SQUARKS AND SPHALERONS

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Abstract

Electric charge and color breaking minima along third generation squark directions do appear in the Minimal Supersymmetric Standard Model for particular regions of the corresponding supersymmetric parameters $A_t$, $\mu$, $\tan \beta$, $m_Q^2$ and $m_U^2$. We have studied possible instabilities of electroweak sphalerons along non-trivial squark configurations. We have found that instabilities along the latter lie in the region of parameter space where the electroweak minimum is not the global one. Thus charge and color conservation imply that the standard electroweak sphaleron is not destabilized along squark directions.
1. The existence of sphalerons (static and unstable solutions to classical field equations) in the SU(2) gauge theory with Higgs fields in the fundamental representation has been known since long ago [1], triggering the hope of generating the baryon asymmetry of the Universe [2] at the electroweak phase transition [3]. The value of the sphaleron energy, an essential ingredient for baryogenesis mechanisms [4, 5, 6, 7], has been the object of detailed numerical calculations in the Standard Model [8, 9, 10], and translates into an upper bound on the Higgs boson mass well below the experimental lower limit. This bound comes from the weakness of the phase transition in the Standard Model.

The minimal supersymmetric extension of the Standard Model (MSSM) [11] is a very appealing candidate, both from the theoretical [1] and the experimental [2] side, to describe physics at the electroweak scale. The strength of the phase transition in the MSSM has been extensively studied in the literature [12, 13, 14]. Recently it has been proven that the phase transition can be considerably strengthened in the MSSM in the presence of light supersymmetric partners of the right-handed top quark (stops) and small values of $\tan \beta$ [15, 16, 17], opening thus the window for baryogenesis at the electroweak phase transition. These results have been recently confirmed by an explicit calculation of the sphaleron solution and the sphaleron energy in the MSSM, at zero and finite temperature, including all relevant one-loop radiative corrections [18].

The sphaleron solution presented in Ref. [18] is the natural generalization of the Standard Model sphaleron solution and involves the two Higgs doublets and the SU(2) gauge fields. However, the main feature of supersymmetric models is the existence of a plethora of new scalar fields, the sfermions, which might affect the sphaleron solution if non-trivial configurations turn out to be energetically favoured. In particular, for the region of supersymmetric parameters where the phase transition is strengthened, there are electric charge and color breaking (CCB) [19] local minima along the left-handed and/or right-handed stop directions, which might destabilize the standard sphaleron in the MSSM and then endanger previous results based on standard sphalerons and the analysis of the phase transition.

In this letter will study the effect of possible non-trivial sfermion configurations on the sphaleron solution and its energy. We find that as far as electric charge and color are not broken along squark directions, i.e. as far as the standard electroweak minimum along the neutral component of the Higgs fields remains as the global minimum, the trivial (zero) configuration for sfermion fields is always energetically favoured.

2. Neglecting sfermion mixing between different generations (as suggested by constraints from flavour changing neutral current processes), we can simplify the problem by considering just one generation of squarks and sleptons. On the other hand, as a first approximation, we will fix to zero the slepton fields. The consistency of this fixing is protected by a global symmetry and will be justified a posteriori by our results. In fact we will work with third generation squarks, whose large Yukawa couplings would allow tunneling [20] from the symmetric vacuum at temperatures above the electroweak phase transition temperature, or from the standard non-symmetric electroweak vacuum, for

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1Yielding a 'technical' solution to the hierarchy problem.
2LEP precision measurements leading to gauge coupling unification.
temperatures below, to a possible lower CCB vacuum. We will also work in the approximation of taking \( g_1 = 0 \) so that the \( U(1)_Y \) gauge field \( B_\mu \) can be consistently set to zero. This allows a spherically symmetric ansatz. So, we are looking for a sphaleron-like solution described by

\[
W_\mu, \quad H_1 = \begin{bmatrix} H_1^0 \\ H_1^1 \end{bmatrix}, \quad H_2 = \begin{bmatrix} H_2^+ \\ H_2^0 \end{bmatrix}
\]

and

\[
Q_L = \begin{bmatrix} U_L \\ D_L \end{bmatrix}, \quad U^c_R, \quad D^c_R.
\]

