Dynamical Solution to Supersymmetric CP Problem with Vanishing B Parameter

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Abstract

The CP violation gives rise to severe restriction of soft breaking terms in supersymmetric standard models. Among them, constraints on the holomorphic soft mass of Higgs doublets (the B parameter) are difficult to satisfy due to the other inherent problem in the Higgs potential; the µ problem. In this letter, it is argued that these CP and µ problems can be rather relaxed provided that B is vanishing at high-energy scale. A generic mechanism and some examples of model are presented to dynamically realize this condition by introducing gauge singlet fields.
1. Introduction and Idea

Supersymmetry is one of the most attractive extensions of the standard model (SM). It has a variety of interesting properties in phenomenological and theoretical viewpoints, which can bring novel approaches to unresolved problems of the SM. For example, grand unification with supersymmetry predicts low-energy values of gauge couplings precisely consistent with the experimental data [1]. In the others, the weak/Planck mass hierarchy is stabilized against radiative corrections due to the non-renormalizable nature of superpotential terms [2].

However, supersymmetry is not a symmetry experimentally observed at low-energy scale and must be broken by soft breaking terms, which do not reintroduce quadratic divergences and not spoil mass hierarchies [3]. These soft supersymmetry breaking terms consist of gaugino masses, scalar trilinear couplings, and scalar masses. Thus the number of physical couplings is dramatically increased compared to an exactly supersymmetric case or that of the SM. Moreover, generic forms of soft breaking terms tend to lead phenomenological disasters such as too large rate of flavor-changing rare processes like $\mu \rightarrow e\gamma$ as well as large CP violation, which is the subject of this letter. To avoid these problems and make models viable, supersymmetry-breaking dynamics is required to have some special properties, and various such mechanisms have been proposed in the literature. Physical mass spectrum can be explicitly calculated for a fixed mechanism of supersymmetry breaking and mediation to the SM sector.

In general, CP-violation measurements provide one of the severest constraints for supersymmetry-breaking couplings. An important point is that CP violation occurs even in the absence of flavor violation. Working with a rather strong hypothesis of flavor universality for soft terms, there still remain two types of CP-violating phases which cannot be rotated away by field redefinition. One is the phase of scalar trilinear couplings ($A$ terms) relative to that of gaugino masses, and another is that of holomorphic bilinear terms of scalars ($B$ terms). These phases (in the basis where gaugino masses are real) are required to be nearly real by the experimental results such as non-observation of sizable electric dipole moments of the neutron and leptons [4]. To satisfy the constraints seems to need some non-trivial, involved mechanisms or unnatural fine-tuning of model parameters. This is the CP problem in supersymmetric extensions of the SM [5].

As for $A$ terms, it is found that phase constraints may be somewhat weak. This is partly because $A$ terms contribute to the electric dipole moments of leptons, which highly restrict models, only through the neutralino diagrams with small parameters ($U(1)$ gauge coupling and a ratio of two Higgs vacuum expectation values (VEVs)).
Furthermore initial complex phases are reduced by renormalization-group running down to low energy. From a viewpoint of model building, however, supersymmetry-breaking $A$ terms are correlated with Yukawa couplings. It is a non-trivial task to reproduce the experimentally observed values of fermion mass hierarchy and CP violation in the Kaon system while there is no new source of CP violation besides that of Yukawa couplings.

Notice that for high-scale supersymmetry breaking like gravity mediation, some kind of implements should be usually introduced to suppress direct couplings between supersymmetry-breaking and visible sectors. The ‘separation’ of the two sectors is inevitable, for example, to avoid dangerous flavor-changing operators radiatively generated in Kähler potential. It is found that such suppression mechanisms also reduce $A$ terms and then their CP-violating phases. In this letter, we assume this kind of separation for simplicity. On the other hand, low-energy supersymmetry breaking scenarios like gauge mediation also separate the two sectors by introducing messenger fields, and predict vanishing $A$ terms at leading order. Accordingly, in both these cases, low-energy nonzero values of $A$ parameters are generated via renormalization by gaugino masses. Hence relative phases of $A$ in the basis where gaugino masses are real turn out to be zero, which is suitable for obtaining CP-conserving results in supersymmetric SM.

