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Effect of bare mass on the Hosotani mechanism

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Abstract

It is pointed out that the existence of bare mass terms for matter fields changes gauge symmetry patterns through the Hosotani mechanism. As a demonstration, we study an $SU(2)$ gauge model with massive adjoint fermions defined on $M^4 \otimes S^1$. It turns out that the vacuum structure changes at certain critical values of mL , where m (L) stands for the bare mass (the circumference of S^1). The gauge symmetry breaking patterns are different from models with massless adjoint fermions. We also consider a supersymmetric $SU(2)$ gauge model with adjoint hypermultiplets, in which the supersymmetry is broken by bare mass terms for the gaugino and squark fields instead of the Scherk–Schwarz mechanism.

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1. Introduction

The Hosotani mechanism is one of the most important dynamical phenomena when one considers physics with compactified extra dimensions [1,2]. The component gauge field for the compactified direction becomes dynamical degrees of freedom and can develop vacuum expectation values, which are related to phases of the Wilson line integrals along the compactified direction. Wilson line is an order parameter for the gauge symmetry breaking. Quantum corrections in the extra dimension are crucial for the mechanism and gauge symmetry can be broken dynamically, reflecting the topology of the extra dimension.

It is important to note that the mechanism is essentially governed by infrared physics. In order to study vacuum structures of the theory, one usually studies the effective potential for the Wilson line phases. The Wilson line is a global quantity, so that the effective potential is free from ultraviolet effects because they are local. This suggests that massive particles, at first glance, do not contribute to the effective potential for the phases. Thus, the gauge symmetry breaking patterns seem not to be affected by the massive particle.

In this Letter we demonstrate that massive particles, on the contrary, can affect the gauge symmetry breaking patterns through the Hosotani mechanism. In fact, we will show that in gauge models with massive fermions, the gauge symmetry breaking patterns change at certain critical values of the bare mass for the fermion. This is never observed in the case of massless fermions.

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We also study supersymmetric gauge models, in which the supersymmetry is broken explicitly by bare mass terms, instead of imposing the twisted boundary condition for the S^1 direction by Scherk and Schwarz [3]. We introduce a bare mass term for the gaugino(squark) field in a vector(hyper)multiplet. We will show that, depending on the relative magnitude between the two mass parameters, the gauge symmetry breaking patterns become different from those by the model with Scherk–Schwarz mechanism of supersymmetry breaking.

Since the full analysis is beyond the scope of this Letter, we restrict our consideration to a simple $SU(2)$ (supersymmetric) gauge model with N_f massive adjoint fermions (hypermultiplets) defined on $M^4 \otimes S^1$, where M^4 (S^1) stands for four-dimensional Minkowski space–time (a circle with the circumference being L). We are interested in the effects of bare mass terms for matter fields on the gauge symmetry breaking patterns.

2. Gauge model with massive adjoint matter

Let us start with the nonsupersymmetric $SU(2)$ gauge model with N_f massive adjoint fermions defined on $M^4 \otimes S^1$. Following the standard procedure, the effective potential for the constant background gauge field $gL(A_y) = \text{diag}(\theta_1, \theta_2) = \text{diag}(\theta, -\theta)$ is given by [2,4]

$$V_{\text{eff}} = -(5-2) \frac{\Gamma(\frac{5}{2})}{\pi^{\frac{5}{2}}} \frac{1}{L^5} \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \frac{1}{n^5} \cos(n(\theta_i - \theta_j)) + (N_f 2^{[\frac{5}{2}]}) \frac{2}{(2\pi)^{\frac{5}{2}}} \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \left(\frac{m}{nL}\right)^{\frac{5}{2}} K_{\frac{5}{2}}(mLn) \cos(n(\theta_i - \theta_j)). \tag{1}$$

Here m stands for a gauge invariant bare mass for the adjoint fermion. The first line in Eq. (1) comes from the gauge and ghost fields and the second line is the contribution from the N_f massive adjoint fermions. The function $K_{\frac{5}{2}}(y)$ is the modified Bessel function and is expressed in terms of the elementary function,

$$K_{\frac{5}{2}}(y) = \left(\frac{\pi}{2y}\right)^{\frac{1}{2}} \left(1 + \frac{3}{y} + \frac{3}{y^2}\right) e^{-y}. \tag{2}$$