The lagrangian density for these fields is given by

\[
\mathcal{L} = -\frac{1}{4} W^a_{\mu\nu} W^{a\mu\nu} + (D_\mu H_1)^\dagger (D^\mu H_1) + (D_\mu H_2)^\dagger (D^\mu H_2) + (D_\mu Q_L)^\dagger (D^\mu Q_L)
\]

\[
+ (\partial_\mu U^c_R)^\dagger (\partial^\mu U^c_R) + (\partial_\mu D^c_R)^\dagger (\partial^\mu D^c_R) - V_{\text{eff}}(H_1, H_2, Q_L, U^c_R, D^c_R)
\]

where \( W^a_{\mu\nu} \) stands for the \( SU(2) \) field strength and \( V_{\text{eff}} \) is the effective potential. We can expand the effective potential (at zero temperature) as

\[
V_{\text{eff}} = V_0(H_1, H_2, Q_L, U^c_R, D^c_R) + V_1(H_1, H_2, Q_L, U^c_R, D^c_R) + \cdots
\]

where \( V_0 \) is the tree-level potential and \( V_1 \) contains the one-loop radiative corrections. The tree-level potential can be decomposed as:

\[
V_0 = V_{\text{SU(2)}}^D + V_{\text{SU(3)}}^D + V_F + V_{\text{soft}}
\]

with

\[
V_{\text{SU(2)}}^D = \frac{g_2^2}{8} \left[ (H_1^\dagger H_1)^2 + (H_2^\dagger H_2)^2 + (Q_L^\dagger Q_L)^2
\right.
\]

\[
- 2 \left( (H_1^\dagger H_1)(H_2^\dagger H_2) + (H_1^\dagger H_1)(Q_L^\dagger Q_L) + (H_2^\dagger H_2)(Q_L^\dagger Q_L) \right)
\]

\[
+ 4 \left( |H_1^\dagger H_2|^2 + |H_1^\dagger Q_L|^2 + |H_2^\dagger Q_L|^2 \right]
\]

and

\[
V_{\text{SU(3)}}^D = \frac{g_2^2}{2} \left[ \frac{1}{3} (Q_L^\dagger Q_L)^2 + \frac{1}{3} (U^c_R U^c_R)^2 + \frac{1}{3} (D^c_R D^c_R)^2
\right.
\]

\[
- (Q_L U^c_R)^\dagger (Q_L U^c_R) + \frac{1}{3} (Q_L^\dagger Q_L)(U^c_R U^c_R)
\]

\[
- (Q_L D^c_R)^\dagger (Q_L D^c_R) + \frac{1}{3} (Q_L^\dagger Q_L)(D^c_R D^c_R)
\]

\[
+ (U^c_R D^c_R)^\dagger (U^c_R D^c_R) - \frac{1}{3} (U^c_R U^c_R)(D^c_R D^c_R)
\]
the SU(2) and SU(3) D-terms. $V_F$ is determined from the superpotential

$$W = h_t Q_L \cdot H_2 U_R^c + h_b Q_L \cdot H_1 D_R^c + \mu H_1 \cdot H_2,$$

and reads:

$$V_F = |h_b Q_L D_R^c - \mu H_2|^2 + |h_t Q_L U_R^c + \mu H_1|^2$$

$$+ |h_t H_2 U_R^c + h_b H_1 D_R^c|^2 + h_t^2 |Q_L \cdot H_2|^2 + h_b^2 |Q_L \cdot H_1|^2$$

The soft-breaking terms are:

$$V_{soft} = m_1^2 H_1^\dagger H_1 + m_2^2 H_2^\dagger H_2 + m_3^2 (H_1 \cdot H_2 + h.c.)$$

$$+ m_Q^2 Q_L^\dagger Q_L + m_U^2 U_R^\dagger U_R^c + m_D^2 D_R^\dagger D_R^c$$

$$+ (h_t A_t Q_L \cdot H_2 U_R^c + h_b A_b Q_L \cdot H_1 D_R^c + h.c.)$$

where $m_1^2$, $m_2^2$ and $m_3^2$, can be traded in favour of the supersymmetric parameters, tan $\beta$ and $m_A$, using the minimization conditions of the radiatively corrected effective potential. Finally, for the one-loop effective potential one should use its expression in the DR \[21\] renormalization scheme:

