Domain Wall Junction in $\mathcal{N} = 2$ Supersymmetric QED in four dimensions

Kazuya Kakimoto * and Norisuke Sakai †

Department of Physics, Tokyo Institute of Technology
Tokyo 152-8551, JAPAN

Abstract

An exact solution of domain wall junction is obtained in $\mathcal{N} = 2$ supersymmetric (SUSY) QED with three massive hypermultiplets. The junction preserves two out of eight SUSY. Both a (magnetic) Fayet-Iliopoulos (FI) term and complex masses for hypermultiplets are needed to obtain the junction solution. There are zero modes corresponding to spontaneously broken translation, SUSY, and $U(1)$. All broken and unbroken SUSY charges are explicitly worked out in the Wess-Zumino gauge in $\mathcal{N} = 1$ superfields as well as in components. The relation to models in five dimensions is also clarified.

*e-mail address: kakimoto@th.phys.titech.ac.jp
†e-mail address: nsakai@th.phys.titech.ac.jp
1 Introduction

In recent years, models with extra dimensions have attracted much attention [1], [2]. In this brane-world scenario, our world is assumed to be realized on extended topological defects such as domain walls or junctions. On the other hand, supersymmetry (SUSY) provides the most promising idea to build realistic unified theories beyond the standard model [3]. Brane-world scenario in the supersymmetric theories can provide an opportunity for a realistic model building on walls and/or junctions. Moreover, it can offer a possible explanation of SUSY breaking [4]–[9], in particular by means of the coexistence of walls [10], [11]. SUSY has been useful to obtain solutions of walls and junctions as BPS states, which preserve a part of SUSY [12].

Domain walls can conserve half of the SUSY, and are called $\frac{1}{2}$ BPS states. They have been extensively studied in globally supersymmetric theories [13], [14]. More recently, an exact BPS wall solution in supergravity theories has been constructed in four dimensions [15] and in five dimensions [16]. We need to consider topological defects such as junctions of walls, to consider a fundamental theory in space-time dimensions higher than five. The domain wall junctions have been studied [17]–[24] and can preserve a quarter of original SUSY. An exact analytic solution of the junction has been obtained in the $\mathcal{N} = 1$ SUSY field theories in four dimensions [20]. Possibility of junction solution has also been explored in supergravity [23]. The exact solution has been useful to unravel several unexpected properties of domain wall junctions. The new Nambu-Goldstone fermion modes associated with the junction is found to be non-normalizable [21]. The new central charge associated with the junction was once considered to be a mass of the junction. However the exact solution showed that the central charge contributes negatively to the energy of the junction [20], [21]. Therefore it should more properly be interpreted as a binding energy of the walls which meet at the junction. As another topological defect with co-dimension two, an exact solution of vortices on $S^2$ has also been obtained before [25].

The SUSY theories in dimensions higher than four are required to have at least eight supercharges. Theories with eight SUSY are often called $\mathcal{N} = 2$ SUSY theories even in five or six dimensions, since they have twice as many SUSY charges compared to the simple SUSY theories in four dimensions. BPS wall solutions have been constructed in the $\mathcal{N} = 2$ SUSY nonlinear sigma models [26]–[30]. Lump and Q-lump solutions preserving $1/8$ and $1/4$ SUSY, respectively, have also been considered [31], [32]. On the other hand, the BPS wall junction has been constructed in linear [20], [21] and nonlinear sigma models [24] only in $\mathcal{N} = 1$ SUSY models in four dimensions.

The first analytic solution of the BPS junction has been obtained for an $\mathcal{N} = 1$ $U(1) \times U(1)$ gauge theory with six charged and one neutral chiral scalar fields with minimal kinetic terms [20], which was constructed as a toy model for the $\mathcal{N} = 2$ $SU(2)$ gauge theory with one flavor [33]. Subsequently it was realized that one can get rid of the vector multiplet by identifying six charged chiral scalar fields pair-wise into three chiral scalar fields. One still obtains the same junction solution as a BPS solution [21] in this model with three “charged” and one “neutral” chiral scalar fields with minimal kinetic terms (linear sigma model), without gauge field at all (Wess-Zumino model). It has also been shown that one can obtain the same solution in an
$\mathcal{N} = 1$ nonlinear sigma model with only single “neutral” chiral scalar field, by eliminating the other three “charged” chiral scalar fields appropriately [24]. In all these solutions, one finds that the “neutral” chiral scalar field plays a central role in constructing the junction solution. On the other hand, a neutral scalar field is contained in the $\mathcal{N} = 2$ vector multiplet in the case of the $\mathcal{N} = 2$ SUSY QED. Therefore it is tempting to embed the $\mathcal{N} = 1$ gauge theory and its junction solution into the $\mathcal{N} = 2$ SUSY QED.

The purpose of this paper is to give an exact analytic solution for the BPS domain wall junction in an $\mathcal{N} = 2$ SUSY QED with three massive hypermultiplets. This is the first example of an exact junction solution in $\mathcal{N} = 2$ SUSY theories. By explicitly working out eight SUSY transformations, we show that the junction solution preserves two out of eight SUSY, namely it is a $\frac{1}{4}$ BPS state. Although the solution have many similarities with the previously obtained $\frac{1}{4}$ BPS junction solution in $\mathcal{N} = 1$ SUSY theory, the resulting spectrum of the low-energy effective theory is richer. For instance, we observe that there are zero modes corresponding to spontaneously broken $U(1)$ global symmetries [29], [30]. Similarly to our previous solution in $\mathcal{N} = 1$ theory [20], the Nambu-Goldstone modes on the junction background are not normalizable. As pointed out in Ref. [21], it may be possible to obtain a normalizable wave function when it is embedded into supergravity as explored in Ref. [23]. We also show that the same eight SUSY transformations can be derived from a nontrivial dimensional reduction of the $\mathcal{N} = 2$ SUSY QED in five dimensions.

The $\mathcal{N} = 2$ SUSY theories with vector and hypermultiplets were introduced by Fayet using an automorphism of SUSY algebra [34]. He used both $\mathcal{N} = 1$ superfield and component formalisms. The $\mathcal{N} = 1$ superfield formalism makes only four SUSY manifest, but has been useful also to write down massless nonlinear sigma models [35]. Harmonic superspace formalism can make all eight SUSY manifest [36]–[39] and has been used to formulate $\frac{1}{2}$ BPS equations to obtain BPS walls [27], [40]. Even in the harmonic superspace formalism, however, it has been useful to use the Wess-Zumino gauge to clarify the physical field content of the theory [27]. The Wess-Zumino gauge in the component formalism allows us to construct all the eight SUSY transformations explicitly. We also find that the action in terms of component fields can be assembled into $\mathcal{N} = 1$ superfield formalisms making four out of eight SUSY manifest in two ways. Namely we can rewrite the same action in terms of two different superfields. One of them makes a set of four SUSY manifest, and the other makes the set of remaining four SUSY manifest. Of course we cannot make eight SUSY manifest in any one of the $\mathcal{N} = 1$ superfield formalisms. We shall here employ the $\mathcal{N} = 1$ superfield formalism [5], [41]–[44] as well as the component formalism both in the Wess-Zumino gauge.

We find it essential to allow complex mass parameters in order to obtain a junction solution. The $\mathcal{N} = 2$ SUSY theories are often derivable by means of a dimensional reduction from five and/or six dimensions [15]. In this spirit, we also show that these $\mathcal{N} = 2$ SUSY transformations can be understood in terms of a massive $\mathcal{N} = 2$ theory in five dimensions. Since the massive theory in five dimensions can be obtained by a nontrivial dimensional reduction à la Scherk and Schwarz [46] in one spacial direction, the mass parameter should be real. Therefore we find that it is difficult to extend our junction solution in the eight SUSY theory to a junction solution of $\mathcal{N} = 2$ SUSY theory in five or six dimensions within the context of our multi-flavor QED. If we make a nontrivial dimensional reduction for two spacial directions from six dimensions, we can
obtain complex mass parameters. Therefore \( \mathcal{N} = 2 \) SUSY theories in four dimensions can have complex mass parameters which allow the junction solutions.

In sect. 2 our model of \( \mathcal{N} = 2 \) SUSY massive multiflavor QED is introduced and BPS equations are derived as a minimum energy condition, and are shown to conserve one out of four SUSY in the \( \mathcal{N} = 1 \) superfield formalism. In sect. 3 an exact junction solution is obtained as a solution of \( \frac{1}{4} \) BPS equations of \( \mathcal{N} = 1 \) superfield formalism. Zero modes are also briefly analyzed. In sect. 4 the remaining SUSY transformations are found by means of an automorphism of SUSY algebra. We also show that our model is invariant under the eight SUSY transformations and our BPS junction solution preserves two out of eight SUSY. Sect. 5 is devoted to relate the eight SUSY transformations in four-dimensions from \( \mathcal{N} = 2 \) SUSY transformations in five dimensions.

\section{\( \mathcal{N} = 2 \) SUSY QED and BPS equations}

As one of the simplest models with eight SUSY, we consider \( \mathcal{N} = 2 \) SUSY model with local \( U(1) \) gauge symmetry in four dimensions with the gauge coupling constant \( g \). If an \( \mathcal{N} = 1 \) SUSY vector superfield \( V_+ \) is combined with an \( \mathcal{N} = 1 \) SUSY chiral scalar superfield \( \Phi_+ \), an \( \mathcal{N} = 2 \) SUSY vector multiplet is obtained. In order to distinguish the four SUSY from the remaining four SUSY which will appear later, we denote the \( \mathcal{N} = 1 \) superfield here by a suffix +. Combining \( \mathcal{N} = 1 \) SUSY chiral scalar superfields \( Q_{a+} \) with \( U(1) \) charge +1, and \( \tilde{Q}_{a+} \) with \( U(1) \) charge −1 gives an \( \mathcal{N} = 2 \) hypermultiplet. The suffix \( a = 1, \ldots, n \) denotes flavor. The \( \mathcal{N} = 2 \) SUSY allows us to introduce the mass \( m_a \) of the hypermultiplet for each flavor. Since our gauge symmetry is \( U(1) \), the electric \( c \in \mathbb{R} \) and the magnetic \( b \in \mathbb{C} \) FI parameters can also be introduced without violating the \( \mathcal{N} = 2 \) SUSY [34]. Assuming a minimal kinetic term for the \( \mathcal{N} = 2 \) vector and hypermultiplets, we thus obtain the \( \mathcal{N} = 2 \) SUSY massive multiflavor QED. Using \( \mathcal{N} = 1 \) superfield formalism, the Lagrangian is given by\(^1\)

\[
\mathcal{L} = \frac{1}{4g^2} \left( W^a W^a_{\bar{q}} \right) + \frac{1}{2g^2} \left( \Phi_+ \right) + \sum_{a=1}^{n} \left( Q^a_e \right) \left( Q^a + \tilde{Q}^a_{\bar{e}} e^{2V_+} \right) + h.c.
\]

\[
-2cV_+ \left( \sum_{a=1}^{n} \left( \phi^a - m_a \right) Q^a + \tilde{Q}^a_{\bar{a}} e^{-2V_+} \right) + b \Phi_+ \left( \phi^a - m_a \right) + h.c.
\]

