m^2_3 = \frac{g^2 + g'^2}{2g^2} \Delta^{(c)} m^2_3,
$$

$$
\Delta^{(n)} \lambda_{5,6,7} = \frac{(g^2 + g'^2)^2}{2g^4} \Delta^{(c)} \lambda_{5,6,7},
$$

We correct here Eq. (26) of Ref. [14] where the proportionality coefficient of $\Delta^{(c)} \lambda_{5,6,7}$ was written as $(g^2 + g'^2)/2g^2$. 

\[3\]
where now $y_\chi^2 = \mu^2/T^2$ and all the integrals are fermionic. Derivatives are defined, as usual,

$$I_f = V_f' \equiv \frac{d}{dy_f^2} V_f(y_f^2). \quad (32)$$

3. We analyze now the effect described in the previous section. The parameter $m_3^2$ will be traded by $m_A^2$, the physical mass of the pseudoscalar which includes all radiative corrections (at zero-temperature). None of conditions (2) is satisfied at tree level (in particular, $\lambda_5 = 0$ and $\delta = n\pi$), neither after including one-loop radiative corrections at zero-temperature, if we fix $m_A$ beyond its present experimental limit [5]. However, they can be satisfied when radiative corrections at finite-temperature are included. In particular, $\lambda_a$, $(a = 5, 6, 7)$ acquires a renormalized value at finite temperature

$$\lambda_a(T) = \sum_{i=s,h,c,n} \Delta^{(i)} \lambda_a \quad (33)$$

as given by Eqs. (12-14), (19-21), (28-29) and (31). Also $m_3^2$ gets renormalized as

$$m_3^2(T) = m_3^2 + \sum_{i=s,h,c,n} \Delta^{(i)} m_3^2 \quad (34)$$

from Eqs. (11), (18), (27) and (30). We have computed, in (33) and (34), $\Delta^{(s,c,n)}$ numerically, from the integrals (7) and (10), while $\Delta^{(h)}$ is computed using its finite temperature expansion, because in the Higgs sector the corresponding value of the integral (7) is well defined only through its high-$T$ expansion.

A glance at the previous section (Eqs. (12), (19), (28) and (31)) shows that $\Delta^{(s)} \lambda_5 \leq 0$, while $\Delta^{(c,n)} \lambda_5 \geq 0$, so that the contribution of fermions (charginos and neutralinos) have to compensate that of squarks, if no other contributions were considered, so that one would obtain in the plane $(\mu, \tilde{m})$ a lower bound contour, as in Ref. [8]. We have analyzed the effects of the Debye screening on the latter contour and found it gets shifted towards smaller ($\sim 5\%$) values of $\tilde{m}$. We can conclude from this that plasma corrections in the squark sector are not very important. Now the contribution from the Higgs sector (where plasma effects are essential to make the effective potential at $H_1^0 = H_2^0 = 0$ real) is negative, and has an infrared singularity ($\propto -(T/\bar{m}_-)$), responsible for the failure of perturbative expansion [11], for values of $(\mu, \bar{m})$ such that $\bar{m}_- = 0$ (long-dashed lines in Figs. 1-3). This creates in the plane $(\mu, \bar{m})$ an upper bound contour which severely constrains the region of the parameter space where $\lambda_5 > 0$. This behaviour is exemplified in Figs.1-3 where $\lambda_5 = 0$ is plotted in the plane $(\mu, \bar{m})$ for different values of $T$, $m_A$ and the supersymmetric parameters [8]. We can see from Figs. 1b, 2b and 3b that the lower and upper contours provide a combined region where $\lambda_5 > 0$. In the absence of mixing $A_t = 0$, Figs. 1a, 2a and 3a, there is no contribution coming from the squark sector and, consequently, the lower contour disappears. Only the upper contour (from the Higgs sector) constrains the region where $\lambda_5 > 0$. If we compare

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4We have checked that the theory remains perturbative, i.e. that $\bar{\beta}_{\pm} \equiv g^2 T/\bar{m}_{\pm} < 1$, in the region where $\lambda_5 > 0$ in Figs. 1-3.
the region where \( \lambda_5(T) > 0 \) for different temperatures, e.g. Figs. 1 and 2, we see it decreases with decreasing \( T \), and for any given point in the plane \((\mu, \bar{m})\) the condition \( \lambda_5 > 0 \) will stop being satisfied at a given temperature.

