Constraints from Inflation and Reheating on Superpartner Masses

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Abstract

Flat directions in the Minimal Supersymmetric Standard Model (MSSM) can become unstable due to radiative corrections, and the global minimum of the (zero temperature) potential can lie at large values of the squark and/or slepton fields. Here we show that, in inflationary models of early cosmology, the universe is very likely to be in the domain of attraction of this global minimum at the end of inflation. While the minimum at the origin of field space is global at sufficiently high temperatures, depending on details of the model, the universe may be trapped in the non-zero minimum until it becomes the global minimum at low temperatures. Parameter values leading to this scenario are therefore ruled out.
In an earlier paper \[1\], we pointed out that certain mass patterns for the superpartners result in a radiatively-corrected scalar potential whose global minimum is at large values of the squark and/or slepton fields. Because these minima are in general electric charge and color breaking, we concluded that these mass patterns were therefore disallowed, or at least required new physics beyond the standard model. However, Riotto and Roulet \[2\] argued that the desired minimum (with zero squark and slepton VEVs and the usual nonzero Higgs VEVs), while not global at zero temperature, is generically the global minimum at high temperature. Furthermore, if, for some reason, the universe begins in the domain of attraction of this minimum, then it will get stuck there, since the probability to tunnel to the global minimum (after the temperature drops sufficiently) is much too small. A similar argument was made in \[3\].

The assumption made (either explicitly \[2\] or implicitly \[3\]) is that the universe will begin in the global minimum of the temperature dependent scalar potential. This, however, is extremely unlikely to be the case along flat directions in field space after a period of inflation \[4\]. During inflation, fluctuations in scalar fields which parameterize flat directions, i.e. those with supersymmetry breaking squared masses $m^2 \ll H^2$ increase with time according to

$$\langle \phi^2 \rangle = \frac{1}{4\pi^2} H^3 \tau,$$  

where $H$ is the Hubble constant during the inflationary epoch. Subsequently, the long wavelength modes of these fluctuations behave as a classical background field with amplitude $\varphi_0 \simeq \langle \phi^2 \rangle^{1/2}$. More precisely, the value of $\varphi_0$ we compute in this way is the RMS width of a probability distribution (which is approximately Gaussian) for the value of $\varphi$ in a horizon volume. Solving the horizon and flatness problems requires that the duration $\tau$ of the inflationary era satisfies $H\tau \gtrsim 60$. In very general models of inflation which involve a single mass scale $m_I$, we have $H\tau \simeq (M_P/m_I)^2$, where $M_P$ is the Planck mass. The value of $m_I$ determines the size of the observed fluctuations in the microwave background, and their measured value fixes $(m_I/M_P)^2 \simeq \text{few} \times 10^{-8} \ [3]$, which then implies $H\tau \gtrsim 10^8 \ [7]$. The case of interest for us will be that of a squark or slepton field $\varphi$ (which is not the inflaton field) with a slowly varying potential $V(\varphi)$ corresponding to an almost flat direction. We define the field-dependent squared mass $m^2(\varphi) = (1/2)\partial^2 V(\varphi)/\partial \varphi^2$. The field stops evolving according to eq.\([4]\) when $m^2(\varphi)t \simeq 3H/2$. Then we find \[3\]

$$(\varphi_0^{\text{max}})^2 = \min \left\{ \frac{1}{4\pi^2} H^3 \tau, \frac{3}{8\pi^2} \frac{H^4}{m^2(\varphi)} \right\}.$$  

(2)

For $H^2 \simeq m_I^2/M_P^4 \simeq 10^{-15} M_P^2$, and for (nearly) flat directions characterized by a mass scale $m(\varphi) \simeq 10^{-16} M_P$, the first term in eq. \([4]\) is smaller than the second, and the growth of
fluctuations is therefore limited by the duration $\tau$ of the inflationary era. This generally results in a value of $\varphi_0$ which is much larger than $H$, $\varphi_0 \simeq 10^{-4} M_P$.