(2.1)

where the \( \mathcal{N} = 1 \) vector multiplet \( V_+ \) and the chiral scalar multiplet \( \Phi_+ \) are multiplied by the gauge coupling \( g \) to make the \( \mathcal{N} = 2 \) SUSY more easily visible. The coupling of \( \Phi_+ \) with the hypermultiplets \( Q_{a+}, \tilde{Q}_{a+} \) in the last line of Eq. (2.1) is dictated by the requirement of the \( \mathcal{N} = 2 \) SUSY. If the mass parameters are absent \( m_a = 0 \), the Lagrangian is invariant under the following global \( U(n) \) transformations:

\[
Q_{a+} \rightarrow Q'_{a+} = Q_{a+} g_{ba}, \quad \tilde{Q}_{a+} \rightarrow \tilde{Q}'_{a+} = (g^\dagger)_{ab} \tilde{Q}_{b+}, \quad \Phi_+ \rightarrow \Phi_+, \quad V_+ \rightarrow V_+, \quad g \in U(n).
\]

(2.2)

\(^1\)We use mostly the conventions of Wess and Bagger[7] for the \( \mathcal{N} = 1 \) superfields, spinor and other notations.
The subgroup $U(1)$ of $U(n) = U(1) \times SU(n)$ is gauged. The mass parameters $m_a$ break the remaining global symmetry $SU(n)$ to $U(1)^{n-1}$. If $b = 0$ in addition to $m_a = 0$, the $\mathcal{N} = 1$ superfield Lagrangian (2.1) appears to have another global $U_N$ to form an $\mathcal{N} = 1$ vector superfield $V_+$. Summarizing, our model with generic values of $m_a$ has the following $U(1)^n$ symmetries, which are consistent with the $\mathcal{N} = 2$ SUSY:

$$Q_{+a} \rightarrow Q'_{+a} = e^{i\alpha_a}Q_{+a}, \quad \tilde{Q}_{+a} \rightarrow \tilde{Q}'_{+a} = e^{i\beta_a}\tilde{Q}_{+a}, \quad \Phi_+ \rightarrow \Phi'_+ = e^{-i\beta - i\alpha} \Phi_+,$$

(2.3)

with $V_+$ invariant. This invariance respects $\mathcal{N} = 1$, but is inconsistent with the $\mathcal{N} = 2$ SUSY, since the chiral scalar field $\Phi_+$ should have the same transformation as the vector multiplet $V_+$ to form an $\mathcal{N} = 2$ vector multiplet. Summarizing, our model with generic values of $m_a$ has the following $U(1)^n$ symmetries, which are consistent with the $\mathcal{N} = 2$ SUSY:

$$Q_{+a} \rightarrow e^{i\alpha_a}Q_{+a}, \quad \tilde{Q}_{+a} \rightarrow e^{-i\alpha_a}\tilde{Q}_{+a}, \quad \Phi_+ \rightarrow \Phi_+, \quad V_+ \rightarrow V_+.$$  

(2.4)

The diagonal $U(1)$ ($\alpha_1 = \cdots = \alpha_n$) is a local gauged symmetry. Other $U(1)^{n-1}$ groups constrained by $\sum_{a=1}^n \alpha_a = 0$ are global symmetries.

To make the physical content of the theory more transparent, we shall use the Wess-Zumino gauge for the $\mathcal{N} = 1$ vector superfield $V_+$. Then the $\mathcal{N} = 1$ vector superfields can be expanded in terms of Grassmann number $\theta_+$ into component fields

$$V_+(x, \theta_+, \bar{\theta}_+) = -\theta_+ \sigma^m \bar{\theta}_+ v_m(x) + i\theta_+^2 \bar{\theta}_+ \lambda(x) - i\theta_+^2 \bar{\theta}_+ \lambda(x) + \frac{1}{2} \theta_+^2 \bar{\theta}_+^2 X_3(x),$$

(2.5)

where $v_m$, $\lambda$, and $X_3$ are gauge field, gaugino, and auxiliary field, respectively. The $\mathcal{N} = 1$ chiral scalar superfields can also be expanded into components using $y^m = x^m + i\theta_+ \sigma^m \bar{\theta}_+$ as usual:

$$\Phi_+(y, \theta_+) = \phi(y) + \sqrt{2} \theta_+ (-i\sqrt{2} \psi(y)) + \theta_+^2 (X_1(y) + iX_2(y))$$

(2.6)

$$Q_{+a}(y, \theta_+) = q_a(y) + \sqrt{2} \theta_+ \psi_{q_a}(y) + \theta_+^2 F_a(y)$$

(2.7)

$$\tilde{Q}_{+a}(y, \theta_+) = \tilde{q}_a(y) + \sqrt{2} \theta_+ \psi_{\tilde{q}_a}(y) + \theta_+^2 \tilde{F}_a(y).$$

(2.8)

where the scalar fields are denoted by a small letter corresponding to the superfields, such as positively charged scalar $q_a$ as the first component of the superfield $Q_{+a}$. Let us note that the suffix $+$ is not carried by component fields, but is carried only by superfields, which are the functions of the associated Grassmann number $\theta_+$.

In terms of component fields, the bosonic part of this Lagrangian becomes

$$\mathcal{L}_{\text{boson}} = -\frac{1}{4g^2} v_{mm} v^{mn} + \frac{1}{2g^2} (X_3)^2 - \frac{1}{2g^2} |\partial m \phi| + \frac{1}{2g^2} |X_1 + iX_2|^2$$

$$+ \sum_{a=1}^n \left[ |F_a|^2 + |\tilde{F}_a|^2 - |D_m q_a|^2 - |D_m \tilde{q}_a|^2 + X_3 q_a^* q_a - \tilde{q}_a^* \tilde{q}_a \right] - cX_3$$

$$+ \sum_{a=1}^n \left[ (\phi - m_a) q_a \tilde{F}_a + (\phi - m_a) F_a q_a + (X_1 + iX_2) q_a \tilde{q}_a \right] - b \left( X_1 + iX_2 \right)$$

$$+ \sum_{a=1}^n \left[ (\phi^* - m_a^*) q_a^* \tilde{F}_a^* + (\phi^* - m_a^*) F_a^* q_a^* + (X_1 - iX_2) q_a^* \tilde{q}_a^* \right] - b^* \left( X_1 - iX_2 \right),$$

(2.9)
where the field strength \( v_{mn} \) and the covariant derivatives \( D_m \) are defined by
\[
v_{mn} = \partial_m v_n - \partial_n v_m, \quad D_m q_a \equiv (\partial_m + iv_m)q_a, \quad D_m \bar{q}_a \equiv (\partial_m - iv_m)\bar{q}_a, \tag{2.10}\]
respectively. The entire Lagrangian including the fermions will be given in sect.4 where the full \( \mathcal{N} = 2 \) SUSY will be clarified. We see that scalar field \( q_a \) with the \( U(1) \) charge +1 and \( \bar{q}_a \) with charge −1 have a complex mass \( m_a \). Since a complex mass common to all the flavors can be absorbed by shifting the neutral complex scalar field \( \phi \), these mass parameters can always be chosen to satisfy
\[
\Sigma_{a=1}^n m_a = 0. \tag{2.11}\]
The real FI parameter \( c \) of the \( D \)-term is usually called the electric FI parameter, and the complex parameter \( b \) appearing in the F-term is called the magnetic FI parameter [34].

The SUSY auxiliary fields \( X_1, X_2, X_3, F_a, \bar{F}_a \) can be eliminated by solving their algebraic equations of motion
\[
X_3 = -g^2 \left\{ \sum_{a=1}^n (|q_a|^2 - |\bar{q}_a|^2) - c \right\} \tag{2.12}
\]
\[
X_1 + iX_2 = -2g^2 \left( \sum_{a=1}^n q_a \bar{q}_a - b^* \right) \tag{2.13}
\]
\[
F_a = - (\phi - m_a^*) \bar{q}_a \tag{2.14}
\]
\[
\bar{F}_a = - (\phi - m_a^*) q_a^*. \tag{2.15}
\]

Then, the Lagrangian is given entirely in terms of physical fields
\[
\mathcal{L}_{\text{boson}} = -\frac{1}{4g^2}v_{mn}v^{mn} - \frac{1}{2g^2} |\partial\phi|^2 - \sum_{a=1}^n \left( |Dq_a|^2 + |D\bar{q}_a|^2 \right) - 2g^2 |\sum_{a=1}^n q_a \bar{q}_a - b|^2
\]
\[
- \sum_{a=1}^n \left| \phi - m_a \right|^2 \left( |q_a|^2 + |\bar{q}_a|^2 \right) - \frac{g^2}{2} \left( \sum_{a=1}^n (|q_a|^2 - |\bar{q}_a|^2) - c \right)^2. \tag{2.16}\]

A similar model has been considered previously in a different context [28–30].

SUSY vacua are given by vanishing auxiliary fields: \( X_1 = X_2 = X_3 = 0 \), and \( F_a = \bar{F}_a = 0 \). In the generic case of distinct complex mass parameters \( m_a \neq m_b \) for \( a \neq b \), we find precisely \( n \) isolated SUSY vacua (and no other vacua). We denote the modulus and phase of the magnetic FI parameter \( b \) by two real parameters \( h > 0 \) and \( \beta \) as
\[
b = h^2 e^{i\beta}. \tag{2.17}\]

The \( i \)-th vacuum is characterized by nonvanishing values of \( q_a, \bar{q}_a \) and \( \phi \):
\[
\phi = m_a, \quad q_a = \sqrt{\frac{c^2 + 4h^4 + c}{2}} e^{i\alpha_a}, \quad \bar{q}_a = \sqrt{\frac{c^2 + 4h^4 - c}{2}} e^{i(\beta - \alpha_a)}, \tag{2.18}\]
with vanishing values for the remaining hypermultiplets \( q_b^* = \bar{q}_b^* = 0 \) \( (a \neq b) \). At the \( a \)-th vacuum, the phase \( \alpha_a \) is fixed breaking a \( U(1) \) symmetry which is a linear combination of local gauged \( U(1) \) and other global \( U(1)^{n-1} \) generators. Because of Higgs mechanism, gauge boson should become massive in the vacuum. However, there still remains \( U(1)^{n-1} \) global symmetries \( \alpha_e, \ b \neq a \) unbroken as given in Eq. (2.4).