In the region where \( \lambda_5(T) > 0 \) the condition \( |\cos \delta(T)| < 1 \) where

\[
\cos \delta(T) = \frac{m_0^2(T) - \lambda_6(T)v_1^2(T) - \lambda_7(T)v_2^2(T)}{4\lambda_5(T)v_1(T)v_2(T)} \tag{35}
\]

is also required to have SCPB. This condition is satisfied for values of \( v(T) \) \((v^2(T) = v_1^2(T) + v_2^2(T))\) such that \( v_-(T) \leq v(T) \leq v_+(T), \) where

\[
v_{\pm}(T) = \frac{m_0^2(T)}{\cos^2 \beta(T)} \left\{ \lambda_6(T) + \lambda_7(T) \tan^2 \beta(T) \mp 4\lambda_5(T) \tan \beta(T) \right\}^{-1} \tag{36}
\]

For a given point in the plane \((\mu, \bar{m})\) the values of \( v_{\pm}(T) \) are very close to each other and depend on \( \tan \beta(T) \). For a first-order phase transition the values of \( \tan \beta(T) \) and \( \langle v(T) \rangle \) at the critical temperature are dynamically fixed by all the parameters of the supersymmetric theory. In the absence of a complete analysis of the phase transition in the MSSM \([14]\), we will take \( \tan \beta(T) = \tan \beta \), which we know is a good approximation at least for large values of \( m_A \), to get an estimate of the values of \( v_{\pm} \) for the different points in the plane \((\mu, \bar{m})\). At very high temperatures, \( \langle v(T) \rangle = 0 \) and the condition (2) will not be satisfied in general for a fixed point in the plane \((\mu, \bar{m})\). At the critical temperature \( T_c \) the field will go from the origin to \( \langle v(T_c) \rangle \) and produce a maximal CP violation \((\Delta \delta = \pm \pi)\) if \( 0 < v_-(T_c) < v_+(T_c) < \langle v(T_c) \rangle \) for the chosen point in the plane \((\mu, \bar{m})\).

We have plotted in Figs. 1-3 level contours corresponding to different values of \( 0 \leq v_-(T) \leq T \). If the phase transition is strong enough, i.e. \( \langle v(T) \rangle \sim T \), then the whole region with \( v_-(T) \leq T \) and \( \lambda_5 > 0 \) satisfies the condition of CP violation. If the phase transition is stronger (weaker), i.e. \( \langle v(T) \rangle > T \) \((\langle v(T) \rangle < T)\), then the allowed region showed in Figs. 1-3 will be further enlarged (reduced). For values of \( \langle \tan \beta(T) \rangle \) different from \( \tan \beta \), the allowed regions in Figs. 1-3 would get distorted, but still a wide region would appear.

4. In conclusion we have estimated the region of maximal CP violation at the critical temperature of the electroweak phase transition in the space of supersymmetric parameters. We have studied for simplicity the degenerate case of \( m_Q = m_U \) in the stop sector, and \( M_1 = M_2 = \mu \) in the gaugino/higgsino sector. We have found that the Higgs sector changes dramatically previous results on the same effect. Nevertheless the mechanism can work in a wide region of the parameter space without any fine-tuning, and trigger easily maximal CP violation during the first order phase transition. However this requires, as a general tendency, small values of \( m_A \) and large values of \( \tan \beta \), while preliminary results on the phase transition in the MSSM \([13, 14]\) seem to point towards the opposite direction in order not to wash out the generated baryon asymmetry. The analysis of the non-degenerate case \( m_Q \neq m_U, M_1 \neq M_2 \neq \mu \) (now in progress) seems necessary to find out whether this mechanism can be realistically used to generate baryon asymmetry at the phase transition in the MSSM.
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**Fig. 1** 1a) Contours of \( \bar{m}_- = 0 \) (long-dashed line), \( \lambda_5 = 0 \) (horizontal solid line), \( v_-(T)/T = 1 \) (outer vertical solid lines) and \( v_-(T)/T = 0 \) (inner vertical solid lines) for \( m_t = 160 \text{ GeV}, \ T = 150 \text{ GeV}, \ m_A = 50 \text{ GeV}, \tan \beta = 8 \) and \( A_t = 0 \); 1b) The same as in 1a) but for \( A_t = 50 \text{ GeV} \). Here the contour of \( \lambda_5 = 0 \) is closed on the left and the contours of \( v_-(T)/T = 1 \) (outer vertical solid lines) and \( v_-(T)/T = 0.6 \) (short-dashed line) are shown.

**Fig. 2** The same as in Fig. 1 but with \( T = 100 \text{ GeV} \) and \( \tan \beta = 18 \). In 2b) the short-dashed line corresponds to \( v_-(T)/T = 0.7 \).

**Fig. 3** The same as in Fig. 1 but for \( m_A = 75 \text{ GeV} \) and \( \tan \beta = 18 \). In 3b) the short-dashed line corresponds to \( v_-(T)/T = 0.5 \).