SUSY Breaking by stable non-BPS configurations

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Abstract

A simple mechanism for SUSY breaking is proposed due to the coexistence
of BPS domain walls. It requires no messenger fields nor complicated SUSY
breaking sector on any of the walls. We assumed that our world is on a
BPS domain wall and that the other BPS wall breaks the SUSY preserved by
our wall. We obtain an \(\mathcal{N} = 1\) model in four dimensions which admits an
exact solution of a stable non-BPS configuration of two walls. The stability
is assured by a topological quantum number associated with the winding on
the field space of the topology of \(S^1\). We propose that the overlap of the
wave functions of the Nambu-Goldstone fermion and those of physical fields
provides a practical method to evaluate SUSY breaking mass splitting on our
wall thanks to a low-energy theorem. This is based on our recent works hep-
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1 Introduction

Supersymmetry (SUSY) provides the most realistic models to solve the hierarchy problem in unified theories [1]. One of the most important issues in model building of SUSY unified theories has been for some years how to understand the SUSY breaking in our observable world. Many models of SUSY breaking uses some kind of mediation of the SUSY breaking from the hidden sector to our observable sector.

Recently the “Brane World” scenario has become quite popular where our four-dimensional spacetime is realized on the wall embedded in a higher dimensional spacetime [2, 3]. In order to discuss the stability of such a wall, it is often useful to consider SUSY theories as the fundamental theory. Moreover, SUSY theories in higher dimensions are a natural possibility in string theories. These SUSY theories in higher dimensions have 8 or more supercharges, which should be broken partially if we want to have a phenomenologically viable SUSY unified model in four dimensions. Such a partial breaking of SUSY is nicely obtained by topological defects [4]. Domain walls or other topological defects preserving part of the original SUSY in the fundamental theory are called the BPS states in SUSY theories. Walls have co-dimension one and typically preserve half of the original SUSY, which are called 1/2 BPS states [5, 6, 7]. Junctions of walls have co-dimension two and typically preserve a quarter of the original SUSY [8, 9].

The new possibility offered by the brane world scenario stimulated studies of SUSY breaking. Recently we have proposed a simple mechanism of SUSY breaking due to the coexistence of different kinds of BPS domain walls and proposed an efficient method to evaluate the SUSY breaking parameters such as the boson-fermion mass-splitting by means of overlap of wave functions involving the Nambu-Goldstone (NG) fermion [10]–[12]. We have exemplified these points by taking a toy model in four dimensions, which allows an exact solution of coexisting walls with a three-dimensional effective theory [10]. Although the first model is only meta-stable, we were able to show approximate evaluation of the overlap allows us to determine the mass-splitting reliably. More recently, we have constructed a stable non-BPS configuration of two walls in an \( \mathcal{N} = 1 \) supersymmetric model in four dimensions to demonstrate our idea of SUSY breaking due to the coexistence of BPS walls. We have also extended our analysis to more realistic case of four-dimensional effective theories and examined the consequences of our mechanism in detail [11].

Our proposal for a SUSY breaking mechanism requires no messenger fields, nor complicated SUSY breaking sector on any of the walls. We assume that our world is on a wall and SUSY is broken only by the coexistence of another wall with some distance from our wall. The NG fermion is localized on the distant wall and its overlap with the wave functions of physical fields
on our wall gives the boson-fermion mass-splitting of physical fields on our wall thanks to a low-energy theorem \[13\].

The purpose of this paper is to illustrate our idea of SUSY breaking due to the coexistence of BPS walls by taking a simple soluble model with a stable non-BPS configuration of two walls and to extend our analysis to more realistic case of four-dimensional effective theories. We work out how various soft SUSY breaking terms can arise in our framework. Phenomenological implications are briefly discussed. We also find that effective SUSY breaking scale observed on our wall becomes exponentially small as the distance between two walls grows. The NG fermion is localized on the distant wall and its overlap with the wave functions of physical fields on our wall gives the boson-fermion mass-splitting of physical fields on our wall thanks to a low-energy theorem. We have proposed that this overlap provides a practical method to evaluate the mass-splitting in models with SUSY breaking due to the coexisting walls.