Thus, if $\varphi_0$ is as large or larger than the location of the global minimum of $V(\varphi)$, we are overwhelmingly likely to find ourselves, after inflation, in the domain of attraction of this minimum. This fact (that scalar fields along flat directions find themselves far from the origin in field space after inflation) has long been recognized and is the basis of the Affleck-Dine mechanism of baryogenesis [8, 9, 7], which requires large initial values of squark and slepton fields. This scenario is particularly attractive because one can in general find flat directions in field space in supersymmetric grand unified theories. That is, these theories naturally contain directions in which the scalar potential is absolutely flat $V(\varphi) = 0$ up to supersymmetry breaking effects which induce masses of order the supersymmetry breaking scale. $\varphi$ is some combination of squark and slepton and Higgs fields in which the F- and D-terms in the scalar potential vanish. If the expectation value of $\varphi$ is non-zero (and large) initially (e.g. due to inflation as described above), the subsequent evolution of certain flat direction in a supersymmetric GUT can be shown to give rise to a baryon asymmetry.

It has recently been pointed out however, that the simple picture of driving scalar fields to large vacuum values along flat directions during inflation is dramatically altered in the context of supergravity [10]. During inflation, the Universe is dominated by the vacuum energy density, $V \sim H^2 M_P^2$. The presence of a non-vanishing and positive vacuum energy density indicates that supergravity is broken and soft masses of order of $H$ are generated [11]. The implications of such terms for the charge and color breaking minima and constraints on the MSSM parameters have also been recently addressed [12]. In minimal supergravity it is quite easy to see that such mass terms are generated. In general, the scalar potential in a supergravity model is described by a Kähler potential $G$ [13] and in minimal supergravity we define the Kähler potential by

$$G = \phi_i^\dagger \phi^i + \ln |W(\phi)|^2$$

(3)

where $\phi^i$ represents all scalar fields in the theory and $W(\phi)$ is the superpotential. This results in a scalar potential of the form

$$V = e^G \left[ |W_i + \phi_i^\dagger W|^2 / |W|^2 - 3 \right]$$

(4)

where $W_i = \partial G / \partial \phi^i$. A positive vacuum energy density, $V > 0$, needed for inflation, breaks supergravity and both the exponential and the term enclosed in brackets in (4) must be non-vanishing. In fact, the exponential $e^G$ is the order parameter for supergravity breaking. Included among the $\phi^i$’s is the flat direction $\varphi$, and by inspection of (3) one finds a mass term
$e^G \varphi \varphi^*$. A large mass term of this type precludes the possibility of Affleck-Dine baryogenesis. It was noted in [10] that A-D baryogenesis could be revived in non-minimal models in which such a mass terms arises with an opposite sign having the effect of driving $\varphi$ to extremely large VEVs. In fact, it was shown in [14], that in supergravity models which possess a Heisenberg symmetry [15], including no-scale models of supergravity [16, 17], supersymmetry breaking makes no contribution to scalar masses, leaving supersymmetric flat directions flat at tree-level. One-loop corrections in general lift the flat directions, but naturally give small negative squared masses $\sim -g^2 H^2/(4\pi)^2$ for all flat directions that do not involve the stop. In the context of these theories, we therefore expect initial field values for $\varphi$ to be quite large.

In the types of models discussed above, reheating is generally quite inefficient. In models in which the inflation is coupled only gravitationally to the observable (gauged) sector, the reheating temperature is in fact very low [7], $T_R \sim \alpha^2 m_i^3/M_P^2 \sim 10^5 \text{GeV}$. However, there has been a great deal of recent activity and progress in understanding the quantitative details of the process of particle production after cosmological phase transitions and the implications of this for reheating after inflation [18, 19, 20, 21, 22], and we would like to briefly discuss the relevant aspects of these developments so as to address their implications for our work here. One of the key issues is whether or not there is symmetry restoration due to the effects of quantum fluctuations [19]. It is now clear that the occurrence of this phenomenon is model dependent [19, 21, 22]. It has been pointed out [22] that symmetry restoration effects may help localize the fields we are interested in close to the origin. However, these effects can occur in the models of interest only if specific couplings between the inflaton, an additional scalar field, and the flat-direction field $\varphi$ are present and assumed to be large [22]. In a general model, particularly one in which the inflaton has only gravitational interactions with other fields, symmetry restoration is not expected to occur.