The Hamiltonian corresponding to the Lagrangian (2.16) is given by

\[
\mathcal{H} = \frac{1}{2g^2}(v_{01}^2 + v_{02}^2 + v_{03}^2 + v_{12}^2 + v_{13}^2 + v_{23}^2) + \frac{1}{2g^2} \left[ |\partial_0 \phi|^2 + |\partial_1 \phi|^2 + |\partial_2 \phi|^2 + |\partial_3 \phi|^2 \right] \\
+ \sum_{a=1}^{n} \left[ |D_0 q_a|^2 + |D_1 q_a|^2 + |D_2 q_a|^2 + |D_3 q_a|^2 + |D_0 \bar{q}_a|^2 + |D_1 \bar{q}_a|^2 + |D_2 \bar{q}_a|^2 + |D_3 \bar{q}_a|^2 \right] \\
+ \sum_{a=1}^{n} \left| \phi - m_a \right|^2 \left( |q_a|^2 + |\bar{q}_a|^2 \right) + g^2 \left\{ \sum_{a=1}^{n} (|q_a|^2 - |\bar{q}_a|^2) - c \right\}^2 + 2g^2 \sum_{a=1}^{n} q_a \bar{q}_a - b^2. \tag{2.19} \]

Since the static domain wall junctions has nontrivial dependence only in two-dimensional spatial coordinates, we shall look for field configurations as a function of \( x^1 \) and \( x^2 \) coordinates and introduce complex coordinates \( z = x^1 + i x^2 \), \( \bar{z} = x^1 - i x^2 \), \( \partial_z = \frac{1}{2} (\partial_1 - i \partial_2) \), and \( \partial_{\bar{z}} = \frac{1}{2} (\partial_1 + i \partial_2) \). We also wish to maintain \( 1 + 1 \) dimensional Lorentz invariance in the \( x^0, x^3 \) plane. Therefore we need to require

\[
v_0 = v_3 = 0, \quad v_{01} = v_{02} = v_{03} = v_{13} = v_{23} = 0, \tag{2.20}\\
\partial_0 \phi = 0, \quad D_0 q_a = 0, \quad D_0 \bar{q}_a = 0, \tag{2.21}\\
\partial_3 \phi = 0, \quad D_3 q_a = 0, \quad D_3 \bar{q}_a = 0. \tag{2.22} \]

In order to find a minimum energy configuration for a given boundary condition, we form complete squares \([19],[20]\) in the energy density functional \( \mathcal{E} \) by introducing an arbitrary phase \( \Omega \), |\( \Omega \)| = 1

\[
\mathcal{E} = \frac{1}{2g^2} \left[ v_{12} + g^2 \left\{ \sum_{a=1}^{n} (|q_a|^2 - |\bar{q}_a|^2) - c \right\}^2 + \frac{2}{g^2} |\partial_z \phi - g^2 \Omega (\sum_{a=1}^{n} q_a \bar{q}_a - b^*)|^2 \right] \\
+ 4 \sum_{a=1}^{n} \left| D_z q_a - \frac{1}{2} \Omega (\phi^* - m_a) q_a^* \right|^2 + 4 \sum_{a=1}^{n} \left| D_z \bar{q}_a - \frac{1}{2} \Omega (\phi^* - m_a) q_a \right|^2 \\
+ c v_{12} + \partial_z (2 \Omega^* (\sum_{a=1}^{n} (\phi^* - m_a) q_a \bar{q}_a - b \phi)) + \partial_{\bar{z}} (2 \Omega^* (\sum_{a=1}^{n} (\phi^* - m_a) q_a \bar{q}_a - b \phi^*)) \\
+ \sum_{a=1}^{n} \left[ \partial_z (q_a D_z q_a - q_a (D_z q_a)^*) + \partial_z (q_a D_z \bar{q}_a - \bar{q}_a (D_z \bar{q}_a)^*) \right] \\
+ \sum_{a=1}^{n} \left[ \partial_{\bar{z}} (q_a D_z q_a)^* - q_a^* D_z q_a + \partial_{\bar{z}} (q_a D_z \bar{q}_a)^* - q_a^* D_z \bar{q}_a \right] \\
+ \frac{1}{2g^2} \partial_z (\phi^* \partial_z \phi - \phi \partial_z \phi^*) + \frac{1}{2g^2} \partial_{\bar{z}} (\phi \partial_{\bar{z}} \phi^* - \phi^* \partial_{\bar{z}} \phi). \tag{2.23} \]

6
The last four lines are total derivatives that give surface terms when integrated over the entire $x^1, x^2$ plane. Since all the remaining terms are the complete squares, we find that the integrated energy over $x^1, x^2$ plane of the field configuration is always larger than the surface terms, which are completely determined by the boundary condition at spatial infinity. This bound is called the BPS bound and is saturated by requiring the complete squares to vanish

$$v_{12} = -g^2 \left\{ \sum_{a=1}^{n} (|q_a|^2 - |\bar{q}_a|^2) - c \right\} \quad (2.24)$$

$$\frac{1}{g^2} \frac{\partial \phi}{\partial z} = \Omega \left( \sum_{a=1}^{n} q_a^* \bar{q}_a - b^* \right) \quad (2.25)$$

$$2D_z q_a = \Omega (\phi^* - m^*_a) \bar{q}_a \quad (2.26)$$

$$2D_z \bar{q}_a = \Omega (\phi^* - m^*_a) q_a. \quad (2.27)$$

These first order differential equations are called the BPS equations. Since the surface terms depend on $\Omega$, the phase factor $\Omega$ can be chosen to obtain the best bound. Let us also note that the minimum energy configurations automatically satisfy the equations of motion \[19\], \[20\].

Since the Lagrangian (2.1) with $\mathcal{N} = 1$ superfield exhibits the $\mathcal{N} = 1$ SUSY manifestly, we can formulate the condition of partial conservation of SUSY. We will see that the above minimum energy conditions (2.24)–(2.27) are precisely the conditions to conserve one out of four SUSY. We need to consider only the SUSY transformations of fermions, since only bosonic fields can have nonvanishing values. The $\mathcal{N} = 1$ SUSY transformations of gaugino is given by \[47\]

$$\delta \xi_+ \lambda = \sigma^{mn} v_{mn} \xi_+ + iX_3 \xi_+$$

$$= \begin{bmatrix} v_{03} - iv_{12} + iX_3 & v_{01} + v_{13} - iv_{23} - iv_{02} \\ v_{01} - v_{13} - iv_{23} + iv_{02} & -v_{03} + iv_{12} + iX_3 \end{bmatrix} \begin{bmatrix} \xi_+ \\ \xi_+ \end{bmatrix}. \quad (2.28)$$

If we require that a part of SUSY corresponding to the upper component $\xi_+ = 0$, we find \[20\]

$$v_{12} = X_3, \quad v_{03} = 0, \quad v_{01} = v_{13}, \quad v_{23} = v_{02}. \quad (2.29)$$

Using the algebraic equation of motion for the auxiliary field (2.12), the minimum energy condition (2.24) for the vector multiplet is precisely the same as the condition of partial SUSY conservation condition (2.29). Similarly the $\mathcal{N} = 1$ SUSY transformations of fermions in chiral scalar multiplets are given by \[47\]

$$\delta \xi_+ (-i\sqrt{2} \psi) = i\sqrt{2} \sigma^m \partial_m (\phi - m_a) \xi_+ + \sqrt{2} (X_1 + iX_2) \xi_+, \quad (2.30)$$

$$\delta \xi_+ \psi_{qa} = i\sqrt{2} \sigma^m \mathcal{D}_m q_a \xi_+ + \sqrt{2} F_a \xi_+, \quad (2.31)$$

$$\delta \xi_+ \bar{\psi}_{\bar{q}a} = i\sqrt{2} \sigma^m \mathcal{D}_m \bar{q}_a^* \xi_+ + \sqrt{2} \bar{F}_a \xi_. \quad (2.32)$$

For these transformations, we express derivatives in terms of complex coordinates, assuming $x_1, x_2$ dependence only

$$\sigma^m \partial_m = (\sigma^+ + i\sigma^2) \frac{1}{2} (\partial_1 - i\partial_2) + (\sigma^+ - i\sigma^2) \frac{1}{2} (\partial_1 + i\partial_2),$$

$$= 2\sigma^+ \partial_z + 2\sigma^- \partial_{\bar{z}}, \quad (2.33)$$
where $\sigma^+ \equiv (\sigma^1 + i\sigma^2)/2$, $\sigma^- \equiv (\sigma^1 - i\sigma^2)/2$. If we require conservation of only one out of four SUSY specified by

$$-\Omega \sigma^+ \xi_+ = i \xi_+ \quad \text{and} \quad \sigma^- \xi_+ = 0,$$

we obtain [20]

$$\delta_{\xi_+}(-i\sqrt{2}\psi) = i\sqrt{2}\sigma^m \partial_m (\phi - m_a) \xi_+ + \sqrt{2}(X_1 + iX_2)\xi_+,$$

$$= i2\sqrt{2}(\sigma^+ \partial_x + \sigma^- \partial_{\bar{x}})(\phi - m_a) \xi_+ + \sqrt{2}(X_1 + iX_2)\xi_+,$$

$$= \sqrt{2}\xi_+ \Omega^{-1} (2\partial_x (\phi - m_a) + \Omega(X_1 + iX_2)). \quad (2.35)$$

The algebraic equations of motion for auxiliary fields are given in terms of the superpotential $P$

$$X_1 + iX_2 = -\left(\frac{\partial P}{\partial \phi}\right)^*, \quad F_a = -\left(\frac{\partial P}{\partial q_a}\right)^*, \quad \tilde{F}_a = -\left(\frac{\partial P}{\partial \bar{q}_a}\right)^* \quad (2.39)$$

This superpotential $P$ as a function of scalar fields is given in our case by

$$P = \sum_{a=1}^n (\phi - m_a) q_a \bar{q}_a - b\phi. \quad (2.40)$$

Using the superpotential (2.40) and the algebraic equations of motion for auxiliary fields (2.39), we see that the minimum energy conditions (2.25)–(2.27) are precisely the same as the conditions for the conservation of one out of four SUSY (2.38).

### 3 Domain wall junction

In order to obtain an exact solution of junctions, we shall embed the known solution to a solution of BPS equations (2.24)–(2.27) for our $\mathcal{N} = 2$ SUSY massive multiflavor QED. We are making one complex structure out of three manifest by using the $\mathcal{N} = 1$ superfield formalism. Although three FI parameters $c$, $b_1 = \text{Re} \, b$, and $b_2 = \text{Im} \, b$ form an $SU(2)_R$ triplet, the choice of particular
complex structure made the $SU(2)_R$ symmetry not visible. In this circumstance, we find it convenient to choose the FI parameters $c$ and $b$ in Eq. (2.17) as
\[ c = 0, \quad b = h^2 \in \mathbb{R} \quad (\beta = 0). \] (3.1)

Then the $a$-th vacuum values of fields in Eq. (2.18) become
\[ \phi = m_a, \quad q_b = h e^{i\alpha_b \delta_{ba}}, \quad \tilde{q}_b = h e^{-i\alpha_b \delta_{ba}}. \] (3.2)

Since the known junction solution was obtained [20] with the $Z_3$ symmetry for SUSY vacua and with a relation between the vacuum values of charged chiral scalar fields $q_a, \tilde{q}_a$ and the neutral scalar fields $\phi$, we should require $n = 3$ flavors with $Z_3$ symmetry, and a relation between the mass scales of $m_b$ and the Fayet-Iliopoulos term $h$. Altogether we assume the following particular values for the parameters of our model:
\[ m_b = \frac{2gh}{\sqrt{3}} e^{i\frac{2\pi}{3}b}, \quad b = 1, 2, 3. \] (3.3)

The resulting vacua are illustrated in the complex $\phi$ plane in Fig. 1.