$$V_1 = \frac{1}{64 \pi^2} \text{Str} \left[ M^4 \left( \log \frac{M^2}{Q^2} - \frac{3}{2} \right) \right]$$

where $Q$ is the renormalization scale. We have taken the approximation where only Higgs background fields are considered in the one-loop correction, $V_1 = V_1(H_1, H_2)$. This correction can be absorbed in a redefinition of the tree-level parameters in the Higgs sector \[22\] and hence provides large corrections to Higgs boson masses. On the other hand the tree-level potential along the squark directions is dominated by the strong $g_3$ and top-quark $h_t$ Yukawa couplings. Since the minima of the potential will always be located at field values of the order of the weak scale, we do not expect large radiative corrections for them and we can safely put to zero the background fields $Q_L$, $U_R^c$ and $D_R^c$ in (11).

Choosing the temporal, radial gauge we can write the static, spherically symmetric ansatz

$$W_a^q(\vec{x}) = \frac{2[1 - f(r)]}{g_2 r^2} \epsilon_{ajk} x_k$$

$$H_1(\vec{x}) = \bar{h}_1(r) \begin{bmatrix} 1 \\ 0 \end{bmatrix}$$

$$H_2(\vec{x}) = \bar{h}_2(r) \begin{bmatrix} 0 \\ 1 \end{bmatrix}$$

$$Q_L(\vec{x}) = \bar{q}(r) \begin{bmatrix} 1 \\ 0 \end{bmatrix}$$

$$U_R^c(\vec{x}) = \bar{u}(r)$$

$$D_R^c(\vec{x}) = \bar{d}(r)$$

$$\bar{u}(r)$$

$$\bar{d}(r)$$
for the sphaleron, where \( r^2 = x^2 + y^2 + z^2 \) and all functions are real. Notice that we have not included gluons, while \( Q_L \) and \( U^c_R \) are aligned on a unique SU(3) direction, which provides a self consistent ansatz.

Replacing the ansatz \([14]\) in \([4]\) we obtain in a straightforward fashion the scalar potential as a function of the background fields \( \tilde{h}_1, \tilde{h}_2, \tilde{q} \) and \( \tilde{u} \):

\[
V_{\text{eff}} = V_{\text{eff}}(\tilde{h}_1, \tilde{h}_2, \tilde{q}, \tilde{u}).
\]

We have explored the region where the supersymmetric parameters

\[
(m_Q, m_U, m_A, \tan \beta, A_t, \mu)
\]

can take all values allowed by experimental and theoretical constraints. In particular large values of \( m_Q \) are preferred by LEP precision measurements, while large values of \( m_A \), \( m_A \sim m_Q \), small values of \( \tan \beta \), \( \tan \beta \lesssim 3 \), and moderately negative values of \( m_U^2 \), \( m_U^2 > -m_Q^2 \), are favoured by the strength of the electroweak phase transition, and will be focused on in the following discussion. This region has been recently proven to appear naturally from the usual radiative scenario of electroweak symmetry breaking \[23\]. On the other hand it is clear that the presence of CCB minima and hence of possible instabilities of the standard sphaleron along the squark configurations will be highly favoured by \( m_U^2 < 0 \) values.

The analysis of the potential \([13]\) reveals the presence of the standard electroweak minimum, at \( \tilde{q} = \tilde{u} = 0, \tilde{h}_1 = v \cos \beta, \tilde{h}_2 = v \sin \beta \), where \( v = 174.1 \text{ GeV} \) is the standard Higgs vacuum expectation value (VEV), with a depth

\[
V_{\text{ew}} = -\frac{1}{4}\frac{m_H^2 v^2}{g_3^2 \sin \beta m_H}
\]

where \( m_H \) is the lightest Higgs boson mass in the limit \( m_A \to \infty \) \[18\]. For the case we are considering of large values of \( m_A, m_H \) is just the lightest Higgs boson mass.