## 2 BPS equation and topological quantum number

Let us illustrate the BPS equation and topological quantum number for \(\frac{1}{2}\)-BPS state in terms of a simple model with a single chiral scalar field \(\Phi' = (A', \psi', F')\) and a superpotential \(W = \Phi' - \frac{1}{3}\Phi'^3\). After eliminating the auxiliary field \(F\) the bosonic part of the Lagrangian becomes

\[
L = -\partial_\mu A'^* \partial^\mu A' - |1 - A'^2|.
\] (2.1)

We have absorbed possible constants into the normalization of field and coordinates for simplicity. The model has two SUSY vacua at \(A' = \pm 1\). The supertransformation of the fermion \(\psi'\) is given by

\[
\delta \psi' = i\sqrt{2}\sigma^\mu \bar{\epsilon} \partial_\mu A' + \sqrt{2}\epsilon F'.
\] (2.2)

If we choose \(A'\) to depend only on one coordinate, say, \(x^2 = y\), and choose \(\epsilon = i\sigma^2 \bar{\epsilon}\), the half of supersymmetry is conserved by the configuration satisfying the BPS equation

\[
\frac{dA'}{dy} = 1 - A'^2.
\] (2.3)

It admits a wall solution connecting the SUSY vacuum \(-1\) at \(y = -\infty\) to another SUSY vacuum \(+1\) at \(y = \infty\)

\[
A^{(1)}_{cl}(y) = \tanh(y - y_1),
\] (2.4)

where \(y_1\) denotes the position of the wall. Orthogonal linear combination of supercharges are conserved if the anti-BPS equation is satisfied

\[
\frac{dA'}{dy} = -(1 - A'^2),
\] (2.5)
which admits a wall solution connecting the SUSY vacuum \( +1 \) at \( y = -\infty \) to another SUSY vacuum \( -1 \) at \( y = \infty \)

\[
A^{(2)}_{cl}(y) = - \tanh(y - y_2),
\]

(2.6)

where \( y_2 \) denotes the position of the wall. If we combine these two solutions, we obtain a wall anti-wall configuration. In fact we have found exact solution of the equation of motion which is a non-BPS state and gives an example of the SUSY breaking due to the coexistence of BPS and anti-BPS walls [10]. The wall anti-wall configuration is unstable due to the annihilation into vacuum. It is desirable to have a model with stable but non-BPS two wall configuration.

We have found a way to give the topological quantum number. We shall give a topology of \( S^1 \) to field space so that we can have a notion of winding from a compactified base space which is also \( S^1 \). To achieve that goal, we change field variable \( A' \) into a periodic variable \( A \)

\[
A' = \sin A, \quad \Phi' = \sin \Phi.
\]

(2.7)

Then the SUSY vacua occurs at \( A = \pi \left( n + \frac{1}{2} \right) \) with the periodicity \( A = A + 2\pi \). The BPS equation (2.3) becomes

\[
\frac{dA}{dy} = \cos A.
\]

(2.8)

The BPS solution (2.4) is mapped into a solution of this transformed BPS Eq.(2.8)

\[
\sin A_{cl}^{(1)}(y) = \tanh(y - y_1)
\]

(2.9)

connecting the SUSY vacuum \( A = -\pi/2 \) at \( y = -\infty \) to \( A = \pi/2 \) at \( y = \infty \). The solution of the anti-BPS equation connecting the SUSY vacuum \( A = \pi/2 \) at \( y = -\infty \) to \( A = 3\pi/2 \) at \( y = \infty \).

\[
\sin A_{cl}^{(2)}(y) = - \tanh(y - y_2).
\]

(2.10)

Now we can see these solutions can be smoothly connected in the field space since the field \( A = \pi/2 \) at right end point of the BPS wall is the same as the field at the left end point of the anti-BPS wall. This suggests that we may have a non-BPS solution of two wall configuration which is non-BPS. Indeed we found that a simple model with minimal kinetic term provides the BPS equation (2.3) and that there is an exact solution for the non-BPS configuration of two walls which winds around the field space \( A \) once [11].