---

**Figure 1:** Vacua in complex $\phi$ plane.

**Figure 2:** The $Z_3$ junction in real space $z = x^1 + ix^2$.

Combined with the algebraic equation of motion for auxiliary field $X_3$ (2.12) for vector superfield, we can satisfy the BPS equations (2.29) trivially by choosing
\[ v_0(x^1, x^2) = v_3(x^1, x^2) = 0, \quad |q_b(x^1, x^2)| = |\tilde{q}_b(x^1, x^2)|. \] (3.4)

Suggested by this condition, we assume the following relation between values of hypermultiplets [20]
\[ q_b(x^1, x^2)e^{-i\alpha_b} = \tilde{q}_b(x^1, x^2)e^{i\alpha_b} \in \mathbb{R}, \] (3.5)
in accordance with the vacuum values (3.2) which should be reached at infinity. This assumption will be justified a posteriori after finding out solutions. We are interested in a BPS junction configuration separating three vacuum domains \( a = 1, 2, 3 \) with \( Z_3 \)-symmetry, where the third vacuum is placed at infinity along the positive real axis as illustrated in Fig.2. This configuration corresponds to the choice of the phase factor \( \Omega = -1 \) \cite{20}, \cite{21}. Now the remaining BPS equations for chiral scalar multiplets (hypermultiplets and the chiral scalar \( \phi \) in the \( \mathcal{N} = 2 \) vector multiplet) read

\[
2 \frac{\partial q_b}{\partial z} = \left( \frac{2gh}{\sqrt{3}} e^{i \frac{2\pi}{3} b} - \phi \right)^* q_b, \quad (3.6)
\]

\[
\frac{1}{g^2} \frac{\partial \phi}{\partial z} = h^2 - \sum_{b=1}^{3} |q_b|^2. \quad (3.7)
\]

We define a dimensionless complex coordinate \( \hat{z} \) and real dimensionless fields \( \hat{q}_b \) and \( \hat{\phi} \) by rescaling with the normalization factor associated to the vacuum values as

\[
z = \frac{\sqrt{3}}{2} \frac{1}{gh} \hat{z}, \quad q_b = he^{i\alpha_b} \hat{q}_b, \quad \phi = \frac{2}{\sqrt{3}} gh \hat{\phi}. \quad (3.8)
\]

Then the BPS equations become

\[
2 \frac{\partial \hat{q}_b}{\partial \hat{z}} = \left( e^{i \frac{2\pi}{3} b} \hat{\phi} \right)^* \hat{q}_b, \quad \hat{q}_b \in \mathbb{R}, \quad (3.9)
\]

\[
2 \frac{\partial \hat{\phi}}{\partial \hat{z}} = 3 \left( 1 - \sum_b \hat{q}_b^2 \right). \quad (3.10)
\]

We can now recognize the familiar form of the BPS equation allowing the junction as an exact solution \cite{20}, \cite{24}. Let us define the following auxiliary quantities \( f_b \)

\[
f_b = \exp \left( \frac{1}{2} \left( e^{-i \frac{2\pi}{3} b} \hat{z} + e^{i \frac{2\pi}{3} b} \hat{z}^* \right) \right), \quad b = 1, 2, 3. \quad (3.11)
\]

Following identities can be derived for these auxiliary quantities \cite{20}, \cite{24}

\[
2 \frac{\partial}{\partial \hat{z}} \left( \frac{f_b}{f_1 + f_2 + f_3} \right) = \left( e^{-i \frac{2\pi}{3} b} - \sum_c e^{-i \frac{2\pi}{3} c} f_c \right) \frac{f_b}{f_1 + f_2 + f_3}, \quad (3.12)
\]

\[
2 \frac{\partial}{\partial \hat{z}} \left( \sum_b e^{i \frac{2\pi}{3} b} f_b \frac{1}{f_1 + f_2 + f_3} \right) = 3 \frac{\sum_b f_b^2}{(f_1 + f_2 + f_3)^2} \left( 1 - \sum_b \frac{f_b^2}{(f_1 + f_2 + f_3)^2} \right). \quad (3.13)
\]

Therefore we find the solutions for the BPS equations

\[
\hat{q}_b = \frac{f_b}{f_1 + f_2 + f_3}, \quad \hat{\phi} = \frac{\sum_b e^{i \frac{2\pi}{3} b} f_b}{f_1 + f_2 + f_3}. \quad (3.14)
\]
These solutions can be rewritten in terms of our original variables $q_b$, $\phi$, and $z$,

$$q_b = h e^{i\alpha_b} \frac{f_b}{f_1 + f_2 + f_3}, \quad (3.15)$$

$$\tilde{q}_b = h e^{-i\alpha_b} \frac{f_b}{f_1 + f_2 + f_3}, \quad (3.16)$$

$$\phi = \frac{2gh}{\sqrt{3}} \sum_b e^{i\frac{2\pi}{3}b} f_b, \quad (3.17)$$

where

$$f_b = \exp \left( \frac{2gh}{\sqrt{3}} \left( e^{-i\frac{2\pi}{3}b} z + e^{i\frac{2\pi}{3}b} z^* \right) \right), \quad b = 1, 2, 3. \quad (3.18)$$

By rotating the field configuration by $\frac{2\pi}{3}$, we find that the solution (3.18) is precisely the same field configuration of the previous junction solution in the $\mathcal{N} = 1$ SUSY theory [20] provided the dimensionful parameter $h$ is related to the parameter $\Lambda$ through $h = \sqrt{2}\Lambda$. The field $\phi$ of the junction configuration takes values inside the triangle connecting the three vacua as illustrated in Fig.1. The straight line segments between vacua on the $\phi$ plane correspond to the spatial infinity in real space $z \to \infty$. As an illustration of the asymptotic behavior of the junction configuration, values of fields $q_a$, $a = 1, 2, 3$ choosing $\alpha_a = 0$ are shown along the real axis in Fig.3. The energy density computed analytically in our previous solution of $\mathcal{N} = 1$ SUSY model [20] can easily be converted into our case of $\mathcal{N} = 2$ SUSY QED. The energy density of the junction solution in our $\mathcal{N} = 2$ SUSY model is shown in Fig.4.

Let us discuss zero modes on this junction solution. First we notice that we have two massless boson corresponding to spontaneously broken translation in two directions $x^1, x^2$. As for the

![Figure 3](image-url)
Figure 4: Energy density of the $Z_3$ junction configuration of $\mathcal{N} = 2$ SUSY QED.

two global $U(1)$ symmetries, both of them are broken on the junction solution, and there are two corresponding Nambu-Goldstone bosons. For each domain wall, only two hypermultiplets have nontrivial field configurations. Therefore only one of the two $U(1)$ global symmetries is spontaneously broken and only one Nambu-Goldstone boson associated with the $U(1)$ phase rotation appears [29], [30]. Since the broken $U(1)$ on each wall is a different combination of the two $U(1)$ global symmetries, the associated Nambu-Goldstone boson on each wall is also a different linear combination of these two. This situation is very similar to the property of the Nambu-Goldstone fermions on the junction that was observed previously [20], [21]. As for fermions, six out of eight SUSY are spontaneously broken. Therefore we have six Nambu-Goldstone fermions. On each wall, only four SUSY are broken. Therefore there are four Nambu-Goldstone fermions on each wall. These four Nambu-Goldstone fermions on each wall are different linear combinations of six Nambu-Goldstone fermions on the junction as a whole, since different linear combinations of SUSY charges are broken on each wall. This situation is analogous to the previously obtained junction solution in $\mathcal{N} = 1$ theory [20], [21]. It has been pointed out that the wall provides a model to "localize" gauge bosons [29]. In this context, one should note that one out of two Nambu-Goldstone bosons associated with the two $U(1)$'s mixes with the $U(1)$ gauge boson. We postpone a more detailed analysis of walls and junctions involving the gauge field for subsequent publications.

Let us discuss possible junction configurations in other models. We have obtained an analytic solution of junction provided the mass parameters $m_i$ are tuned to the gauge coupling $g$ and the Fayet-Iliopoulos parameter $b = h^2$ as in Eq.(3.3). However, it is almost clear from continuity that junction configurations should exist even if the mass parameters are perturbed infinitesimally away from the tuned values. On the other hand, junction configuration is not allowed if masses
are on a straight line in complex mass plane, for instance if all masses are aligned on the real axis. Since the tension is given by the absolute value of the difference of superpotential, the possible junction configuration cannot satisfy the condition of mechanical stability \[19\] if masses are on a straight line. We conjecture that there should exist a “line of marginal stability” somewhere between the \(Z_3\) symmetric mass parameters and mass parameters on a straight line. The junction should become unstable across the line of marginal stability. An interesting example of the line of marginal stability for domain walls has been explicitly demonstrated for a similar Wess-Zumino model \[49\]. Lines of marginal stability for monopoles and dyons are well-studied in the \(N = 2\) gauge theories \[50\].

It is also likely that there exist junctions of \(n\)-walls, although it is difficult to work out explicit solutions except in particular nonlinear sigma models \[24\]. Similarly to our previous solution \[21\], the Nambu-Goldstone bosons and fermions on our junction solution are not normalizable. Therefore usual wisdom of a low-energy effective Lagrangian approach cannot be applied easily to our junction solution. However, it has been observed that a graviton zero mode is localized at the junction, if the bulk spacetime is warped \[51\]. This is due to a suppression factor produced by the bulk AdS(-like) spacetime. It may also be possible to exploit this mechanism to obtain normalizable zero modes localized at the junction. The general properties of possible junction solutions has been studied in the presence of gravity \[23\], although no explicit solution has been obtained. This is an interesting subject in future.

4 8 SUSY transformations

In this and the following sections, we return to the general case of \(n\)-flavors. There have been a number of studies to formulate the \(N = 2\) SUSY field theories in five dimensions in terms of \(N = 1\) superfield formalism \[5, 41-44\]. Inspired by these studies \[5, 41, 42\], we will redefine auxiliary fields of chiral scalar fields for a hypermultiplet \(2.7\) and \(2.8\) in order to identify all the eight supersymmetry transformations

\[
F_a = F'_a - (\phi^* - m_a^*) \tilde{q}_a^* , \quad \tilde{F}_a = -\tilde{F}'_a - q_a^* (\phi^* - m_a^*) \tag{4.1}
\]

Instead of Eqs.\(2.3\)-\(2.8\), the component expansions of superfields in powers of Grassmann number \(\theta_+\) now read

\[
V_+ = -\theta_+ \sigma^m \bar{\theta}_+ v_m + i \theta_+ ^2 \bar{\theta}_+ \bar{\lambda} - i \theta_+ ^2 \theta_+ \lambda + \frac{1}{2} \theta_+ ^2 \bar{\theta}_+ ^2 X_3 \tag{4.2}
\]

\[
\Phi_+ = \phi + \sqrt{2} \theta_+ (\phi^* - m_a^*) q_a + \theta_+ ^2 (F'_a - q_a^* (\phi^* - m_a^*)) \tag{4.3}
\]

\[
Q_{+a} = q_a + \sqrt{2} \theta_+ \psi_{q_a} + \theta_+ ^2 (F'_a - q_a^* (\phi^* - m_a^*)) \tag{4.4}
\]

\[
\tilde{Q}_{+a} = \tilde{q}_a + \sqrt{2} \theta_+ \bar{\psi}_{\tilde{q}_a} + \theta_+ ^2 (\tilde{F}'_a - q_a^* (\phi^* - m_a^*)) \tag{4.5}
\]

In terms of these component fields, the full Lagrangian \[241\] is given by

\[
\mathcal{L} = -\frac{1}{4g^2} v_m v^m - \frac{1}{g^2} i \lambda \sigma^m \partial_m \bar{\lambda} + \frac{1}{2g^2} (X_3)^2 + \frac{1}{2g^2} |X_1 + iX_2|^2
\]
This is the full Lagrangian including fermion terms compared to the bosonic one in Eq. (2.9).