However the situation is very different for $B$ terms, and naive separation mechanisms do not work for the CP problem, unlike $A$ terms. In supersymmetric SM, $B$ appears in the holomorphic soft mass of the Higgs doublets. It is well-known that the most delicate issue in dealing with the Higgs $B$ parameter is that it closely relates to other problems in the Higgs potential, which involves the supersymmetric mass parameter $\mu$. The supersymmetric SM Higgs sector has to satisfy the following three conditions for realizing proper phenomenology;

\begin{align}
\mu & \ll \Lambda, \\
B\mu & \lesssim \mu^2, \\
\arg(M^* B) & \simeq 0,
\end{align}

where $\Lambda$ is a cutoff scale of supersymmetric SM (e.g. the GUT or Planck scale), and $M$ is a universal gaugino mass. We here take a convention where the $\mu$ parameter is real. The first condition is nothing but the well-known gauge hierarchy problem (the $\mu$ problem); it is unnatural to have a tree-level $\mu$ parameter of the electroweak scale which is much smaller than the natural scale of theory $\Lambda$. A vanishing $\mu$ may be natural by symmetry argument but has been excluded by the LEP experiment. Supersymmetry
itself can provide a support to stabilize a tree-level value against quantum correction but does not give any reasons why $\mu$ is finite and so small relative to larger fundamental scales. The second condition (2) is also necessary for the correct electroweak symmetry breaking. If this condition is not satisfied, the potential is unbounded from below along some supersymmetric flat direction. It should be noticed that the condition (2) implies a potential difficulty that the supersymmetry-conserving parameter $\mu$ has to be correlated with the supersymmetry-violating one $B$.

The third condition (3) is concerned with CP violation as stated above. It says that nonzero phases of a gaugino mass and the $B$ parameter must be aligned. One could see whether this condition is satisfied or not once one knows about all structures of supersymmetry breaking dynamics. That is, since gaugino masses appear through a gauge kinetic function and on the other hand, the $B$ parameter resides in the Higgs potential, they apparently have no connection to each other. Nevertheless Eq. (3) imposes a Fermi-Bose relation even in supersymmetry-breaking sector. Since there has been few viable mechanisms to arbitrarily control phase values, phase alignment of the terms which originate from different sources is rather unnatural and needs some adjustment of model parameters.

In this way, each condition holds in itself independent problems to solve. Therefore it seems an annoying issue to construct realistic models satisfying all the requirements on the Higgs potential. In practice, various attempts have been done to have successful mechanisms.

For a tree-level $\mu$ term, the required large hierarchy (1) is the foremost problem to be considered. A weak-scale $\mu$ could arise if some symmetries forbid a bare coupling and small symmetry-breaking effects generate an effective $\mu$ term. In this case, generally, a nonzero $B$ parameter is also effectively induced. However its phase is less under control without any additional assumptions and may conflict with the constraint (3) for the CP problem. As we will see below, $B$ often contains a term proportional to the gravitino mass $m_{3/2}$. The magnitude of the gravitino mass is fixed by the requirement that the cosmological constant should vanish. Its phase is, however, not constrained and is in general different from that of the gaugino mass. For example, in the string-inspired supergravity model, the $B$ parameter is given by $B = A - m_{3/2}$ in the dilaton dominant case [7]. In this case, the phase of the $A$ parameter is the same as that of the gaugino mass and given by that of the dilation $F$ term which is generally undetermined. Then the phase of $B$ is uncorrelated to the gaugino mass phase. Other scenarios along this line have also been discussed in [8].

Non-minimal form of Higgs Kähler potential is another origin of the Higgs mass couplings [9]. Nonzero $\mu$ and $B$ parameters are generated via supersymmetry-breaking
effects and then cannot be so hierarchical, which is suitable for the requirement (2). This mechanism, however, may seem to need to introduce tuning of parameters or involved structures of models in order to suppress the CP phase of the $B$ parameter.

In theories where soft mass terms are generated by loop diagrams as in gauge mediation or anomaly mediation \[10, 11\], another severer problem often emerges. This is due to the fact that $\mu$ and $B\mu$ are usually generated at the same loop order. Thus $B = B\mu/\mu$ is given by the supersymmetry breaking scale without loop suppression which is much larger than other mass parameters of super-particles including the Higgs doublets. This conclusion seems quite generic and hard to resolve, and naive attempts such as a tree-level $\mu$ term or non-minimal Kähler potentials discussed above do not work. Some mechanisms to resolve the difficulty were discussed in \[12\].