In this paper we wish to determine the conditions under which gauge symmetries remain broken after inflation due to the presence of additional minima along flat directions. In the problem we are interested in, although for sufficiently high reheat temperatures the minimum at $\varphi = 0$ is a global minimum subsequent to the reheating of the universe after inflation, there is another minimum of the potential at $\varphi \neq 0$ which becomes the global minimum at lower temperatures and hence later in the evolution of the universe. During the epoch when the $\varphi \neq 0$ minimum is the metastable minimum of the potential, it is a relevant and appropriate question to ask whether the tunneling rate is such that we will end up at $\varphi = 0$ due to tunneling. We will examine this rate and hence the possibility of such a tunneling in this paper.
At tree level, the zero temperature potential takes the form

\[ V(\varphi, Q) = \frac{1}{2} m^2(Q) \varphi^2 + \frac{1}{n} M^{4-n} \varphi^n, \tag{5} \]

where the mass-squared \( m^2(Q) \) depends on the renormalization group scale \( Q \), and \( M \) is a new mass scale (such as the Planck mass) which characterizes the strength of non-renormalizable interactions due to physics beyond the minimal supersymmetric standard model (MSSM). Loop corrections, if included to all orders, would remove the \( Q \) dependence of \( V \). We will work with the RG-improved tree-level potential, but attempt to optimize the choice of \( Q \) so as to minimize the contributions of the loop corrections. We treat the case of flat directions, where there is no renormalizable interaction term. The lowest value of \( n \) which is allowed depends on the flat direction of interest \[23\]; for the case \( \tilde{u}_R = \tilde{s}_R = \tilde{b}_R \equiv v(Q) \) which we treated earlier, we have \( n_{\text{min}} = 10 \). For the case \( \tilde{u}_R = \tilde{c}_R = \tilde{s}_R = \tilde{e}_R \equiv v(Q) \), we have \( n_{\text{min}} = 6 \). In our numerical investigation of the tunneling rates, we concentrate on the case \( n_{\text{min}} = 6 \), as the rates will be even lower in the case \( n_{\text{min}} = 10 \). For both flat directions, the main contribution to the running of \( m^2(Q) \) comes from SU(3) interactions, which result in \[1\]

\[ m^2(Q) = m^2_{||} - \frac{2}{\pi} g_3^2(M_3) M_3^2 \ln(Q/M_3) \left\{ \frac{1 + 3g_3^2(M_3) \ln(Q/M_3)/16\pi^2}{[1 + 3g_3^2(M_3) \ln(Q/M_3)/8\pi^2]^2} \right\}, \tag{6} \]

where \( M_3 \) is the gluino mass at the scale \( Q = M_3 \), and \( m^2_{||} \) is the sum of the squark squared masses in the flat direction at a scale \( M_3 \). We will ignore the effects of the electroweak gauge couplings and the Yukawa couplings, which are sub-dominant. At large enough \( Q \), \( m^2(Q) \) becomes negative, and the minimum of the effective potential occurs at \( \varphi = v(Q) \) where

\[ v(Q) = M^{(n-4)/(n-2)}[-m^2(Q)]^{1/(n-2)}. \tag{7} \]

The best value of \( Q \) to choose (that is, the value which is likely to minimize the higher-loop contributions near the minimum) is the solution (if it exists) of \( Q = v(Q) \); if no solution exists, the apparent minimum is likely to be an invalid artifact of the one-loop approximation. Alternatively, we can just choose \( Q = \varphi \) and find the minimum of \( V(\varphi, \varphi) \). A slightly better choice is \[2\] \( Q = (\varphi^2 + M_3^2)^{1/2} \), which prevents problems at small values of \( \varphi \), and this is the scheme we adopt.