Let us note that the effective potential (1) becomes identical to the one with N_f massless adjoint fermions when we take the limit,

$$\lim_{m \rightarrow 0} m^{\frac{5}{2}} K_{\frac{5}{2}}(mLn) = \frac{2^{\frac{3}{2}} \Gamma(\frac{5}{2})}{(nL)^{\frac{5}{2}}}. \tag{3}$$

By defining $z \equiv mL$ and noting that $\Gamma(5/2) = 3\sqrt{\pi}/4$. The effective potential is recast as

$$V_{\text{eff}} = \frac{3}{4\pi^2} \frac{1}{L^5} \sum_{n=1}^{\infty} \frac{1}{n^5} \left(-3 + 4N_f \left(1 + zn + \frac{(zn)^2}{3}\right) e^{-zn}\right) \times 2(1 + \cos(2n\theta)), \tag{4}$$

where we have used Eq. (2).

In order to see how the massive fermion affects the vacuum structure of the model, let us first consider asymptotic behaviors of the effective potential with respect to z . If z is large enough, the fermion contribution to the effective potential is suppressed due to the Boltzmann like factor e^{-zn} in Eq. (4). This is consistent with the observation stated in the introduction that the Hosotani mechanism is governed by infrared physics. Then, the dominant contribution to the potential comes from the gauge sector in the model. Therefore, the vacuum configuration is given, in this limit, by $\theta = 0(\text{mod } \pi)$ [2]. The $SU(2)$ gauge symmetry is not broken in the limit.¹ On

¹ The two configurations $\theta = 0$ and $\theta = \pi$ are physically equivalent and is related by Z_2 symmetry.

the other hand, if we take the massless limit, $z \rightarrow 0$, the effective potential, as we have mentioned above, becomes identical to the one for the case of N_f massless adjoint fermion. It has been known that the vacuum configuration for this case is given by $\theta = \pi/2$ [5]. The $SU(2)$ gauge symmetry is spontaneously broken down to $U(1)$ in this limit. The above observations strongly suggest that there must exist certain critical values of z , at which the gauge symmetry breaking patterns change.

In order to confirm the existence of the expected critical values of z , let us study the stability of the configuration $\theta = 0$ ($\pi/2$) with respect to z , which corresponds to the vacuum configuration in limit of $z \rightarrow \infty$ (0). We set $N_f = 1$ for simplicity. By simple numerical calculations, we find that the sign of the second derivative of the effective potential at $\theta = \pi/2$ (0) changes at $z \simeq 1.54135$ (1.13501) and becomes negative (positive) for larger values of it. Both configurations are stable for $1.13501 \lesssim z \lesssim 1.54135$. If we compare the potential energy for the two configurations in the narrow region, we find that $V_{\text{eff}}(\theta = \pi/2) < (>) V_{\text{eff}}(\theta = 0)$ for $z < (>) z_* \simeq 1.40087$.

These observations mean that the gauge symmetry breaking patterns change at z_* . As far as our numerical analyzes are concerned, there appears two degenerate minima at $z_* \simeq 1.40087$ and its finite height of potential barrier between the configuration $\theta = 0$ and $\theta = \pi/2$. If z becomes smaller than z_* , the minimum of the effective potential locates at $\theta = \pi/2$, while if z becomes larger than z_* , then, the vacuum configuration is given by $\theta = 0 \pmod{\pi}$. Hence, we conclude that²

$$\text{gauge symmetry breaking pattern} = \begin{cases} SU(2) \rightarrow U(1) & \text{for } z < z_*, \\ SU(2) \rightarrow SU(2) & \text{for } z > z_*. \end{cases} \quad (5)$$

We have also computed the critical values of z in case $N_f = 3, 6, 10$. They are given by $z_* \simeq 3.55154, 4.61892, 5.3521$, respectively. If we take $N_f = 100$, the critical value is about $z_* \simeq 8.378$.