In this letter, we would like to stress that there is a simple assumption that rather relaxes these tight constraints on the Higgs mass parameters. That is, at some high-energy scale, the $B$ parameter is vanishing:

$$B = 0.$$  \[(4)\]

One can immediately see that this assumption resolves the problem (3). A low-energy (electroweak) nonzero $B$ term is generated via renormalization group evolution with the boundary condition (4). The running of $B$ is driven by gaugino masses and possibly the large scalar top coupling. The latter is controlled also by the gaugino masses as long as $A$ terms are supposed to be small at high-energy scale. This situation is achieved, for example, with the mechanisms explained before. The gaugino masses thus govern the low-energy $B$ parameter including its phase, and hence the condition (3) turns out to be satisfied. Interestingly, this solution for the CP problem is independent of actual phase values, which is difficult to control. Vanishing $B$ is natural in a sense that some symmetries such as Peccei-Quinn symmetry and possibly $R$ symmetry are restored in this limit.

The remaining problems of Higgs potential, (1) and (2), are simultaneously settled if a nonzero $\mu$ is generated by supersymmetry-breaking parameters (and can be real) \[8\]. Suppose that there is no $\mu$ term in the vacuum without supersymmetry breaking, and breaking effects shift the vacuum resulting in an effective $\mu$ term. The $\mu$ problem (1) is thus solved in a technically natural way. In addition, $B$ is now generated by renormalization-group evolution and becomes of the order of gaugino masses, and the condition (2) is achieved. Interestingly, the solution can be discussed only within the Higgs potential sector and needs not to require knowledge of other parts of theory such

\*This condition was considered in gauge mediation \[13\] as well as in gaugino mediation \[14\].
as gaugino masses. It is a reasonable situation that $B$ is vanishing since mechanisms to control phase values have not been established.

In this letter, we present a mechanism where the solution with the boundary condition (4) is dynamically realized. We will discuss supersymmetry-breaking effects at high-energy scale, but the mechanism presented here is general and can be applied to other supersymmetry-breaking models. Phenomenological implications of this solution to the supersymmetric CP problem will be discussed elsewhere [15].

2. Models

As mentioned before, we suppose that a separation of supersymmetry-breaking and visible sectors occurs. If this is not the case, induced higher-dimensional Kähler terms which link two sectors generate phenomenologically dangerous operators for flavor-changing neutral currents, etc. Such an implement is indeed necessary for models to be viable. A separation may be accomplished, for example, in a geometrical way in higher-dimensional theories [10] or by strong coupling dynamics of superconformal field theories [16]. The latter might be interpreted as a gravity dual of the former. In this letter, we assume such mechanisms for separation, for simplicity, and concentrate on dynamics in the visible sector. With this locality at hand, in particular, scalar trilinear supersymmetry-breaking terms vanish at leading order of perturbation theory. However scalar bilinear terms do not necessarily share the same result. This is because a non-trivial $\mu$-generation mechanism has to be fixed as discussed in the introduction.

Let us consider models in which SM gauge singlet fields generate the Higgs bilinear couplings. Supersymmetry-breaking effect is expressed in terms of the gravitational multiplet. The relevant part of Lagrangian is

$$L = \int d^4 \theta \left[ \phi^\dagger \phi f(S, S^\dagger) + H^\dagger H + \bar{H}^\dagger \bar{H} \right] + \int d^2 \theta \phi \lambda g(S) H \bar{H} + \int d^2 \theta \phi^3 W(S, \cdots),$$

(5)

where $S$ is the singlet which couples to the Higgs doublets through a function $g(S)$ in the superpotential with a coupling constant $\lambda$. Contribution of possible additional (singlet) fields are denoted by $\cdots$. A chiral superfield $\phi$ is the compensator multiplet with Weyl weight 1. In the case we consider, a VEV of $\phi$ is the only source of supersymmetry breaking in the visible sector and is determined by hidden sector dynamics by requiring a vanishing cosmological constant,

$$\phi = 1 + F_\phi \theta^2.$$  

(6)

In the Lagrangian (5), we have rescaled the Higgs fields into a canonical form. It is found convenient to work in this basis, since the physical masses of Higgs fields are
clearly understood and the effects of $\phi$ to singlet fields is easier to evaluate. With appropriate form of potentials, nonzero VEVs of singlet fields are generated after supersymmetry breaking and can provide a solution to the $\mu$ problem.