The temperature-dependent corrections to \( V(Q, \varphi) \) take the form \[24\]

\[ V_T(\varphi) = \frac{T^4}{2\pi^2} \sum_i \epsilon_i \int_0^\infty dq q^2 \ln \left[ 1 - \epsilon_i \exp \left( -[q^2 + m_i^2(\varphi)/T^2]^{1/2} \right) \right], \tag{8} \]
where $T$ is the temperature, the sum is over all particle species whose tree-level masses $m_i$ depend on $\varphi$, and $\epsilon_i = +1$ for a boson and $-1$ for a fermion. For the flat directions we consider, there are a total of 32 particles which couple strongly to the flat direction that contribute to the $\varphi$ dependence of $V_T(\varphi)$: 8 gluons with mass $g_3\varphi$, 8 squarks with mass $(m_0^2 + g_3^2\varphi^2)^{1/2}$, 8 gluinos with mass $\frac{1}{2}[(M_3^2 + 2g_3^2\varphi^2)^{1/2} + M_3]$, and 8 quarks with mass $\frac{1}{2}[(M_3^2 + 2g_3^2\varphi^2)^{1/2} - M_3]$. (Actually, the fermion mass eigenstates are quark/gluino mixtures.) The effects of $V_T(\varphi)$ are significant for $g_3\varphi \lesssim T$, where it can be approximated as

$$V_T(\varphi) \simeq -\frac{2}{3}\pi^2 T^4 + 2g_3^2 T^2 \varphi^2.$$  

There are also particles which couple to the flat directions weakly and via Yukawa interactions, and these particles acquire much smaller masses when $\varphi$ is large than those which couple strongly to $\varphi$. Thus when $\varphi \gtrsim T/g_3$, the number density of the strongly coupled fields becomes exponentially suppressed, but the weakly coupled fields can still contribute to $V_T$, and it is therefore these weakly coupled degrees of freedom which are numerically important in determining the temperature at which these large scale minima disappear. These light degrees of freedom include bino/quark, hypercharge gauge boson/squark, Higgsino/quark and Higgs/squark admixtures, and for sufficiently high temperatures, they contribute an amount

$$\Delta V_T(\varphi) = \left(\frac{1}{4}c_6^2 g_1^2 + \frac{g_2^2}{4M_W^2 \sin^2 \beta} m_+^2 + \frac{g_2^2}{4M_W^2 \cos^2 \beta} m_-^2\right) T^2 \varphi^2$$

(10)

to the effective potential, where $c_6(c_{10}) = \sqrt{20}/9(0)$ and $\{m_+, m_-\} = \{m_c, m_s\}(\{m_u, m_b\})$ for $n_{\text{min}} = 6(10)$. In the $n = 10$ case, $c_{10} = 0$ because there remains an unbroken U(1) gauge symmetry which is a linear combination of hypercharge and a color generator. For flat directions corresponding to $n_{\text{min}} = 6$, the second minimum of $V_T$ (when it exists) occurs at $v \sim 10^{11}$ GeV, for $\sim$ TeV gluino masses. For $n_{\text{min}} = 10$, $v \sim 10^{15}$ GeV.

To calculate the tunneling rate from the $\varphi \neq 0$ minimum to the origin, we need to solve the equation

$$\left(\partial_r^2 + \frac{2}{r} \partial_r\right) \varphi = \frac{\partial}{\partial \varphi} [V(Q, \varphi) + V_T(\varphi)]_{Q=(\varphi+M_Z^2)^{1/2}}$$

(11)

subject to the boundary conditions $\partial_r \varphi|_{r=0} = 0$ and $\varphi(\infty) = v$, where $v$ is the location of the minimum of $V(Q, \varphi) + V_T(\varphi)$. The tunneling rate is then proportional to $\exp(-S_3/T)$, where the three-dimensional action is given by

$$S_3 = 4\pi \int_0^\infty dr \ r^2 \left[\frac{1}{2}(\partial_r \varphi)^2 + V(Q, \varphi) + V_T(\varphi) - V_0\right],$$

(12)