For values of \( \tilde{A}_t \equiv A_t + \mu/\tan \beta < \tilde{A}_t \text{ max} \), with

\[
\frac{\tilde{A}_t^2 \text{ max}}{m_Q^2} = 1 - \frac{g_3 m_H}{\sqrt{3} h_t \sin \beta m_H}
\]

there is a minimum at \( \tilde{h}_1 = \tilde{h}_2 = \tilde{q} = 0, \tilde{u}^2 = \frac{3\tilde{m}_U^2}{g_3^2} \), where \( \tilde{m}_U^2 \equiv -m_U^2 \). The condition that this CCB minimum not be lower than the standard minimum imposes on \( \tilde{m}_U \) the upper bound \( \tilde{m}_U < \tilde{m}_U^{\text{crit}} \) given by \[15\]

\[
\tilde{m}_U^{\text{crit}} = \left( \frac{m_H^2 v^2 g_3^2}{6} \right)^{1/4}
\]

We have plotted in Fig. 1 (long-dashed line) \( \tilde{m}_U^{\text{crit}} \) given by Eq. \([10]\) as a function of \( \tilde{A}_t \) for \( m_Q = m_A = 500 \text{ GeV} \) and \( \tan \beta = 3 \). We can see the curve ends at a value of \( \tilde{A}_t \) given by Eq. \([15]\).

\(^3\)Notice that we have fixed the right-handed squark fields to the trivial configuration \( D_L = D^c_R = 0 \). This result is consistent with the smallness of the Yukawa coupling \( h_b \), that we will neglect from here on, and will be justified a posteriori by the numerical results of this work.
For $\tilde{A}_t > \tilde{A}_{t,\text{max}}$ there is a CCB minimum with all fields $\tilde{h}_1, \tilde{h}_2, \tilde{q}$ and $\tilde{u}$ acquiring VEVs. The location of the minimum can only be computed numerically and $\tilde{m}_{\text{U}}^{\text{crit}}$, at which the CCB minimum becomes degenerate with the standard electroweak minimum is shown in Fig. 1 (short-dashed line) for the same values of the supersymmetric parameters.

In short, in the region below (above) the dashed line the electroweak minimum (the CCB minimum) is the global one, and only below the dashed line is electric charge and color guaranteed to be unbroken. For illustration we have depicted also in Fig. 1 the (thin solid) line where the right-handed squarks are massless. Then only below the thin solid line there exists a (local or global) standard electroweak minimum.

3. The last issue we want to address is the possible (in)stability of the standard sphalerons in the MSSM [18]. It is again clear that the region $m_U^2 < 0$, where the right-handed squark direction is unstable at the origin, favours the presence of such instabilities. One should now substitute the ansatz (12) into the Euler-Lagrange equations,

$$ (D_\nu W^{\mu\nu})^a = \frac{ig_2}{2} \sum_{Y=H_{1,2},Q_L} \left[ Y^\dagger \sigma^a (D^\mu Y) - (D^\mu Y)^\dagger \sigma^a Y \right] $$

$$ D_\mu D^\mu X = -\frac{\partial V}{\partial X^\dagger} $$

where $X = H_{1,2}, Q_L, U, D$, supplied by the appropriate boundary conditions. In some cases we expect to end up with a solution described by vanishing squark fields, whereas in other situations non-vanishing profiles can be preferred. Before embarking ourselves in the complete solution of the equations of motion, we will study the possible standard sphaleron instabilities along the squark field configurations. Hence we will treat the standard (squarkless) sphaleron as a background and we will study the corresponding squark equations in this background. This approximation will be justified a posteriori, and will provide us with an intuitive picture of the parameter space region where squarks are involved in the sphaleron solution.

The energy of the squark fields in the sphaleron background is given by:

$$ E_{\tilde{q}} = 4\pi \int dr r^2 \left[ (\partial_r \tilde{q})^2 + \frac{2}{r^2} (1 - f)^2 \tilde{q}^2 + (\partial_r \tilde{u})^2 + \frac{g_2^2}{8} \left[ \tilde{q}^4 + 2\tilde{q}^2 (\tilde{h}_1^2 - \tilde{h}_2^2) \right] + \frac{g_3^2}{6} (\tilde{q}^2 - \tilde{u}^2)^2 + h^2 \tilde{q}^2 \tilde{u}^2 + 2\mu h_t \tilde{q} \tilde{u} \tilde{h}_1 + h_t^2 \tilde{h}_2^2 \tilde{u}^2 + h_t^2 \tilde{h}_2 \tilde{q}^2 + m_Q^2 \tilde{q}^2 + m^2 \tilde{u}^2 + 2h_t A_t \tilde{q} \tilde{h}_2 \tilde{u} \right] $$