The equations of motion for the singlet auxiliary components are

$$F_\phi \partial_i f + \sum_j F_j \partial_i \partial_j f + \partial_i W^* = 0,$$

(7)

where $\partial_i (\partial_i)$ denotes the derivative with respect to a field $\Phi_i = S, \cdots (\Phi^*_i = S^*, \cdots)$. Here and in the following, we consider the vacua around $H = \bar{H} = 0$. In this case, any details of Higgs fields such as Kähler form are irrelevant to the potential analyses. By integrating out the auxiliary components and minimizing the scalar potential of singlet fields, one finds the VEVs of scalar components satisfy the following equations:

$$0 = \frac{\partial V}{\partial \Phi_i} = -2F_\phi \partial_i W + \sum_j \partial_i F^*_j \sum_k F_k \partial_k \partial_j f,$$

(8)

where $F_i$’s have been replaced with the scalar components through the equations (7). The factor 2 in the right-handed side of Eq. (8) is related to the fact that superpotential terms have Weyl weight 3. The Higgs mass parameters $\mu$ and $B$ are given by

$$\mu = \lambda g, \quad -B\mu = \lambda (F_\phi g + F_S \partial S g).$$

(9)

Here two comments are in order. First it is found from Eq. (9) that a tree-level $\mu$ term, namely, $g = \text{const}$ leads to $B = -F_\phi$. This is too a large value for electroweak symmetry breaking to be turned on as long as soft terms are induced at loop level. Secondly, we do not include a non-minimal Higgs Kähler term, $\int d^4 \theta \frac{g^2}{\phi} HH$, which could give another origin of the Higgs mass couplings [9]. As mentioned in the introduction, it does not give a complete solution to the Higgs mass problems without introducing fine-tuning of couplings or involved model structures; $\mu \propto F_\phi^2$ and $B = F_\phi$, and generally, the phases of $B\mu$ and gauginos are different. In the following, we simply assume the minimal form of Kähler potential, $\partial_i \partial_j f = \delta_{ij}$ as a good first-order approximation. Note, however, that even when higher-order Kähler terms suppressed by powers of the fundamental scale are taken into account, they are irrelevant in the analyses unless models contain VEVs very close to the fundamental scale.

Our mechanism for a vanishing $B$ term introduces two singlet fields, here we call $S$ and $N$. Let us consider the following form of superpotential terms:

$$\int d^2 \theta \phi \lambda g(S) H \bar{H} + \int d^2 \theta \phi^3 \left[\lambda' N g^2(S) + W(S, \cdots)\right],$$

(10)

where $g(S)$ and $W(S)$ are arbitrary functions of $S$, and $\lambda$ and $\lambda'$ are coupling constants. From the equations of motion (7) and (8), one can see that the above superpotential
leads to \( B = 0 \) at the potential minimum with respect to the \( N \) field. In fact, Eq. (8) for \( \Phi_i = N \) now reads

\[
0 = \frac{\partial V}{\partial N} = -2\lambda' g(F_\phi + F_S \partial_S g).
\]

Then the \( B \) parameter (9) vanishes as long as the VEV of \( g \) is nonzero, that is, a nonzero \( \mu \) parameter is generated. Actual values of the couplings are irrelevant to the result.

A key ingredient is that the superpotential has at most linear dependence of \( N \). With the above suitable form of the \( N^1 \) term, the two equations, the minimization by \( N \) (i.e., Eq. (8) with the minimal Kähler) and the \( B \) parameter in (9), take the same form besides nonzero overall factors. In this case, a separation that suppresses direct couplings between the hidden and visible sectors plays an important role. For example, if we have a Kähler term

\[
\int d^4 \theta Z(X, X^\dagger) N^1 N,
\]

where \( X \) is the field responsible to supersymmetry breaking in the hidden sector, it induces a tree-level soft mass of \( N \) and may destabilize our vacuum of \( B = 0 \). Separation mechanisms we assume throughout this letter is crucial to derive our solution as well as for suppressing other flavor problems.

Thus the desired boundary condition \( B = 0 \) is dynamically realized with the superpotential (10). Notice that \( W \) does not play any role in obtaining a vanishing \( B \) term. It is clear from the above derivation that the only requirement which \( W \) must satisfy is that it does not depend on the \( N \) field. Therefore additional dynamics can be incorporated into \( W \) (and also the Kähler of \( S \)) to have preferable VEVs of singlet fields while the result \( B = 0 \) is still preserved. Furthermore, as for a polynomial \( g(S) \), it might be curious to have the \( Ng^2(S) \) term in the Lagrangian (11). However as stressed above, \( B = 0 \) is guaranteed as long as the \( W \) part does not contain the \( N \) field. Accordingly it is rather easy to have \( Ng^2(S) \) with polynomial \( g \) if we start with the Lagrangian

\[
\int d^2 \theta \phi \lambda T H \bar{H} + \int d^2 \theta \phi^3 [\lambda' NT^2 + \kappa U(T - g(S)) + W(S, \cdots)],
\]  