\footnote{We thank Alessandro Strumia for pointing this out to us.}
where $V_0$ is a constant which cancels off the potential energy at $\varphi = v$. $S_3$ can be computed numerically, but to gain intuition we note that at small $\varphi$ the potential is dominated by $V_T(\varphi)$; the desired solution is then

$$\varphi(r) = A \frac{\sinh(2g_3 T r)}{r}$$

(13)

where $A$ is a constant. Near the minimum, the potential can be approximated as

$$V(\varphi) \simeq \frac{1}{2} m^2 (\varphi - v)^2$$

(14)

where $m^2 \sim |m^2(v)|$ is positive. Near $\varphi = v$ the desired solution is

$$\varphi = v - B \frac{\exp(-mr)}{r}$$

(15)

where $B$ is another constant. If we now make the drastic approximation that the full potential is given by eq.(9) for $\varphi < \varphi_c$ and by eq.(14) for $\varphi > \varphi_c$, with $\varphi_c$ given by the value of $\varphi$ where the two approximate potentials equal each other, then it is possible to solve for $A$, $B$, and $S_3$ analytically. For $T \lesssim v$ and $m \ll T^2/v$, we find $S_3/T \sim v^3/T^3$. Thus we must have $T$ near $v$ in order to have a fast enough tunneling rate.

On the other hand, for $T$ near $v$, the field-dependent part of $V_T$ dominates the zero-temperature potential at $\varphi = v$, and for $T$ larger than some critical temperature $T_{crit}(m_{||}, M_3)$, the second minimum (corresponding to non-zero $v$) along the flat direction is removed. In fig. 1, we plot contours of $T_{crit}(m_{||}, M_3)$, in units of $10^7$ GeV, for the $n_{\text{min}} = 6$ flat direction $\tilde{u}_R = \tilde{c}_R = \tilde{s}_R = \tilde{b}_R \equiv v(Q)$. We have taken $\tan \beta = 2$. The $x$-axis specifies the gluino mass $M_3(M_3)$, while the $y$-axis labels $\hat{m}_{\tilde{q}}(\hat{m}_{\tilde{q}})$, where $\hat{m}_{\tilde{q}}^2 = (m_{||}^2 - m_{\tilde{e}_R}^2)/3$ is the average squark mass in the flat direction for $m_{\tilde{e}_R}^2 \ll m_{\tilde{q}}$. Above the solid curve, which lies along the line $\hat{m}_{\tilde{q}} \approx 0.76 M_3$ for large $M_3$, the zero temperature potential $V(\varphi, Q)$ (and hence the full potential $V(\varphi, Q) + V_T$) has no second minimum ($T_{crit}$ drops off very rapidly to zero above the top contour). We see that if the reheating temperature $T_R$ is $\lesssim 10^7$ GeV, the finite temperature effective potential after reheating will contain large scale minima for the same parameter values as for the zero-temperature effective potential. For larger $n_{\text{min}}$, $v$ is larger, along with the correspondingly larger critical temperatures. In fig. 2, we plot contours of $T_{crit}(m_{||}, M_3)$, in units of $10^9$ GeV, for the $n_{\text{min}} = 10$ flat direction $\tilde{u}_R = \tilde{s}_R = \tilde{b}_R \equiv v(Q)$, and we see that for this direction we require larger reheating temperatures, on the order of a few $\times 10^9$ GeV, in order to remove the $\varphi \neq 0$ minima. In this case the average squark mass $\hat{m}_{\tilde{q}}$ is defined by $\hat{m}_{\tilde{q}}^2 = m_{||}^2/3$.