Non-vanishing squark field configurations will be associated to negative values of the previous functional. We can split it as follows:

$$ E_{\tilde{q}} = \delta^2 E_{\tilde{q}} + \delta^4 E_{\tilde{q}} $$

5
where the first piece includes quadratic terms in the squark fields and the second one collects the remaining, quartic, ones. Notice that $\delta^4 E_q \geq 0$ and then a decreasing of the energy is possible only in the case where the quadratic terms are negative. These are given by:

$$\delta^2 E_q = 4\pi \int dr \left[ r\bar{q}(r) \quad r\bar{u}(r) \right] \mathcal{M}^2 \left[ \begin{array}{l} r\bar{q}(r) \\ r\bar{u}(r) \end{array} \right]$$

where

$$\mathcal{M}^2 = \begin{bmatrix}
-\partial_r^2 + \frac{2}{r^2} (1 - f)^2 + \frac{g_2^2}{4} (\tilde{h}_1^2 - \tilde{h}_2^2) + h_l^2 \tilde{h}_2^2 + m_{\tilde{q}}^2 & h_t (\mu \tilde{h}_1 + A_l \tilde{h}_2) \\
\hfill \hfill & h_t (\mu \tilde{h}_1 + A_l \tilde{h}_2) \quad -\partial_r^2 + h_l^2 \tilde{h}_2^2 + m_{\tilde{t}}^2
\end{bmatrix}$$

To determine the instabilities of the standard sphaleron along the squark direction one must diagonalize the quadratic operator in (20). Each eigenvector of $\mathcal{M}^2$ with a negative eigenvalue

$$\mathcal{M}^2 \left[ \begin{array}{l} r\bar{q}(r) \\ r\bar{u}(r) \end{array} \right] = -\omega^2 \left[ \begin{array}{l} r\bar{q}(r) \\ r\bar{u}(r) \end{array} \right]$$

generates an instability.

We have studied the spectrum of the operator (21) numerically, using a discretized representation of the quadratic energy functional. The MSSM sphaleron is characterized by just one scale $\sim M_W$, since it is mainly controlled by the $W$ mass and the lightest Higgs boson mass, the latter never being hierarchically larger than the former. When squarks are included, a second scale, $M$, could emerge in our solution. It would be roughly given by the inverse radius of the region where squarks are non-vanishing. The value of this scale depends on the spectrum of our MSSM model, but we expect it to be in the range $M_W \lesssim M \lesssim M_{\text{SUSY}}$, where $M_{\text{SUSY}}$ is typically $\lesssim 1 \text{ TeV}$. In order to cover both scales, we will follow the approach of Ref. [10] and use a uniform discretization of the variable:

$$s = \ln \left[ \frac{1 + Mr}{1 + M_W r} \right] / \ln(M/M_W)$$

where $s$ goes from 0 to 1 when $r$ varies from 0 to infinity. We have done our analysis using 200 points and taking $M$ values within the above mentioned range. We have verified, by increasing the number of points and comparing, that the lower eigenvalue is calculated with an error smaller than 0.001 in $M_W^2$ units.

We have plotted in Fig. 1, for the same values of the supersymmetric parameters $(\tan \beta, m_A, m_Q)$ as before the (thick solid) line along which there appears a negative eigenvalue along a non-trivial squark configuration. So above (below) the thick solid line the standard sphaleron in the MSSM is unstable (stable). We have stopped the instability line when it intersects the thin solid line, where the standard electroweak minimum disappears.

This situation is different in the Standard Model, where the quartic coupling of the Higgs field is, in principle, arbitrary, so that it can be arbitrarily heavier than the $W$ boson, in which case a second scale, $M_H$, appears.
We have not shown the region with $m_{U}^{2} > 0$, but it becomes clear from the behaviour of the thick solid line in Fig. 1 that there is no instability region for $m_{U}^{2} > 0$ and any value of $\tilde{A}_{t}$.