To obtain the $\mathcal{N} = 1$ SUSY transformations in the Wess-Zumino gauge, one has to combine an ordinary SUSY transformation with an accompanying gauge transformation to preserve the Wess-Zumino gauge. Let us consider the four SUSY transformations $\delta \xi_+^a$ given by an infinitesimal Grassmann number $\xi_+$ in the direction of $\theta_+$ in Eqs. (4.12)–(4.15). The $\mathcal{N} = 2$ vector multiplet represented by superfields (4.2) and (4.3) transforms under the infinitesimal SUSY $\delta \xi_+$ along the left-handed spinor $\theta_+$ in the Wess-Zumino gauge as (47)

$$
\delta \xi_+, v^m = i\bar{\xi}_+ \sigma^m \lambda + i\xi_+ \sigma^m \bar{\lambda},
$$

$$
\delta \xi_+, \lambda = \sigma^{mn} \psi_{mn} \xi_+ + iX_3 \xi_+, \quad (4.8)
$$

$$
\delta \xi_+, X_3 = \bar{\xi}_+ \bar{\sigma}^m \partial_m \lambda - \xi_+ \sigma^m \partial_m \bar{\lambda},
$$

$$
\delta \xi_+, \phi = \sqrt{2} \xi_+ (-i \sqrt{2} \psi),
$$

$$
\delta \xi_+ (-i \sqrt{2} \psi) = i\sqrt{2} \sigma^m \partial_m \phi \xi_+ + \sqrt{2}(X_1 + iX_2) \xi_+, \quad (4.11)
$$

$$
\delta \xi_+ (X_1 + iX_2) = i\sqrt{2} \xi_+ \bar{\sigma}^m \partial_m (-i \sqrt{2} \psi). \quad (4.12)
$$

Similarly we obtain the supersymmetry transformation rules for hypermultiplets in the Wess-Zumino gauge

$$
\delta \xi_+, q_a = \sqrt{2} \xi_+ \psi_{qa}, \quad (4.13)
$$

$$
\delta \xi_+, \psi_{qa} = i\sqrt{2} \sigma^m \mathcal{D}_m q_a \xi_+ + \sqrt{2}(F_a^* - (\phi^* - m_a^*) \bar{q}_a^*) \xi_+, \quad (4.14)
$$

$$
\delta \xi_+, (F_a^* - (\phi^* - m_a^*) \bar{q}_a^*) = i\sqrt{2} \xi_+ \sigma^m \mathcal{D}_m \psi_{qa} + 2i \bar{\xi}_+ \bar{\lambda} q_a, \quad (4.15)
$$

14
\[ \delta \xi_+ \bar{q}_a = \sqrt{2} \xi_+ \psi \bar{q}_a, \quad (4.16) \]
\[ \delta \xi_+ \psi \bar{q}_a = i \sqrt{2} \sigma^m \partial_m \bar{q}_a \xi_+ + \sqrt{2} (\bar{F}_a^\prime - (\phi^* - m^*_a) q_a^*) \xi_+, \quad (4.17) \]
\[ \delta \xi_+ (\bar{F}_a^\prime - (\phi^* - m^*_a) q_a^*) = i \sqrt{2} \bar{\xi}_+ \bar{\sigma}^m \partial_m \psi \bar{q}_a - 2i \bar{\xi}_+ \bar{\lambda} \bar{q}_a, \quad (4.18) \]

Let us now consider the remaining four SUSY transformations \( \delta \xi_\) also along the left-handed spinor \( \theta_\) which are not manifest by the above \( N = 1 \) superfield formalism \((2.1)\) or the associated component formalism \((4.6)\) in the Wess-Zumino gauge. Please note that Grassmann numbers \( \xi_-, \theta_\) are left-handed chiral spinors, similarly to \( \xi_+, \theta_\) and not to be confused with the right-handed spinors. To work out \( \delta \xi_- \), we shall follow the classical method of Fayet \cite{34}. The two sets of left-handed chiral spinor Grassmann numbers \( \theta_\) and \( \theta_- \) should form a doublet under an internal \( SU(2)_R \) transformation \( M \) whose representation matrix is denoted by a \( 2 \times 2 \) matrix \( M \)

\[ M \begin{pmatrix} \theta_+ \\ \theta_- \end{pmatrix} M^{-1} = M \begin{pmatrix} \theta_+ \\ \theta_- \end{pmatrix}. \quad (4.19) \]

It is enough to consider a discrete transformation, for instance, an \( SU(2)_R \) rotation around second axis by \( \pi \)

\[ M_0 = \exp(i \pi I_2), \quad M_0 = i \sigma_2 = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}. \quad (4.20) \]

Following Fayet, we demand that \( v_m, \phi, \psi_{q_a} \), and \( \psi_{\bar{q}_a} \) be \( SU(2)_R \) singlets, and that

\[ \begin{pmatrix} \lambda \\ \psi \end{pmatrix}, \quad \begin{pmatrix} q_a \\ \bar{q}_a \end{pmatrix}, \quad \begin{pmatrix} -\bar{q}_a \\ q_a^* \end{pmatrix}, \quad (4.21) \]

be \( SU(2)_R \) doublets. The equations of motion \((2.12)\) and \((2.13)\) show that

\[ \begin{pmatrix} X_1 \\ -X_2 \\ X_3 \end{pmatrix} \quad (4.22) \]
transforms as an \( SU(2)_R \) triplet. The equations of motion for auxiliary fields in Eqs. \((2.14)\), \((2.15)\) for the hypermultiplets become

\[ F'_a^* = F_a^* + (\phi^* - m_a) \bar{q}_a = 0, \quad \bar{F}'_a = -\bar{F}_a - q_a^* (\phi^* - m^*_a) = 0, \quad (4.23) \]

which suggest that the auxiliary fields in hypermultiplets generate \( SU(2)_R \) doublets

\[ \begin{pmatrix} (m^*_a - \phi^*) q_a - \bar{F}_a^* \\ (m^*_a - \phi^*) \bar{q}_a + F_a^* \end{pmatrix} \sim \begin{pmatrix} q_a \\ \bar{q}_a \end{pmatrix}, \quad \begin{pmatrix} -\bar{q}_a (m^*_a - \phi^*) - F_a^* \\ q_a^* (m_a^* - \phi^*) - \bar{F}_a^* \end{pmatrix} \sim \begin{pmatrix} -\bar{q}_a \\ q_a^* \end{pmatrix}. \quad (4.24) \]

By applying the discrete \( SU(2)_R \) internal transformation \((4.20)\) with the assignment \((4.21)\)–\((4.24)\) to the first four SUSY \((4.7)\)–\((4.18)\) represented by an infinitesimal \( \xi_+ \), we can find the
another four SUSY $\delta_\xi -$ corresponding to an infinitesimal Grassmann number $\xi_-$. We obtain $\delta_\xi -$ transformations of the $\mathcal{N} = 2$ vector multiplet in the Wess-Zumino gauge

$$\delta_\xi - v^m = i\xi_- \sigma^m \bar{\psi} + i\bar{\xi}_- \bar{\sigma}^m \psi$$  \hspace{1cm} (4.25)
$$\delta_\xi - \psi = \sigma^{mn} v_{mn} \xi_- - iX_3 \xi_-.$$  \hspace{1cm} (4.26)
$$\delta_\xi - (-X_3) = -\xi_- \sigma^m \partial_m \bar{\psi} + \bar{\xi}_- \bar{\sigma}^m \partial_m \psi$$ \hspace{1cm} (4.27)
$$\delta_\xi - \phi = \sqrt{2} \xi_- (\sqrt{2} i \lambda),$$ \hspace{1cm} (4.28)
$$\delta_\xi - (\sqrt{2} i \lambda) = i \sqrt{2} \sigma^m \partial_m \phi \xi_- + \sqrt{2} (-X_1 + iX_2) \xi_-,$$ \hspace{1cm} (4.29)
$$\delta_\xi - (-X_1 + iX_2) = i \sqrt{2} \xi_- \sigma^m \partial_m (\sqrt{2} i \lambda).$$ \hspace{1cm} (4.30)

The transformation laws for the $\mathcal{N} = 2$ vector multiplet (4.17)–(4.12) with respect to the Grassmann numbers $\theta_+$ and (4.25)–(4.30) with respect to the Grassmann numbers $\theta_-$ are illustrated in Fig.5. For hypermultiplet, $\delta_\xi -$ transformations in the Wess-Zumino gauge is given by

$$\delta_\xi - \tilde{q}_a^* = \sqrt{2} \xi_- \psi_{q_a}, \hspace{1cm} (4.31)$$
$$\delta_\xi - \psi_{q_a} = i \sqrt{2} \sigma^m \mathcal{D}_m \tilde{q}_a \xi_- + \sqrt{2} \left( F_{a}^{\prime*} + (\phi^* - m^*_a) q_a \right) \xi_-,$$ \hspace{1cm} (4.32)
$$\delta_\xi - \left( F_{a}^{\prime*} + (\phi^* - m^*_a) q_a \right) = i \sqrt{2} \xi_- \bar{\sigma}^m \mathcal{D}_m \psi_{q_a} + 2 i \bar{\xi}_- \bar{\psi}_{q_a},$$ \hspace{1cm} (4.33)
$$\delta_\xi - q_a^* = \sqrt{2} \xi_- (-\psi_{q_a}), \hspace{1cm} (4.34)$$
$$\delta_\xi - (-\psi_{q_a}) = i \sqrt{2} \sigma^m \mathcal{D}_m q_a \xi_- + \sqrt{2} (-F_{a}^{\prime*} + (\phi^* - m^*_a) \bar{q}_a) \xi_-,$$ \hspace{1cm} (4.35)
$$\delta_\xi - (-F_{a}^{\prime*} + (\phi^* - m^*_a) \bar{q}_a) = i \sqrt{2} \xi_- \bar{\sigma}^m \mathcal{D}_m (-\psi_{q_a}) - 2 i \bar{\xi}_- \bar{\psi}_{q_a}.$$ \hspace{1cm} (4.36)

The transformation law for the $\mathcal{N} = 2$ hypermultiplet (4.13)–(4.18) with respect to the Grassmann numbers $\theta_+$ and (4.31)–(4.33) with respect to the Grassmann numbers $\theta_-$ are illustrated in Fig.6
The following superfields using another set of Grassmann number \( \theta \)
\( \delta \) manifests
\( \xi \) in (4.25)–(4.33) manifest

To summarize the transformation property (4.25)–(4.33) under \( \delta \), it is convenient to define the following superfields using another set of Grassmann number \( \theta \)

\[
V_- = -\theta_- \sigma^m \bar{\theta} \, v_m + i \theta_2 \bar{\theta} \, \psi - i \bar{\theta}_2 \, \theta \, \psi + \frac{1}{2} \theta_2 \bar{\theta}_2 (-X_3),
\]
\[
\Phi_- = \phi + \sqrt{2} \theta_2 (\sqrt{2} i \lambda) + \theta_2 (-X_1 + iX_2),
\]
\[
Q_{-a} = \bar{q}_a^* + \sqrt{2} \theta_- \psi_{qa} + \theta_-^2 (\bar{F}_a^* + (\phi^* - m_a^*) q_a),
\]
\[
\bar{Q}_{-a} = q_a^* + \sqrt{2} \theta_- (-\psi_{qa}) + \theta_-^2 (-\bar{F}_a^* + \bar{q}_a (\phi^* - m_a^*)).
\]