After inflation, the field $\varphi$ parameterizing the flat direction rolls from $\varphi_0$, given by (2), toward the the nearest minimum of the potential. Because the time-scale for the field
to reach this minimum is much shorter than the time for decay of the inflaton and the subsequent thermalization of its decay products (assuming, as above, that the inflaton is gravitationally coupled to ordinary matter), the evolution of $\varphi$ is determined by the zero-temperature effective potential. The initial conditions at reheating are therefore determined by the value of $\varphi_0$ relative to the maximum of $V$ separating the $\varphi = 0$ and $\varphi \neq 0$ minima. If it is to the left of the barrier, $\varphi$ will roll to the origin and remain there and we recover the scenario described by [2]. If $\varphi_0$ is to the right of the barrier, it will roll to the $\varphi \neq 0$ minimum (unless it is so far to the right that it picks up a sufficient velocity so as to carry it over the barrier and settle once again at the origin). When non-renormalizable interactions for flat directions are included, the maximum value for $\varphi_0$ is limited by the effective mass at large $\varphi$. For our $n = 6$ case, the minimum is at roughly $\varphi \sim 10^{11}$ GeV, and the position of the barrier depends on $m_\parallel$ and $M_3$, but is typically much smaller than $v$, while $\varphi_0^{\text{max}} \lesssim 10^{14}$ GeV. Thus we expect to be to the right of the barrier, and we have determined that so long as $\varphi_0 < 1.5 \times 10^{14}$ GeV, the field will settle in the broken vacuum. For $n = 10$, the minimum is shifted up to about $\varphi \sim 10^{15}$ GeV and is comparable to the maximum value of $\varphi_0$, while again the maximum of the barrier is a much lower scale. Thus if the zero-temperature potential for $\varphi$ has a large-scale minimum, this is where we expect the field to sit at the beginning of reheating. If $T_R < T_{\text{crit}}(m_\parallel, M_3)$, the position of the minimum of $V_{\text{eff}}$ is unaffected by the temperature corrections, and after reheating $\varphi$ will be trapped in the symmetry-breaking minimum.

It is possible that thermal effects can excite $\varphi$ over the potential barrier and allow nucleation of bubbles of the symmetric vacuum. In order to determine whether $\varphi$ remains trapped in the $\varphi \neq 0$ minimum, we numerically compute the finite temperature transition rate from the $\varphi \neq 0$ minimum to the true minimum (of the finite-temperature potential) at $\varphi = 0$. For fixed temperature $T$, we use the full finite-temperature corrections (8) to the effective potential to compute the bounce solution to (11) and the resulting three-dimensional action (12). The tunneling probability per volume per unit time is [25]

$$\Gamma \sim T^4 e^{-S_3/T}.$$  

The fraction of space remaining in the broken phase is $f = e^{-P}$ [28], where

$$P \sim \frac{M_{\text{Pl}}^4}{T_0^3} \int_{T_d}^{T_R} T^{-2} \left(1 - \frac{T_0}{T}\right)^3 e^{-S_3/T} dT < \frac{M_{\text{Pl}}^4}{T_d T_0^3} e^{-S_3/T_R},$$

and where $T_0$ is the current temperature of the universe and $T_d$ is the temperature at which the two minima become degenerate. For $n_{\text{min}} = 6$ flat directions, $T_d$ is of order $10^6 - 10^7$ GeV,
while for $n_{\text{min}} = 10$ flat directions $T_d$ is of order $10^8 - 10^9$ GeV; for reheat temperatures less than this, there is of course no tunneling. The expression on the right-hand side of (17) can be a gross overestimate, depending on how fast $S_3$ drops off as the temperature falls. As $P$ is exponentially sensitive to $S_3$, we can take $P = 1$ as a critical value, such that $P < 1$ implies no nucleation of the symmetric phase. We then find that $\varphi$ remains trapped in the large-scale minimum as long as

$$\frac{S_3}{T_R} > 217 + \ln \left( \frac{10^9 \text{GeV}}{T_R} \right) + \ln \left( \frac{T_R}{T_d} \right)$$

(18)

For fixed values of $(m_{||}, M_3)$, we have looked for values of $T_R$ satisfying $T_R < T_{\text{crit}}(m_{||}, M_3)$, but violating (18). We find that $S_3$ rises very quickly as $T_R$ drops below $T_{\text{crit}}$, and we were unable to resolve a single case where (18) does not hold by better than an order of magnitude. That is, $T_R$ must be extremely close to $T_{\text{crit}}$ in order for (18) to be violated (in which case the integral in (17) would need to be evaluated numerically anyway, yielding a much lower value for $P$ than given by the right-hand-side of (17)). (In [2], tunneling from $\varphi = 0$ to $\varphi = v$ was found to occur over non-negligible regions of the parameter space at small $m_{||}$, where the width of the barrier was small ($< 1$ TeV); in this case, the barrier widths are much larger).