\footnote{Even if we worked with a non-minimal, involved Kähler potential \( f \), the same result can be obtained by additional conditions for \( f \): \( \partial_N \partial_S f = \partial_N^2 f = 0 \).}

\footnote{We assume that there is no tadpole of \( N \), whose existence may disturb our vacuum. Such a tadpole could be prevented by (\( R \)) symmetries, which requires some modification of the model. Symmetries would also restrict the superpotential in desired forms. For example, an \( R \) symmetry under which \( N \) has a charge +2 leads to a linear dependence of \( N \) in the superpotential.}
where $T$ and $U$ are the gauge singlets. With this action, $B = 0$ is also achieved. On the other hand, integrating out $T$ and $U$, we have just the Lagrangian \[1\] with polynomial $g$. It could be easier for the above form of action to follow from symmetries.

With the condition $B = 0$ at hand, a low-energy $B$ term is generated by gaugino mass effects in renormalization group evolution. The phase of $B$ is therefore automatically aligned with that of gaugino masses, and suppression of CP violation is accomplished. Another part of our solution to the Higgs mass problems is that the $\mu$ term is induced by supersymmetry breaking dynamics. In the present case, this corresponds to the effects of the compensator field $F_\phi$ in the visible sector. The VEV of the scalar component of $S$, which gives rise to the $\mu$ parameter, is fixed by minimizing potential with respect to $S$. The $\mu$ problem \[2\] can be solved with an appropriate form of the functions $g(S)$ and $W(S)$. Deviations from the minimal Kähler form of $S$ could also stabilize the $S$ field. It should be noted that detailed forms of $g$ and $W$ do not affect the above argument for the vanishing $B$ parameter.

Let us study several examples of $g$ and $W$, and estimate orders of magnitude of effectively induced $\mu$ parameter. Our aim here is not to present complete theories but to give simple toy models in order to study structures of the mechanism. At first, as a special case, consider $g = S$ and $W = S^3$, that is, all the superpotential terms are cubic and renormalizable ones. In this case, however, the VEV of $S$ is undetermined by the supersymmetry breaking $F_\phi$. This is because the compensator $\phi$ can be absorbed by the rescaling $S_\phi \to S, \cdots$, and hence does not give supersymmetry-breaking potential at tree level. Therefore dimensionful parameters and/or higher-dimensional operators have to be introduced in the Lagrangian.

[ Model 1 ]

First we consider the model

$$
\begin{align*}
g &= S, \\
W &= m_S S^2,
\end{align*}
$$

where $y$ is the coupling constant of $O(1)$. A tree-level $\mu$ term may be forbidden by imposing a discrete symmetry $S \to -S$ and $\bar{H} \to -\bar{H}$. This symmetry also forbids a generation of tadpole operator for the $S$ field as well as the dimension 5 operators for nucleon decay. In the supersymmetric limit $F_\phi = 0$, $S$ is forced to be zero. Then at the minimum of the potential including supersymmetry breaking, the VEV of $S$ flows to a nonzero value $\langle S \rangle \simeq \frac{1}{\lambda'} (m_S F_\phi)^{1/2}$. The VEV of $N$ is determined by the equation of motion for $N$ and is given by $\langle N \rangle \simeq \frac{1}{\lambda'} m_S$. We note that around this vacuum, radiatively induced soft masses are negligibly small and do not disturb the potential
analysis here. Integrating out the high-energy dynamics, the effective \( \mu \) parameter is generated as
\[
\mu \simeq \frac{\lambda}{\lambda'} (m_S F_\phi)^{1/2},
\]
which is a geometric mean of \( m_S \) and the gravitino mass \( F_\phi \).

The scale of \( \mu \) depends on both \( m_S \) and \( F_\phi \). In the case of high-energy supersymmetry breaking, i.e. the heavy gravitino, \( m_S \) must be chosen as a TeV scale for a small \( \mu \) parameter. A natural way to achieve this is to slightly modify the Kähler form of \( S \) [9]. If Kähler potential contains the quadratic term of \( S \),
\[
\int d^4 \theta \phi^\dagger \phi S^2 + \text{h.c.},
\]
induces a mass \( m_S \simeq F_\phi^* \). (This deformation of Kähler potential also induces a holomorphic soft mass of \( S \) but does not affect the result (15).) Thus \( \mu \) is roughly on the correct order of magnitude; \( \mu \simeq F_\phi \). When soft terms are loop induced as in anomaly mediation, some tuning of couplings is still required. Other mechanisms discussed in the literature [8] can also be utilized to obtain a supersymmetry-breaking order of mass \( m_S \).