Let us comment on the adjoint Higgs scalar, which is originally the component gauge field for the S^1 direction. The mass term for the scalar, which is massless at the tree level, is generated through the quantum correction in the extra dimensions [1]. The mass term is given by estimating the second derivative of the effective potential (4) at the absolute minimum. Our numerical analyzes tell us that the Higgs scalar is always massive for each gauge symmetry breaking pattern in Eq. (5).

3. Supersymmetric gauge model

In this section let us study the $\mathcal{N} = 1$ supersymmetric $SU(2)$ gauge model with N_f adjoint hypermultiplets defined on $M^4 \otimes S^1$. The bare mass terms, which are gauge invariant, are introduced in such a way that they break the supersymmetry. Here we are interested in how the gauge symmetry breaking patterns are modified by such the supersymmetry breaking terms and comparing results with those by the model with the Scherk–Schwarz mechanism of supersymmetry breaking.

In the model we have a five-dimensional vectormultiplet $\mathcal{V} = (A_{\hat{\mu}}, \Sigma, \lambda_D)$, where $A_{\hat{\mu}}$ is the five-dimensional gauge potential and λ_D (Σ) stands for a Dirac spinor³ (a real scalar). Here we call λ_D gaugino. We also have N_f adjoint hypermultiplets $\mathcal{H} = (\psi_D, \phi_i)$, where ψ_D is a Dirac spinor and $\phi_{i(=1,2)}$ is a complex scalar called squark. In case of the Scherk–Schwarz mechanism of supersymmetry breaking, the gaugino and squark masses are shifted by an unique nontrivial phase associated with the $SU(2)_R$ symmetry, so that they have different mass terms from their superpartners in four dimensions. The supersymmetry is broken by the unique phase. Here instead of resorting to the Scherk–Schwarz mechanism, we add the gauge invariant bare mass for the gaugino (m_g) and squarks (m_s), as one of the examples, to break the supersymmetry.

² The phase transition at z_* is the first order.

³ Let us note that n Dirac spinors are equivalent to $2n$ symplectic (pseudo)Majorana spinors. λ_D can be decomposed into two Majorana spinors λ, λ' in four dimensions.

Following the standard procedure to calculate the effective potential for the constant background gauge field $gL\langle A_y \rangle = \text{diag}(\theta_1, \theta_2) = \text{diag}(\theta, -\theta)$, we obtain that⁴

$$\begin{aligned}
 V_{\text{eff}} = & (-3 - 1) \frac{\Gamma(\frac{5}{2})}{\pi^{\frac{5}{2}}} \frac{1}{L^5} \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \frac{1}{n^5} \cos(n(\theta_i - \theta_j)) \\
 & + \frac{2}{(2\pi)^{\frac{5}{2}}} (2^{\lfloor \frac{5}{2} \rfloor}) \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \left(\frac{m_g}{nL}\right)^{\frac{5}{2}} K_{\frac{5}{2}}(z_g n) \cos(n(\theta_i - \theta_j)) \\
 & - \frac{2}{(2\pi)^{\frac{5}{2}}} (4N_f) \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \left(\frac{m_s}{nL}\right)^{\frac{5}{2}} K_{\frac{5}{2}}(z_s n) \cos(n(\theta_i - \theta_j)) \\
 & + \frac{\Gamma(\frac{5}{2})}{\pi^{\frac{5}{2}}} (2^{\lfloor \frac{5}{2} \rfloor} N_f) \frac{1}{L^5} \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \frac{1}{n^5} \cos(n(\theta_i - \theta_j)), \tag{6}
 \end{aligned}$$

where we have defined $z_g \equiv m_g L$, $z_s \equiv m_s L$. The first and second lines in Eq. (6) come from the vectormultiplet. The third and fourth lines are the contributions from the N_f adjoint hypermultiplets. Here we have assumed that the two complex scalars in the hypermultiplet have a common bare mass m_s . The bare mass $m_g(m_s)$ explicitly breaks the supersymmetry as it should. In fact, if we take the limit of $m_g(m_s) \rightarrow 0$ and utilizing (3), the effective potential (6) vanish and the original $\mathcal{N} = 1$ supersymmetry in five dimensions is restored.