We can therefore say that if the reheat temperature $T_R$ is less than the critical temperature $T_{\text{crit}}$ (shown in Figure (1)), we expect the field $\varphi$ parameterizing the flat direction to become trapped in the large-scale symmetry-breaking minimum, and to remain there until the large-scale minimum becomes the global minimum at low temperatures. This effectively rules out regions of $(m_{||}, M_3)$ parameter space where such minima exist, or at the very least implies new physics beyond the standard model.

From fig. 2, we see that typical values of $T_{\text{crit}}$ for the $n_{\text{min}} = 10$ flat direction are about a few times $10^9$ GeV. We can ask whether reheat temperatures of this magnitude are reasonable. There are of course strong limits on the reheat temperature in supersymmetric theories due to gravitino production. If the gravitino mass is related to the supersymmetry breaking scale (though this need not be the case in no-scale supergravity [27]) and $m_{3/2} \sim 1$ TeV, then the production of gravitinos during the reheat process and their subsequent decay will lead to the photodestruction of deuterium and $^4$He. This allows one [28] to set a limit on the reheat temperature which is comparable to $T_{\text{crit}}$, i.e. $T_R \lesssim \text{a few } \times 10^9$ GeV, though this limit can be evaded for a significantly larger gravitino mass. We note that the gravitino decay bound was re-examined in the light of parametric resonance by Allahverdi and Campbell [29]. They concluded that these bounds cannot be evaded even in the presence of parametric resonance.

In summary, we find that after an inflationary epoch, fields parameterizing flat directions
in the SUSY scalar potential will settle into large-scale ($> 10^{10}$ GeV) minima, if they exist, and will remain there through reheating until the minima become global minima at low temperatures. This effectively confirms the qualitative conclusions of \[1\], albeit with some new model-dependence. In particular, we conclude that the SUSY particle mass relationships corresponding to the region under the solid line in fig. 1 are inconsistent with inflationary models, such as those in no-scale supergravity, in which 1) the inflaton is gravitationally coupled to the observable sector, or in any model in which the reheat temperature is less than that given by the contours in fig. 1 or fig. 2, depending on which flat direction is populated during inflation, and 2) there are no positive $O(H^2)$ contributions to the scalar masses during inflation. We note that under these conditions, parametric resonance will not localize the fields we are interested in close to the origin. The case of the $n_{\text{min}} = 6$ flat directions we have considered is the most conservative, in the sense that the regions in the parameter space which yield large scale charge and color breaking minima can be made acceptable by lower reheat temperatures than in any of the other $n_{\text{min}} > 6$ flat directions involving squarks. The $n_{\text{min}} = 10$ flat directions require much higher reheating temperatures. In general, however, different flat directions are not necessarily mutually flat, and a large VEV in an $n_{\text{min}} = 6$ direction may preclude large-scale minima in other directions. Conservatism then requires that we restrict our conclusions to the $n_{\text{min}} = 6$ case, although there is no reason that inflation should populate $n_{\text{min}} = 6$ directions in preference to directions with larger VEVs.

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Figure Captions

Fig. 1) Contours of constant $T_{\text{crit}}$, in units of $10^7$ GeV, as a function of the gluino mass $M_3$ and the average squark mass $\hat{m}_q = \sqrt{(m_{||}^2 - m_{\tilde{e}_R}^2)/3}$, for the $n_{\text{min}} = 6$ flat direction $\tilde{u}_R = \tilde{c}_R = \tilde{s}_R = \tilde{e}_R \equiv v(Q)$. Above the top solid curve, the zero-temperature potential (and hence the full potential $V(\varphi, Q) + V_T$) has no second minimum.

Fig. 2) Contours of constant $T_{\text{crit}}$, in units of $10^9$ GeV, as a function of the gluino mass $M_3$ and the average squark mass $\hat{m}_q = \sqrt{m_{||}^2/3}$, for the $n_{\text{min}} = 10$ flat direction $\tilde{u}_R = \tilde{s}_R = \tilde{b}_R \equiv v(Q)$. Above the top solid curve, the zero-temperature potential (and hence the full potential $V(\varphi, Q) + V_T$) has no second minimum.
Figure 1
Figure 2