We can now rewrite the Lagrangian (4.6) to make the second set of four SUSY transformations \( \delta \) manifest

\[
\mathcal{L} = -\frac{1}{4g^2} v_{mn} v^{mn} - \bar{\psi} \sigma^m \partial_m \psi + \frac{1}{2g^2} (-X_3)^2 + \frac{1}{2g^2} |X_1 - iX_2|^2
\]
\[
-\frac{1}{2g^2} |\partial_m \phi|^2 - \frac{1}{g^2} \bar{\lambda} \sigma^m \partial_m \lambda + c(-X_3) + b(-X_1 - iX_2) + b^*(-X_1 + iX_2)
\]
\[
+ \sum_{a=1}^n \left[ (-\bar{F}_a^* - q_a^* (\phi - m_a))(-\bar{F}_a^* - (\phi^* - m_a^*) q_a) - |\mathcal{D}_m \bar{q}_a|^2 \right]
\]
\[
- i \bar{\psi}_{qa} \sigma^m \mathcal{D}_m \psi_{qa} + \sqrt{2} \bar{q}_a \psi_{qa} \psi_{qa} + (-X_3) |\bar{q}_a|^2
\]
\[
+ \sum_{a=1}^n \left[ (F_a^* - \bar{q}_a (\phi^* - m_a^*)) (F_a^* - (\phi - m_a) \bar{q}_a^*) - |\mathcal{D}_m q_a|^2 \right]
\]
\[
- i \psi_{qa} \sigma^m \mathcal{D}_m \psi_{qa} + \sqrt{2} (\bar{q}_a \psi_{qa} \bar{q}_a + \psi_{qa} \psi_{qa}) (-X_3) |q_a|^2
\]
\[
+ \sum_{a=1}^n \left[ (\phi - m_a) (q_a^* (-\bar{F}_a^* - (\phi^* - m_a^*) q_a) + (F_a^* - \bar{q}_a (\phi^* - m_a^*) q_a^*) \right]
\[
+ (X_1 - i X_2) \eta_a^* \bar{\eta}_a - \psi_{\bar{q}_a} (\phi - m_a) \psi_{q_a} - i \sqrt{2} \psi_{\bar{q}_a} \lambda \bar{\eta}_a + i \sqrt{2} \psi_{q_a} \lambda \eta_a \\
+ \sum_{a=1}^{n} \left[ (\phi^* - m_a) \{ q_a (F_\alpha - (\phi - m_a) \bar{q}_a^\dagger) + (-\bar{F}_\alpha - q_a (\phi - m_a)) q_\alpha \} \\
+ (X_1 + i X_2) \bar{q}_a q_a - \bar{\psi}_{q_a} (\phi^* - m_a) \psi_{\bar{q}_a} + i \sqrt{2} \bar{\psi}_{\bar{q}_a} \lambda \bar{\eta}_a - i \sqrt{2} \bar{\psi}_{q_a} \lambda \eta_a \right].
\] (4.41)

We can finally assemble the above component form of the Lagrangian to an \( N = 1 \) superfield formalism using the second set of Grassmann number \( \theta_\pm \) in Eqs.(4.37)–(4.40), in contrast to those superfields with \( \theta_\pm \) in Eqs.(4.2)–(4.5):

\[
\mathcal{L} = \frac{1}{4 g^2} \left( W^\alpha W^\dagger_\alpha \right)_{\theta_2^2} + \frac{1}{2 g^2} \Phi_\dagger_\alpha \Phi_\alpha \left| \theta_2^2 \right| + \sum_{a=1}^{n} \left( Q_{-a} e^{2V-} Q_{-a} + \tilde{Q}_{-a} e^{-2V-} \tilde{Q}_{-a} \right) \left| \theta_2^2 \right| + 2 c V_{-} \left| \theta_2^2 \right| + \left( \sum_{a=1}^{n} \left( \Phi_{-} - m_a \right) Q_{-a} \tilde{Q}_{-a} \right) \left| \theta_2^2 \right| + b^* \Phi_{-} \left| \theta_2^2 \right| + \text{h.c.},
\] (4.42)

Let us note that the Lagrangian is not invariant under the automorphism \( SU(2)_R \) if FI parameters are present. Obviously the Lagrangian is invariant under the eight SUSY transformations when it is invariant under the first set of four SUSY and the discrete \( SU(2)_R \) transformation \( \mathfrak{M}_0 \). However, it is important to realize that the Lagrangian represented by (4.6) and (4.42) is invariant under all the eight SUSY transformations (4.7)–(4.18) and (4.25)–(4.33) irrespective of the values of the FI parameters \( c \) and \( b \). This is because the difference of the Lagrangian transformed by \( \mathfrak{M}_0 \) and the original one is given by the FI terms which are transformed to total derivative by SUSY transformations \( \delta_{\xi_+} \) and \( \delta_{\xi_-} \). Therefore the action is invariant as usual for the SUSY theories.

Theories with eight SUSY like our model have been shown to possess three complex structures \cite{48,26}. Our formulation in terms of two sets of Grassmann numbers \( \theta_\pm \) does not make this property manifest. In fact our choice of \( c = 0, b \in \mathbb{R} \) breaks \( SU(2)_R \) symmetry and particular complex structure has been selected. However, this particular choice of complex structure will turn out to be useful for the analysis of our model and solution. For instance, we will show that we can choose one of the two conserved SUSY directions from \( \theta_+ \), and the other from \( \theta_- \).

Since all the eight SUSY transformations are clarified, we are now in a position to determine precisely how many SUSY charges out of these eight are conserved by our solution (3.15)–(3.17) of the domain wall junction. We have already found in Eq.(2.34) that one out of four SUSY \( \delta_{\xi_-} \) is conserved. We need to examine another four SUSY \( \delta_{\xi_-} \). Let us first consider fermions \( \psi_{q_a} \) and \( \psi_{\bar{q}_a} \) in the \( N = 1 \) chiral scalar superfields \( \tilde{Q}_{-a} \) in Eq.(4.39) and \( Q_{-a} \) in Eq.(4.40) for the hypermultiplets:

\[
\delta_{\xi_-} \psi_{q_a} = i \sqrt{2} s^m \partial_q \bar{q}^\dagger \xi_- + \sqrt{2} (\phi^* - m_a^*) q_a \xi_- \\
= i 2 \sqrt{2} \left( \sigma^+ \partial_q \bar{q}^\dagger \xi_- + \sigma^- \partial_{\bar{q}}^\dagger \xi_- - i \frac{1}{2} (\phi^* - m_a^*) q_a \xi_- \right) \\
+ i 2 \sqrt{2} (\sigma^+ \xi_- + i \xi_-) \partial_q \bar{q}^\dagger_a + i 2 \sqrt{2} \sigma^- \xi_- \partial_{\bar{q}}^\dagger_a,
\] (4.43)
\[
\delta_{\xi_-} (\psi_{\bar{a}}) = i \sqrt{2} \sigma^m \partial_m \bar{q}_a \xi_- + \sqrt{2} (\phi^* - m_a^*) \bar{q}_a \xi_- \\
= i 2 \sqrt{2} (\sigma^+ \partial_q a \xi_- + \sigma^- \partial_{\bar{q}} a \xi_- - i \frac{1}{2} (\phi^* - m_a^*) \bar{q}_a \xi_-) \\
= i 2 \sqrt{2} \sigma^- \partial_q a \xi_- + i 2 \sqrt{2} (\sigma^+ \bar{q}_a + i \xi_-) \partial_{\bar{q}} a. 
\] (4.44)

Therefore the conserved SUSY direction is given by

\[
\sigma^+ \bar{\xi}_- = -i \xi_- \quad \text{and} \quad \sigma^- \bar{\xi}_- = 0. 
\] (4.45)

We have to study the SUSY transformation of the remaining fermions \( \lambda \) and \( \bar{\psi} \) in the \( \mathcal{N} = 1 \) chiral scalar superfield \( \Phi_- \) and \( V_- \) using the BPS equations (2.29) and (2.38) with \( \Omega = -1 \) and algebraic equations of motion for auxiliary fields (2.39). Moreover, we note that \( X_2 = X_3 = 0 \) in our solution. Then we find the same SUSY directions are conserved for \( \lambda \)

\[
\delta_{\xi_-} (i \sqrt{2} \lambda) = i \sqrt{2} \sigma^m \partial_m \phi \xi_- + \sqrt{2} (-X_1 + i X_2) \xi_- \\
= i 2 \sqrt{2} (\sigma^+ \partial_q a + \sigma^- \partial_{\bar{q}} a) \phi \xi_- + \sqrt{2} (-X_1 + i X_2) \xi_- \\
= i 2 \sqrt{2} (\sigma^+ \xi_- + i \xi_-) \partial_{\bar{q}} \phi + i 2 \sqrt{2} \sigma^- \xi_- \partial_{\bar{q}} \phi = 0, 
\] (4.46)

with the same infinitesimal SUSY transformation parameters (4.45). The right-hand side of \( \xi_- \) transformation of \( \psi \) in Eq. (4.26) involve only \( v_m \) and \( X_3 \) which vanish in our solution. Therefore they conserve all the SUSY transformations trivially.

Summarizing our results for all the eight SUSY, we see that there are two conserved directions in the Grassmann parameter

\[
\sigma^+ \xi_+ = i \xi_+ \quad \text{and} \quad \sigma^- \xi_+ = 0, 
\] (4.47)
\[
\sigma^+ \xi_- = -i \xi_- \quad \text{and} \quad \sigma^- \xi_- = 0. 
\] (4.48)

Namely we have determined the two conserved directions

\[
\xi_{+1} = i (\xi_{+1})^*, \quad \xi_{+2} = 0, \quad \xi_{-1} = -i (\xi_{-1})^*, \quad \xi_{-2} = 0 
\] (4.49)

where we set

\[
\xi_+ = \begin{pmatrix} \xi_{+1} \\ \xi_{+2} \end{pmatrix}, \quad \bar{\xi}_+ = \begin{pmatrix} \bar{\xi}_{+1} \\ \bar{\xi}_{+2} \end{pmatrix} = \begin{pmatrix} (\xi_{+2})^* \\ -(\xi_{+1})^* \end{pmatrix}, 
\] (4.50)
\[
\xi_- = \begin{pmatrix} \xi_{-1} \\ \xi_{-2} \end{pmatrix}, \quad \bar{\xi}_- = \begin{pmatrix} \bar{\xi}_{-1} \\ \bar{\xi}_{-2} \end{pmatrix} = \begin{pmatrix} (\xi_{-2})^* \\ -(\xi_{-1})^* \end{pmatrix}. 
\] (4.51)

We have now established that our domain wall solution preserves two out of eight SUSY.

5 Dimensional reduction from five dimensions

The \( \mathcal{N} = 1 \) superfield formalism is of course most useful to study \( \mathcal{N} = 1 \) SUSY theories in four dimensions. It can also be used to describe the \( \mathcal{N} = 2 \) SUSY theories in four dimensions.
making only four out of eight SUSY manifest. In order to describe \( \mathcal{N} = 2 \) SUSY theories in five dimensions, however, we need to sacrifice the five-dimensional Lorentz invariance \([5],[11]–[14]\).