\[1\] Since the 1/2-BPS solution is intrinsically real, we can interprete the nonlinearity of the right-hand side of the BPS equation in two ways: Either as the derivative of a superpotential (our model), or as the inverse of the nontrivial Kähler metric in a nonlinear sigma model as in [14]. Here we choose a simpler possibility.
3 Stable non-BPS configuration of two walls

In order to illustrate our basic ideas, we consider three dimensional domain walls in four-dimensional spacetime. Our model reads

$$L = \bar{\Phi}\Phi|_{\theta^2} + W(\Phi)|_{\bar{\theta}^2} + h.c., \quad W(\Phi) = \frac{\Lambda^3}{g^2}\sin\left(\frac{g}{\Lambda}\Phi\right). \quad (3.1)$$

We have introduced a scale parameter $\Lambda$ with a mass-dimension one and a dimensionless coupling constant $g$, both of which are real positive. Choosing $y = X^2$ as the extra dimension and compactify it on $S^1$ of radius $R$. Other coordinates are denoted as $x^m$ ($m = 0, 1, 3$), i.e., $X^\mu = (x^m, y)$. The bosonic part of the model is

$$L_{bosonic} = -\partial^\mu A^* \partial_\mu A - \frac{\Lambda^4}{g^2}\left|\cos\left(\frac{g}{\Lambda}A\right)\right|^2. \quad (3.2)$$

The target space of the scalar field $A$ has a topology of a cylinder. This model has two vacua at $A = \pm \pi\Lambda/(2g)$, both lie on the real axis.

In the limit $R \to \infty$, we have a BPS domain wall solution

$$\sin\frac{g}{\Lambda}A_{cl}^{(1)}(y) = \tanh(\Lambda(y - y_1)), \quad (3.3)$$

which interpolates the vacuum at $A = -\pi\Lambda/(2g)$ to that at $A = \pi\Lambda/(2g)$ as $y$ increases from $y = -\infty$ to $y = \infty$ and conserves the real two component SUSY charge $Q_{\alpha}^{(1)}$ which can be regarded as supercharges in three dimensions. We have also an anti-BPS wall solution

$$\sin\frac{g}{\Lambda}A_{cl}^{(2)}(y) = -\tanh(\Lambda(y - y_2)), \quad (3.4)$$

which interpolates the vacuum at $A = \pi\Lambda/(2g)$ to that at $A = 3\pi\Lambda/(2g) = -\pi\Lambda/(2g)$ and preserves another real two component supercharge $Q_{\alpha}^{(2)}$. Here $y_1$ and $y_2$ are integration constants and represent the location of the walls along the extra dimension. The four-dimensional supercharge $Q_\alpha$ is a sum of these two supercharges $Q_\alpha = \frac{1}{\sqrt{2}}(Q_{\alpha}^{(1)} + iQ_{\alpha}^{(2)})$. Each wall breaks a half of the bulk supersymmetry and all of the bulk supersymmetry will be broken if these walls coexist.

For this model, we have found an exact solution of the non-BPS two wall configuration which is stable due to the winding number: $\pi(S^1) = Z$. Such a configuration should be a solution of the equation of motion,

$$\partial^\mu \partial_\mu A + \frac{\Lambda^3}{g}\sin\left(\frac{g}{\Lambda}A^*\right)\cos\left(\frac{g}{\Lambda}A\right) = 0. \quad (3.5)$$

A general real static solution of Eq.(3.3) that depends only on $y$ is found to be

$$A_{cl}(y) = \frac{\Lambda}{g}\text{am}\left(\frac{\Lambda}{k}(y - y_0), k\right), \quad (3.6)$$
Figure 1: The profile of the classical solution $A_{cl}(y)$. The dotted lines $A = -\pi \Lambda/(2g)$ and $A = 3\pi \Lambda/(2g)$ are identified.

where $k$ and $y_0$ are real parameters and the function $am(u, k)$ denotes the amplitude function, which is defined as an inverse function of $u(\varphi) = \int_0^\varphi \frac{d\theta}{\sqrt{1-k^2 \sin^2 \theta}}$. If $k < 1$, the solution $A_{cl}(y)$ is a monotonically increasing function with

$$A_{cl} \left( y + \frac{4kK(k)}{\Lambda} \right) = A_{cl}(y) + 2\pi \frac{\Lambda}{g}. \quad (3.7)$$