For low-energy supersymmetry breaking, a gravitino mass is much smaller than the weak scale, \( F_\phi \sim F_{\text{hid}}/M_{\text{Pl}} \). Here \( \sqrt{F_{\text{hid}}} \) is a supersymmetry-breaking scale in hidden sector which is separated by taking \( \sqrt{F_{\text{hid}}} \) as much lower than the fundamental scale. As a consequence, \( A \) terms are reduced. From Eq. (15), in this case, a correct order of \( \mu \) parameter is realized if \( m_S \) is slightly lower than the Planck scale. For example, with \( \sqrt{F_{\text{hid}}} \sim 100 \) TeV, \( m_S \) is around the GUT scale \( \sim 10^{16} \) GeV for a weak-scale \( \mu \) parameter. This example gives a solution to the \( \mu \) problem in gauge mediation scenarios. The well-known trouble of too a large \( B \) parameter is avoided in the present case by realizing the condition \( B = 0 \) dynamically. The \( \mu \) term is generated at tree level but the \( B \) term at loop level.

[ Model 2 ]

Next let us discuss an example with higher-dimensional operators (superpotential terms whose mass dimensions are greater than three). The superpotential is given by
\[
g = S^n, \quad W = yS^{n+2},
\]
where the exponent \( n \) is not necessarily an integer but satisfies \( n > 1 \) unless the coupling \( \lambda' \) is too small. (The \( n = 1 \) case corresponds to the conformal limit discussed above where \( S \) is not determined by supersymmetry-breaking effects.) The \( S \) direction is stabilized by the potential also in this case. Minimizing the potential, the \( \mu \) parameter
\[\text{Alternatively one can take } \lambda \text{ small enough to make } \mu \text{ in the correct order of magnitude. This does not, however, provide a natural explanation of the } \mu \text{ parameter.}\]
in the vacuum is found to be

$$\mu \approx \frac{1 - n}{2n(n + 2) \lambda} \lambda F_\phi. \quad (17)$$

We find that \( \mu \) is proportional to \( F_\phi \) irrespective of \( n \), and therefore this model gives a solution to the \( \mu \) problem relevant to high-energy supersymmetry breaking like gravity mediation and other scenarios where the gravitino mass is around the electroweak scale \[17\]. This example shares essential features with the Model 1. For example, a tree-level \( \mu \) can be forbidden with a discrete symmetry, \( S^n \rightarrow -S^n \) and \( \bar{H} \rightarrow -\bar{H} \) (that is broken by higher-dimensional operators).

In case of \( n \gg 2 \), the singlet \( N \) tends to have a rather flat potential and a large value of VEV beyond the Planck scale, though depending on the couplings. A natural way to cure this problem is to add a small soft mass for \( N \), which stabilizes the VEV \( \langle N \rangle \) to an intermediate scale value irrespectively of \( n \). The small soft mass, roughly around a MeV to GeV scale, can be induced by a small deviation from the minimal Kähler form (the exact separation) without conflicting with flavor-changing neutral current constraints. It is also found that introducing soft mass of \( N \) makes the model cosmologically safe, with small decay rate and a negligible contribution to energy density by coherent oscillation of the scalar. We note that the problem could also be removed by adopting other forms of potential \( W \) including relevant fields.

3. Summary

We have shown that the condition \( B = 0 \) for a solution to the supersymmetric CP problem can be dynamically realized in the Higgs potential with gauge singlet fields. With the appropriate terms in superpotential, minimizing scalar potential with respect to singlet fields exactly leads to a vanishing \( B \) parameter. At the same time, the supersymmetric mass \( \mu \) is generated by supersymmetry-breaking effects. A nonzero \( B \) parameter is induced during renormalization-group evolution down to low energy. It is driven by gaugino masses and hence the CP-violating phase of \( B \) parameter vanishes in the basis where gaugino masses are real. This solution to the CP problem does not need any information about gaugino masses and can be discussed only in the Higgs sector. Our solution can be incorporated in various mediation mechanisms, including gauge mediation, gaugino mediation, etc. Vanishing \( A \) and \( B \) terms is shown to be an attractive way to solve the CP and \( \mu \) problems and be attainable with natural physical implications.
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