In terms of components, we can always express the \( \mathcal{N} = 2 \) SUSY theories with all the necessary auxiliary fields to close the algebra off-shell in the Wess-Zumino gauge.

The highest dimension allowed by the \( \mathcal{N} = 2 \) SUSY is six. The hypermultiplet in six dimensions cannot have masses. The five-dimensional theories can be obtained by a dimensional reduction from six dimensions. If we perform a nontrivial dimensional reduction, allowing the momenta in the sixth dimension, we obtain a massive hypermultiplet \([15]\). Therefore the hypermultiplets in five dimensions can have only real mass parameters. The \( \mathcal{N} = 2 \) SUSY theories in four dimensions can be obtained by dimensionally reducing the five-dimensional theories. The real scalar field \( \Sigma \) in five dimensions comes originally from the sixth component of gauge field and the combination \( \Sigma + iv_5 \) becomes a complex scalar in four dimensions, when we consider the reduction to four dimensions. To obtain the four-dimensional theory with complex mass parameters, we should perform a nontrivial dimensional reduction in the \( x^5 \) direction. In this process, the mass terms arise as momenta in the fifth dimension. Conversely we can recover the five-dimensional theory by restoring the fifth dimension from the imaginary part of the complex mass parameter \( m_a \equiv m_{aR} + im_{aI} \) as

\[
\partial_5 q_a = -im_{aI}q_a, \quad \partial_5 \psi_q = -im_{aI}\psi_q. \tag{5.1}
\]

The imaginary part of the complex scalar field \( \phi \) can also be identified as the fifth component of the vector field \( v_5 \) as

\[
\phi = \Sigma + iv_5. \tag{5.2}
\]

We can recover the covariant derivative along the fifth direction as

\[
i(v_5 - m_{aI})q_a = (\partial_5 + iv_5)q_a = \mathcal{D}_5 q_a \tag{5.3}
\]

Therefore the mass terms associated with the quark fields are reduced to covariant derivatives in the fifth dimension

\[
(\phi - m_a)q_a = [\Sigma - m_{aR}] + i(v_5 - m_{aI})]q_a = [\Sigma - m_{aR}] + \mathcal{D}_5 q_a. \tag{5.4}
\]

In the spirit of the \( \mathcal{N} = 1 \) superfield formalisms, we can express the \( \mathcal{N} = 2 \) vector and hypermultiplets in five dimensions by two kinds of superfields similarly to our results \([4.2]–[4.5]\), and \([4.37]–[4.40]\) for four-dimensional \( \mathcal{N} = 2 \) theories. The \( \theta_+ \) superfields are given by

\[
V_+(x, \theta_+, \bar{\theta}_+) = -\theta_+ \sigma^m \bar{\theta}_+ v_m + i\theta_+^2 \bar{\theta}_+ \lambda - i\bar{\theta}_+^2 \theta_+ \lambda + \frac{1}{2} \theta_+^2 \bar{\theta}_+^2 (X_3 - \partial_5 \Sigma), \tag{5.5}
\]

\[
\Phi_+(y, \theta_+) = (\Sigma + iv_5) + \sqrt{2} \theta_+ (-i\sqrt{2}\psi) + \theta_+^2 (X_1 + iX_2), \tag{5.6}
\]

\[
Q_+(a, y, \theta_+) = q_a + \sqrt{2}\theta_+ \psi_{q_a} + \theta_+^2 (F_a' + \mathcal{D}_5 q_a^* - (\Sigma - m_{aR}) q_a^*), \tag{5.7}
\]

\[
\tilde{Q}_+(a, y, \theta_+) = \tilde{q}_a + \sqrt{2}\theta_+ \psi_{\tilde{q}_a} + \theta_+^2 (-\tilde{F}_a' - \mathcal{D}_5 q_a^* - q_a^*(\Sigma - m_{aR})). \tag{5.8}
\]
The $\theta_-$ superfields are given by
\[
V_-(x, \theta_-, \bar{\theta}_-) = -\theta_- \sigma^m \bar{\theta}_- v_m + i \theta_+^2 \bar{\theta}_- \psi - i \bar{\theta}_- \theta_+ \psi + \frac{1}{2} \theta_+^2 \bar{\theta}_-^2 (-X_3 - \bar{\sigma}_3 \Sigma), \tag{5.9}
\]
\[
\Phi_-(y, \theta_-) = (\Sigma + i \tilde{\psi}_3) + \sqrt{2} \theta_- (\sqrt{2} i \lambda) + \theta_-^2 (-X_1 + i X_2), \tag{5.10}
\]
\[
Q_{-a}(y, \theta_-) = \bar{q}_a^* + \sqrt{2} \theta_- \psi^a + \theta_-^2 (\tilde{F}_a^s - D_5 q_a + (\Sigma - m_{aR}) q_a), \tag{5.11}
\]
\[
\bar{Q}_{-a}(y, \theta_-) = q_a^* + \sqrt{2} \theta_- (-\psi^a) + \theta_-^2 (-F_a^s + D_5 q_a + \bar{q} (\Sigma - m_{aR})). \tag{5.12}
\]

Please note that all the fields depend on the coordinate $x^5$ in fifth dimensions, in spite of almost the same appearance as the four-dimensional superfields.

Since the mass term in five dimensions can be obtained as a nontrivial dimensional reduction from six dimensions, we can have only real mass parameter in the Lagrangian $m_{aR}$ \cite{42}, \cite{45}. Therefore we find the Lagrangian in terms of the $\theta_+ \mathcal{N} = 1$ superfields in Eqs. (5.5)– (5.8) as
\[
\mathcal{L} = \frac{1}{4g^2} \left( W^a_\alpha \bar{W}^\alpha_\alpha \right)_{\theta_+^2} + \frac{1}{g^2} \left( \partial_5 V - \frac{\Phi_+^\dagger + \Phi_+}{2} \right)_{\theta_+^2} + \sum_{a=1}^n \left( Q_{+a} e^{2V_+} Q_{+a} + \bar{Q}_{+a} e^{-2V_+} \bar{Q}_{+a} \right) \left( \partial_5 \Phi_+ - m_{aR} \Phi_+ \right)_{\theta_+^2} + \left( \sum_{a=1}^n \bar{Q}_{+a} \left( \partial_5 + \Phi_+ - m_{aR} \right) Q_{+a} \right)_{\theta_+^2} + \text{h.c.}, \tag{5.13}
\]
The same Lagrangian is given in terms of the $\theta_- \mathcal{N} = 1$ superfields in Eqs. (5.9)– (5.12) as
\[
\mathcal{L} = \frac{1}{4g^2} \left( W^a_\alpha \bar{W}^\alpha_\alpha \right)_{\theta_-^2} + \frac{1}{g^2} \left( \partial_5 V - \frac{\Phi_-^\dagger + \Phi_-}{2} \right)_{\theta_-^2} + \sum_{a=1}^n \left( Q_{+a} e^{2V_-} Q_{+a} + \bar{Q}_{+a} e^{-2V_-} \bar{Q}_{+a} \right) \left( \partial_5 \Phi_- + m_{aR} \Phi_- \right)_{\theta_-^2} + \left( \sum_{a=1}^n \bar{Q}_{+a} \left( \partial_5 + \Phi_- + m_{aR} \right) Q_{+a} \right)_{\theta_-^2} + \text{h.c.}, \tag{5.14}
\]

In terms of components, we can express the Lagrangian more symmetrically with respect to the $SU(2)_R$ symmetry
\[
\mathcal{L} = \mathcal{L}_{\text{boson}} + \mathcal{L}_{\text{fermion}} \tag{5.15}
\]
\[
\mathcal{L}_{\text{boson}} = -\frac{1}{4g^2} F_{MN} F_{MN} - \frac{1}{2g^2} (\partial_5 \Sigma)^2 + \frac{1}{2g^2} \left\{ (X_1)^2 + (X_2)^2 + (X_3)^2 \right\} + c(-X_3) + b(-X_1 - iX_2) + b^*(X_1 - iX_2) \\
+ \sum_{a=1}^n \left[ |D_M q_a|^2 - |D_M \bar{q}_a|^2 + |F^v|^2 + |\tilde{F}^v|^2 - (\Sigma - m_{aR})^2 (|q_a|^2 + |\bar{q}_a|^2) \right] \\
+ X_3 (|q_a|^2 - |\bar{q}_a|^2) + (X_1 + iX_2) \bar{q}_a q_a + (X_1 - iX_2) q_a^* \bar{q}_a \tag{5.16}
\]
The scalars and auxiliary fields in the hypermultiplet transform as doublet under the \( \psi^a \) four-component spinor \( SU(2) \) vector multiplet can be assembled into a four-component matrix as

\[
\begin{bmatrix}
0 & \sigma^1 \\
\sigma^2 & 0
\end{bmatrix}
\]

where \( \sigma^m = (1, \vec{\sigma}) \) and \( \sigma^m = (1, -\vec{\sigma}) \). To achieve the \( \mathcal{N} = 2 \) SUSY off-shell formalism, it is convenient to make \( SU(2)_R \) manifest. In the case of vector multiplet, the spinors in five dimensions are most conveniently organized in terms of the symplectic \( (SU(2)) \) Majorana spinors \( \lambda^i, \ i = 1, 2 \), transforming as doublets under the \( SU(2)_R \) symmetry. The \( SU(2) \) Majorana spinor is defined by

\[
\lambda^i = \epsilon^{ij} C \bar{\lambda}^j,
\]

where the charge conjugation matrix \( C \) in five dimensions satisfies \( C \gamma^M C^{-1} = (\gamma^M)^T, \ C^T = -C, \) and \( CC^\dagger = 1 \). An explicit form may be given by \( C = \text{diag}(i\sigma^2, i\sigma^2) \). The spinors in the \( \mathcal{N} = 2 \) vector multiplet can be assembled into a four-component \( SU(2) \) Majorana spinor \( \lambda^i \) as

\[
\lambda^1 = \left( \begin{array}{c} \lambda_\alpha \\ \bar{\psi}_a \end{array} \right), \quad \lambda^2 = \left( \begin{array}{c} \bar{\psi}_a \\ -\lambda_\alpha \end{array} \right),
\]

\[
\bar{\lambda}_1 = \left( \begin{array}{c} \psi^\alpha \\ \bar{\lambda}_a \end{array} \right), \quad \bar{\lambda}_2 = \left( \begin{array}{c} -\lambda^\alpha \\ \bar{\psi}_a \end{array} \right).
\]

The spinor in the hypermultiplets are singlets under the \( SU(2)_R \) symmetry and is assembled into a four-component spinor \( \psi^a \) for each flavor

\[
\psi = \left( \begin{array}{c} \psi^a \\ \bar{\psi}^a \end{array} \right), \quad \bar{\psi} = \left( \begin{array}{c} \bar{\psi}^a \\ -\psi^a \end{array} \right).
\]

The scalars and auxiliary fields in the hypermultiplet transform as doublet under the \( SU(2)_R \) as given in Eqs. (4.21) and (4.24).

The above five-dimensional Lagrangian \((5.15)-(5.17)\) makes it clear that we can have only real mass parameters for hypermultiplets, which is obtained as a momentum in one extra dimension.
(the sixth dimension) through a nontrivial (Scherk-Schwarz) dimensional reduction from six dimensions. On the other hand, we need to have at least three discrete vacua in complex field plane. This situation can be realized in our model through the complex masses of hypermultiplets. This is the reason why we cannot generalize our junction solution to five or six-dimensions.