The effective potential is recast, by using Eq. (2), as

$$V_{\text{eff}} = 4 \left(\frac{3}{4\pi^2} \right) \frac{1}{L^5} \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \frac{1}{n^5} (N_f F(n, z_s) - F(n, z_g)) \times 2(1 + \cos(2n\theta)), \tag{7}$$

where we have defined

$$F(n, z_i) \equiv 1 - \left(1 + nz_i + \frac{(nz_i)^2}{3} \right) e^{-nz_i}, \quad i = g, s. \tag{8}$$

The function $F(n, z_i)$ satisfies $0 \leq F(n, z_i) \leq 1$, where the first (second) equality holds when $z_i \rightarrow 0$ (∞).

Let us first study the case of $z_g = z_s (\equiv z_c)$. The effective potential (7) becomes

$$V_{\text{eff}} = 4 \left(\frac{3}{4\pi^2} \right) \frac{N_f - 1}{L^5} \sum_{i,j=1}^2 \sum_{n=1}^{\infty} \frac{1}{n^5} F(n, z_c) \times 2(1 + \cos(2n\theta)). \tag{9}$$

It is easy to see that the potential vanishes for $N_f = 1$. This is because one can still have $\mathcal{N} = 1$ supersymmetry by recombining the two massive fields into the same multiplet, so that there is one massless and one massive multiplet. And each multiplet is supersymmetric under the $\mathcal{N} = 1$ supersymmetry. As a result, the total action is invariant under the supersymmetry.

The nonvanishing potential is given for $N_f \geq 2$. The supersymmetry is broken by an unique parameter z_c in this case. Taking $0 \leq F(n, z) \leq 1$ into account, the minimum of the potential is always located at $\theta = \pi/2$, independent of the values of $z_c (\neq 0)$. Thus, the $SU(2)$ gauge symmetry is broken to $U(1)$. This is the same result as the one obtained by the Scherk–Schwarz mechanism, in which the unique nontrivial phase associated with $SU(2)_R$

⁴ We have ignored quantum corrections to the vacuum expectation values for the squark fields for simplicity.

symmetry breaks the supersymmetry, and the vacuum configuration is given by $\theta = \pi/2$ [6].⁵ Let us note that the adjoint Higgs scalar in this case is always massive except for $z_c = 0$, where the potential vanishes due to the original $\mathcal{N} = 1$ supersymmetry in five dimensions.

Let us next consider the case $z_g \neq z_s$ with $N_f = 1$. The sign of a function $C \equiv F(n, z_s) - F(n, z_g)$ is important to determine the minimum of the effective potential (7). For $z_s < (>)z_g$, the sign of C is negative (positive), so that the configuration $\theta = 0$ ($\pi/2$) is realized as the vacuum configuration. Therefore, depending on the relative magnitude between z_g and z_s , the vacuum configuration is different and accordingly, the gauge symmetry breaking patterns are different. If $z_g = z_s$, the effective potential vanishes due to the survived supersymmetry explained above. The adjoint Higgs scalar is always massive in this case.

Let us finally study the case $N_f = 2$ with $z_g \neq z_s$. In order to demonstrate the possible effects of the bare mass on the gauge symmetry breaking patterns, we take z_g to be 0.1, 1.0, 10 as an example. For each value of z_g we study the behavior of the effective potential with respect to z_s and find the minimum of the potential. We examine the stability of the configuration $\theta = 0(\text{mod}\pi)$ and $\theta = \pi/2$ with respect to z_s for the given values of z_g by studying the second derivative of the effective potential.

In case $z_g = 0.1$, simple numerical calculations show that the configuration $\theta = 0$ becomes unstable for $z_s \gtrsim 0.0672937 \equiv z_{s1}$, on the other hand $\theta = \pi/2$ becomes stable for $z_s \gtrsim 0.0706947 \equiv z_{s2}$. The configuration that minimizes the effective potential in the narrow region between z_{s1} and z_{s2} is still given by the configuration which breaks the $SU(2)$ gauge symmetry to $U(1)$, though it is not $\theta = \pi/2$. We confirm that by numerical analyzes, the behavior of θ that minimizes the potential in the narrow region is that the θ increases gradually from zero at z_{s1} and approaches to $\pi/2$ at z_{s2} . Thus, we have shown that the bare mass terms for the gaugino and squark can affect the gauge symmetry breaking patterns through the Hosotani mechanism and the phase transition occurs at $z_s = z_{s1}$ for $z_g = 0.1$. Hence, we obtain that⁶

$$\text{gauge symmetry breaking pattern} = \begin{cases} SU(2) \rightarrow SU(2) & \text{for } z_s < z_{s1}, \\ SU(2) \rightarrow U(1) & \text{for } z_s > z_{s1}. \end{cases} \quad (10)$$

It should be noted that the very small values of θ , which is usually of order $O(1)$, is possible in this case. This may affect mass spectrum in four dimensions. We will discuss this point in the last section.