We have also explored general values of the supersymmetric parameters in the region $1 \leq \tan \beta \lesssim 20$ (where we can comfortably neglect the bottom Yukawa coupling $h_{b}$ as compared to $h_{t}$), $100 \text{ GeV} \lesssim m_{A} \lesssim m_{Q}$ and $200 \text{ GeV} \lesssim m_{Q} \lesssim 500 \text{ GeV}$, and we have found a qualitative behaviour similar to that in Fig. 1. Even in the region of large values of $\tan \beta$, where $h_{b}$ can no longer be neglected, our results holds since the possible instability along $D_{L}, D_{R}^{c}$ would be controlled by the entry

$$-\partial_{r}^{2} + h_{b}^{2}h_{1}^{2} + m_{D}^{2}$$

in the corresponding quadratic operator, similar to that in (21), which in turns depends on the value of the bottom-quark mass. Therefore, the present experimental bounds on supersymmetric masses impose $m_{D}^{2} > 0$ which prevents, using the results obtained in this paper, the existence of instabilities along the squark fields $D_{L}, D_{R}^{c}$. Since the previous results obviously apply to slepton fields, our choice of the ansatz where slepton and down-squark fields vanish remains fully justified.

4. Using the previous results, which concern the MSSM standard electroweak sphaleron at zero temperature, we can infer the stability of the standard sphaleron against non-trivial squark configurations at temperatures higher than, or equal to, the critical temperature of the electroweak phase transition, $T_{c} \sim 100 \text{ GeV}$. The argument goes as follows: The main effect of finite temperature corrections in the MSSM sphaleron energy and profiles can be encoded, to a reasonable accuracy, in the (asymptotic) vacuum expectation values of the Higgs fields, i.e. $v_{1}(T)$ and $v_{2}(T)$ [18]. On the other hand, the only modification of the quadratic operator (21) from finite temperature corrections is to replace the squark masses $m_{Q}^{2}, m_{U}^{2}$ by effective thermal masses, $m_{Q}^{2}(T) = m_{Q}^{2} + \Pi_{L}(T)$ and $m_{U}^{2}(T) = m_{U}^{2} + \Pi_{R}(T)$, where the self-energies $\Pi_{L}(T)$ and $\Pi_{R}(T)$ can be found in Ref. [14], and whose precise expressions are not important for the sake of the present discussion. A necessary condition at $T \geq T_{c}$ is that $m_{Q}^{2}(T) > 0$, $m_{U}^{2}(T) > 0$ since otherwise the phase transition would proceed at some temperature $T_{CCB} \geq T_{c}$ towards the finite temperature CCB minimum through a strong first order phase transition [15]. Using the results of the present work we can infer there would be no instability region for the range of temperatures $T \geq T_{c}$, since negative squared masses have been found to be a necessary condition for the existence of instabilities. In other words, the stability of the standard electroweak minimum at the critical temperature implies stability of the standard electroweak sphaleron along squarks configurations. Of course, below $T_{c}$, $m_{Q}^{2}(T)$ and/or $m_{U}^{2}(T)$ can become negative and an instability region will start to grow up and evolve towards the zero temperature result shown in Fig. 1.

5. In conclusion, we have studied possible instabilities of the standard electroweak sphaleron in the MSSM along the most dangerous third generation sfermion configurations, with CCB minima, which might decrease the sphaleron energy. At zero temper-
nature, we have found that instabilities lie in the region of supersymmetric parameters where CCB minima are the global minima. Thus imposing electric charge and color conservation is a sufficient condition to ensure the stability of the standard sphaleron along third generation sfermion directions. Finally, at the electroweak phase transition temperature, imposing the stability of the electroweak minimum prevents the appearance of instabilities along sfermions configurations.

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Figure 1: For $\tan \beta = 3$ and $m_Q = m_A = 500$ GeV: a) Above the thin solid line the electroweak minimum disappears; b) Above (below) the dashed lines electric charge and color are broken (conserved); c) Above (below) the thick solid line the standard sphaleron is unstable (stable) along the squark directions.