Acknowledgments

One of the authors (N.S.) acknowledges useful discussion of $\mathcal{N} = 2$ SUSY theories with Masato Arai, Masashi Naganuma, Muneto Nitta, and Keisuke Ohashi. This work is supported in part by Grant-in-Aid for Scientific Research from the Japan Ministry of Education, Science and Culture 13640269.

References

[1] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. B429 (1998) 263 [hep-ph/9803315]; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. B436 (1998) 257 [hep-ph/9804398].

[2] L. Randall and R. Sundrum, Phys. Rev. Lett. 83 (1999) 3370 [hep-ph/9905221]; Phys. Rev. Lett. 83 (1999) 4690 [hep-th/9906064].

[3] S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150; N. Sakai, Z. f. Phys. C11 (1981) 153; E. Witten, Nucl. Phys. B188 (1981) 513; S. Dimopoulos, S. Raby, and F. Wilczek, Phys. Rev. D24 (1981) 1681.

[4] L. Randall and R. Sundrum, Nucl. Phys. B557 (1999) 79 [hep-ph/9810155].

[5] E.A. Mirabelli and M.E. Peskin, Phys. Rev. D58 (1998) 065002 [hep-th/9712214].

[6] D.E. Kaplan, G.D. Kribs and M. Schmaltz, Phys. Rev. D62 (2000) 035010 [hep-ph/9911293]; Z. Chacko, M.A. Luty, A.E. Nelson and E. Ponton, JHEP 0001 (2000) 003 [hep-ph/9911323].

[7] T. Kobayashi and K. Yoshioka, Phys. Rev. Lett. 85 (2000) 5527 [hep-ph/0008069]; Z. Chacko and M.A. Luty, JHEP 0105 (2001) 067 [hep-ph/0008103].

[8] N. Arkani-Hamed, L. Hall, D. Smith and N. Weiner, Phys. Rev. D63 (2000) 056003 [hep-ph/9911421].

[9] D.E. Kaplan and T.M.P. Tait, JHEP 0006 (2000) 020 [hep-ph/0004200].

[10] N. Maru, N. Sakai, Y. Sakamura, and R. Sugisaka, Phys. Lett. B496 (2000) 98, [hep-th/0009023].

[11] N. Maru, N. Sakai, Y. Sakamura, and R. Sugisaka, Nucl. Phys. B616 (2001) 47, [hep-th/0107204].
[12] E. Witten and D. Olive, Phys. Lett. B78 (1978) 97.

[13] M. Cvetic, F. Quevedo and S.J. Rey, Phys. Rev. Lett. 67 (1991) 1836; M. Cvetic, S. Griffies and S.J. Rey, Nucl. Phys. B381 (1992) 301 [hep-th/9201007]; M. Cvetic, and H.H. Soleng, Phys. Rep. B282 (1997) 159, [hep-th/9604090].

[14] G. Dvali and M. Shifman, Phys. Lett. B396 (1997) 64 [hep-th/9612128]; Nucl. Phys. B504 (1997) 127 [hep-th/9611213]; A. Kovner, M. Shifman, and A. Smilga, Phys. Rev. D56 (1997) 7978 [hep-th/9706089]; A. Smilga and A. Veselov, Phys. Rev. Lett. 79 (1997) 4529 [hep-th/9706217]; B. Chibisov and M. Shifman, Phys. Rev. D56 (1997) 7990; V. Kaplunovsky, J. Sonnenschein, and S. Yankielowicz, Nucl. Phys. B552 (1999) 209 [hep-th/9811195]; B. de Carlos and J. M. Moreno, Phys. Rev. Lett. 83 (1999) 2120 [hep-th/9905165]; G. Dvali, G. Gabadadze, and Z. Kakushadze, Nucl. Phys. B493 (1997) 148, [hep-th/9901032]; J.D. Edelstein, M.L. Trobo, F.A. Brito and D. Bazeia, Phys. Rev. D57 (1998) 7561 [hep-th/9707016]; M. Naganuma and M. Nitta, Prog. Theor. Phys.105(2001)501 [hep-th/0007184].

[15] M. Eto, N. Maru, T. Sakata and N. Sakai, Phys. Lett. B553 (2003) 87, [hep-th/0208127].

[16] M. Arai, S. Fujita, M. Naganuma and N. Sakai, Phys. Lett. B556 (2003) 192, [hep-th/0212175].

[17] E. R. C. Abraham and P. K. Townsend, Nucl. Phys. B351 (1991) 313.

[18] G. Gibbons and P.K. Townsend, Phys. Rev. Lett. 83 (1999) 1727 [hep-th/9905196]; P.M. Safin, Phys. Rev. Lett. 83 (1999) 4249-4252 [hep-th/9907066].

[19] S.M. Carroll, S. Hellerman and M. Trodden, Phys. Rev. D61 (2000) 065001 [hep-th/9905217].

[20] H. Oda, K. Ito, M. Naganuma and N. Sakai, Phys. Lett. 471B (1999) 140 [hep-th/9910095].

[21] K. Ito, M. Naganuma, H. Oda and N. Sakai, Nucl. Phys. B586 (2000) 231 [hep-th/0004188]; Nucl. Phys. Proc. Suppl. 101 (2001) 304 [hep-th/0012182].

[22] A. Gorsky and M. Shifman, Phys. Rev. D61 (2000) 085001 [hep-th/9909015]; Nucl. Phys. B586 (2000) 231 [hep-th/0004188]; D. Bazeia, F.A. Brito, Phys. Rev. D61 (2000) 105019, [hep-th/9912015]; D. Binosi and T. ter Veldhuis, Phys. Lett. B476 (2000) 124 [hep-th/9912081]; S. Nam and K. Olsen, JHEP 0008 (2000) 001 [hep-th/0002176].

[23] S.M. Carroll, S. Hellerman and M. Trodden, Phys. Rev. D62 (2000) 044049, [hep-th/9911083]; A. Chamblin, C. Csaki, J. Erlich, T.J. Hollowood, Phys. Rev. D62 (2000) 044012, [hep-th/0002076]; T. Nihei, Phys. Rev. D62 (2000) 124017, [hep-th/0005014].

[24] M. Naganuma, M. Nitta and N. Sakai, Phys. Rev. D65 (2002) 045016, [hep-th/0108179]; M. Naganuma, M. Nitta and N. Sakai, “BPS Walls and Junctions in $\mathcal{N} = 1$ SUSY Nonlinear Sigma Models”, [hep-th/0210205].
[25] M. Sakamoto and S. Tanimura, Phys. Rev. D65 (2002) 065004, hep-th/0108208.

[26] J.P. Gauntlett, D. Tong and P.K. Townsend, Phys. Rev. D63 (2001) 085001 [hep-th/0007124]; J.P. Gauntlett, R. Portugues, D. Tong and P.K. Townsend, Phys. Rev. D63 (2001) 085002 [hep-th/0008221]; J.P. Gauntlett, D. Tong and P.K. Townsend, Phys. Rev. D64 (2001) 25010 [hep-th/0107228].

[27] M. Arai, M. Naganuma, M. Nitta and N. Sakai, Nucl. Phys. B652 (2003) 35, hep-th/0211103; “BPS wall in $\mathcal{N} = 2$ SUSY nonlinear sigma model with Eguchi-Hanson manifold”, hep-th/0302025; to appear in Graden of Quanta in honor of Hiroshi Ezawa, (World Scientific Pub.).

[28] D. Tong, Phys. Rev. D66 (2002) 025013, hep-th/0202012.

[29] M. Shifman and A. Yung, “Domain Walls and Flux Tubes in $\mathcal{N} = 2$ SQCD: D-Brane Prototypes,” hep-th/0212293.

[30] D. Tong, “Mirror Mirror on the Wall, (On Two-Dimensional Black Holes and Liouville Theory”, hep-th/0303151.

[31] E.R.C. Abraham and P.K. Townsend, Phys. Lett. B291 (1992) 85; ibid. B295 (1992) 225.

[32] M. Naganuma, M. Nitta, and N. Sakai, Grav. Cosmol. 8 (2002) 129, hep-th/0108133.

[33] N. Seiberg and E. Witten, Nucl. Phys. B426 (1994) 19; Nucl. Phys. B431 (1994) 484.

[34] P. Fayet, Nucl. Phys. B113 (1976) 135.

[35] U. Lindström and M. Roček, Nucl. Phys. B222 (1983) 285.

[36] A. Galperin, E. Ivanov, S. Kalitzin, V. Ogievetsky and E. Sokatchev, Class. Quantum Grav. 1 (1984) 469; A. Galperin, E. Ivanov, V. Ogievetsky and E. Sokatchev, Class. Quantum Grav. 2 (1985) 601; 617.

[37] A. Galperin, E. Ivanov, V. Ogievetsky and E. Sokatchev, “Harmonic Superspace”, Cambridge University Press, Cambridge (2001).

[38] B. M. Zupnik, Theor. Math. Phys. 130 (2002) 213 hep-th/0107012.

[39] A. Galperin, E. Ivanov, V. Ogievetsky and P. K. Townsend, Class. Quantum Grav. 3 (1986) 625.

[40] B.M. Zupnik, “Static BPS-conditions in N=2 harmonic superspace”, hep-th/0109113.

[41] N. Arkani-Hamed, L. Hall, D. Smith and N. Weiner, Phys. Rev. D63 (2001) 056003, hep-th/9911421; N. Arkani-Hamed, T. Gregoire and J. Wacker, JHEP 0203 (2002) 055, hep-th/0101233.

[42] A. Hebecker, Nucl. Phys. B632 (2002) 101, hep-ph/0112230.
[43] N. Marcus, A. Sagnotti and W. Siegel, *Nucl. Phys.* **B224** (1983) 159.

[44] W.D. Linch III, M.A. Luty and J. Phillips, “Five dimensional supergravity in $\mathcal{N} = 1$ superspace”, [hep-ph/0209060].

[45] G. Sierra and P.K. Townsend, *Nucl. Phys.* **B233** (1984) 289.

[46] J. Scherk and J.H. Schwarz, *Phys. Lett.* **B82** (1979) 60.

[47] J. Wess and J. Bagger, “Supersymmetry and Supergravity”, (1991), Princeton University Press.

[48] L. Alvarez-Gaumé and D.Z. Freedman, *Comm. Math. Phys.* **91** (1983) 87.

[49] R. Portugues and P.K. Townsend, *Phys.Lett.* **B530** (2002) 227, [hep-th/0112077]; S. Cecotti, and C. Vafa, *Commun.Math.Phys.* **158** (1993) 569, [hep-th/9211097].

[50] A. Ritz, M. Shifman, A. Vainshtein, M. Voloshin, *Phys.Rev.* **D63** (2001) 065018, [hep-th/0006028]; Philip C. Argyres, and K. Narayan, *Int. J. Mod. Phys.* **A16S1C** (2001) 962, [hep-th/0101139]; A. Ritz, A. Vainshtein, *Nucl.Phys.* **B617** (2001) 43, [hep-th/0102121].

[51] N. Arkani-Hamed, S. Dimopoulos, G. Dvali, and N. Kaloper, *Phys.Rev.Lett.* **84** (2000) 586, [hep-th/9907209].