This is the solution that we want. Since the field $A$ is an angular variable $A = A + 2\pi \Lambda/g$, we can choose the compactified radius $2\pi R = 4kK(k)/\Lambda$ so that the classical field configuration $A_{cl}(y)$ contains two walls and becomes periodic modulo $2\pi \Lambda/g$. We shall take $y_0 = 0$ to locate one of the walls at $y = 0$. Then we find that the other wall is located at the anti-podal point $y = \pi R$ of the compactified circle. We have computed the energy of a superposition of the first wall $A_{cl}^{(1)}(y)$ located at $y = y_1$ in Eq. (3.3) and the second wall $A_{cl}^{(2)}(y)$ located at $y = y_2$ in Eq. (3.4). This energy can be regarded as a potential between two walls in the adiabatic approximation and has a peak at $|y_1 - y_2| = 0$ implying that two walls experience a repulsion. This is in contrast to a BPS configuration of two walls which should exert no force between them. Thus we can explain that the second wall is settled at the anti-podal point $y = \pi R$ in our stable non-BPS configuration because of the repulsive force between two walls. Since the repulsive force forces the other wall to oscillate around the anti-podal point when a small fluctuation is added, we have a physical reason to obtain a stable spectrum without any tachyon.
Figure 2: The mode functions for the bosonic modes $a_{R,0}$ and $a_{R,1}$. The solid line represents the profile of $b_{R,0}(y)$ and the dashed line is that of $b_{R,1}(y)$.

In the limit of $R \to \infty$, i.e., $k \to 1$, $A_{cl}(y)$ approaches near $y = 0$ to the BPS configuration $A_{cl}^{(1)}(y)$ with $y_1 = 0$ which preserves $Q^{(1)}$, and near $y = \pi R$ to $A_{cl}^{(2)}(y)$ with $y_2 = \pi R$ which preserves $Q^{(2)}$. The profile of the classical solution $A_{cl}(y)$ is shown in Fig.1. We will refer to the wall at $y = 0$ as “our wall” and the wall at $y = \pi R$ as “the other wall”.

4 Mode expansion and effective Lagrangian

The fluctuation fields around the background $A_{cl}(y)$ can be expanded into modes

$$A(X) = A_{cl}(y) + \frac{1}{\sqrt{2}}(A_{R}(X) + iA_{I}(X)), \quad \Psi_{\alpha}(X) = \frac{1}{\sqrt{2}}(\Psi_{\alpha}^{(1)}(X) + i\Psi_{\alpha}^{(2)}(X)).$$ (4.1)

The four-dimensional fluctuation fields can be expanded as

$$A_{R}(X) = \sum_{p} b_{R,p}(y)a_{R,p}(x), \quad A_{I}(X) = \sum_{p} b_{I,p}(y)a_{I,p}(x),$$

$$\Psi^{(1)}(X) = \sum_{p} f_{p}^{(1)}(y)\psi_{p}^{(1)}(x), \quad \Psi^{(2)}(X) = \sum_{p} f_{p}^{(2)}(y)\psi_{p}^{(2)}(x).$$ (4.2) (4.3)

Exact mode functions and mass-eigenvalues can be found for several light modes of $b_{R,p}(y)$,

$$b_{R,0}(y) = C_{R,0}dn\left(\frac{\Lambda y}{k}, k\right), \quad m_{R,0}^2 = 0,$$
Figure 3: The mode functions for fermionic zero-modes $\psi_0^{(1)}$ and $\psi_0^{(2)}$. The solid line represents the profile of $f_0^{(1)}(y)$ and the dashed line is that of $f_0^{(2)}(y)$.

$$
\begin{align*}
\frac{b_{R,1}(y)}{R} &= C_{R,1} \text{cn} \left( \frac{\Lambda y}{k}, k \right), \quad m_{R,1}^2 = \frac{1 - k^2}{k^2} \Lambda^2, \\
\frac{b_{R,2}(y)}{R} &= C_{R,2} \text{sn} \left( \frac{\Lambda y}{k}, k \right), \quad m_{R,2}^2 = \frac{\Lambda^2}{k^2},
\end{align*}
$$

(4.4)

where functions $\text{dn}(u, k)$, $\text{cn}(u, k)$, $\text{sn}(u, k)$ are the Jacobi’s elliptic functions and $C_{R,p}$ are normalization factors. For $b_{l,p}(y)$, we can find all the eigenmodes

$$
\frac{b_{l,p}(y)}{R} = \sqrt{\frac{2}{\pi R}} e^{i \frac{\Lambda y}{k}}, \quad m_{l,p}^2 = \Lambda^2 + \frac{p^2}{R^2}, \quad (p \in \mathbb{Z}).
$$

(4.5)

The massless field $a_{R,0}(x)$ is the Nambu-Goldstone (NG) boson for the breaking of the translational invariance in the extra dimension. The first massive field $a_{R,1}(x)$ corresponds to the oscillation of the background wall around the anti-podal equilibrium point and hence becomes massless in the limit of $R \to \infty$. All the other bosonic fields remain massive in that limit.

For fermions, only zero modes are known explicitly,

$$
\begin{align*}
f_0^{(1)}(y) &= C_0 \left\{ \text{dn} \left( \frac{\Lambda y}{k}, k \right) + k \text{cn} \left( \frac{\Lambda y}{k}, k \right) \right\}, \quad f_0^{(2)}(y) = C_0 \left\{ \text{dn} \left( \frac{\Lambda y}{k}, k \right) - k \text{cn} \left( \frac{\Lambda y}{k}, k \right) \right\},
\end{align*}
$$

(4.6)

where $C_0$ is a normalization factor. These fermionic zero modes are the NG fermions for the breaking of $Q^{(1)}$-SUSY and $Q^{(2)}$-SUSY, respectively.

Thus there are four fields which are massless or become massless in the limit of $R \to \infty$: $a_{R,0}(x)$, $a_{R,1}(x)$, $\psi_0^{(1)}(x)$ and $\psi_0^{(2)}(x)$. The profiles of their mode functions are shown in Fig.2 and Fig.3. Other fields are heavier and have masses of the order of $\Lambda$. We will concentrate ourselves on the breaking of the $Q^{(1)}$-SUSY, which is approximately preserved by our wall at $y = 0$. So we call the field $\psi_0^{(2)}(x)$ the NG fermion.
We can obtain a three-dimensional effective Lagrangian by substituting the mode-expanded fields into the Lagrangian (3.1), and carrying out an integration over $y$

\[ L^{(3)} = -V_0 - \frac{1}{2} \partial^m a_{R,0} \partial_m a_{R,0} - \frac{1}{2} \partial^m a_{R,1} \partial_m a_{R,1} - \frac{i}{2} \psi_0^{(1)} \phi_0^{(1)} - \frac{i}{2} \psi_0^{(2)} \phi_0^{(2)} - \frac{1}{2} m_{R,1}^2 a_{R,1}^2 + g_{\text{eff}} a_{R,1} \psi_0^{(1)} \psi_0^{(2)} + \cdots, \]

where $\phi \equiv \gamma^m_{(3)} \partial_m$ and an abbreviation denotes terms involving heavier fields and higher-dimensional terms. Here $\gamma$-matrices in three dimensions are defined by $\left( \gamma^m_{(3)} \right) \equiv \left( \sigma^2, i\sigma^3, -i\sigma^1 \right)$ and $V_0$ and $g_{\text{eff}}$ are the vacuum energy and the effective Yukawa coupling

\[ g_{\text{eff}} \equiv \frac{g}{\sqrt{2}} \int_{-\pi R}^{\pi R} dy \cos \left( \frac{g}{\Lambda} A_{cl}(y) \right) b_{R,1}(y) f_0^{(1)}(y) f_0^{(2)}(y) = \frac{g}{\sqrt{2}} \frac{C_0^2}{C_{R,1}} (1 - k^2). \]

The nonvanishing mass term for $a_{R,1}$ shows a mass splitting associated to the SUSY breaking due to the coexistence of BPS and anti-BPS walls. The amount of this mass term can be related to the Yukawa coupling $g_{\text{eff}}$ by means of low energy theorem. This fact provides a powerful method to evaluate the mass splitting between superpartners by evaluating the overlap of mode functions with the NG fermion [10]–[12].

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