We repeat the same analyzes as above for the case $z_g = 1.0$ (10) with $N_f = 2$. The configuration $\theta = 0$ becomes unstable for $z_s \gtrsim 0.618288$ (2.03287) $\equiv z_{s1}$, while the configuration $\theta = \pi/2$ becomes stable for $z_s \gtrsim 0.691531$ (2.47766) $\equiv z_{s2}$. The qualitative behavior of θ that minimizes the effective potential in the narrow region is the same as that in the case $z_g = 0.1$.

The adjoint Higgs scalar in this case can be massless unlike the previous cases. The second derivative of the effective potential evaluated at $\theta = 0$ ($\pi/2$) vanishes for $z_{s1} = 0.0672937$ (0.0706947). Hence, the massless state of the Higgs scalar is possible for the fine tuned values of z_s . In the other cases $z_g = 1.0, 10$, we also have massless state of the adjoint Higgs scalar at the values of z_s , where the second derivative of the potential evaluated at $\theta = 0, \pi/2$ vanishes.

We have seen that the gauge symmetry breaking patterns change due to the existence of the bare mass terms for the gaugino and squark in the model. The $SU(2)$ gauge symmetry is not broken for $z_s < z_{s1}$, on the other hand, $SU(2)$ is broken to $U(1)$ for $z_s > z_{s1}$ for fixed values of z_g in our examples.

⁵ The effective potential for the case of the Scherk–Schwarz mechanism is given by

$$V_{\text{eff}}^{\text{SS}} = \frac{\Gamma(\frac{5}{2})}{\pi^{\frac{5}{2}} L^5} (4N_f - 4) \sum_{n=1}^{\infty} \frac{1}{n^5} [1 - \cos(n\beta)] (2 + 2\cos(2n\theta)).$$

The Hosotani mechanism depends only on matter contents, so that we can quote the results obtained in [6].

⁶ The phase transition is the second order unlike the case of the nonsupersymmetric gauge model studied in Section 2.

If we add a bare mass term for the Dirac spinor ψ_D in the hypermultiplet instead of the squark ϕ_i , the structure of the effective potential is different from Eq. (6). It is easy to see that the $SU(2)$ gauge symmetry is never broken for any nonzero values of the bare masses.

4. Conclusions and discussions

We have demonstrated that the existence of the bare mass affects the gauge symmetry breaking patterns through the Hosotani mechanism. We have explicitly shown that in the nonsupersymmetric $SU(2)$ gauge model with the massive adjoint fermions defined on $M^4 \otimes S^1$, there exist the critical values for $z \equiv mL$, above (below) which the $SU(2)$ gauge symmetry is unbroken (broken). The phase transition is the first order. The asymptotic behavior of the effective potential with respect to z also suggests the existence of the critical values of $z \equiv mL$: if the adjoint fermion is heavy enough, corresponding to $z \rightarrow \infty$, it decouples from the effective potential and the gauge sector of the model dominates the potential. Hence, the $SU(2)$ gauge symmetry is not broken through the Hosotani mechanism. On the other hand, if we take the massless limit of the fermion, that is, $z \rightarrow 0$, the vacuum configuration breaks the $SU(2)$ gauge symmetry to $U(1)$ [5].

We have also studied the supersymmetric gauge model defined on $M^4 \otimes S^1$. Instead of the Scherk–Schwarz mechanism of supersymmetry breaking, we have introduced the bare mass terms for the gaugino in the vectormultiplet and the squark in the hypermultiplet to break the supersymmetry. When the number of the hypermultiplet N_f is equal to one, the critical point is given by $z_g = z_s$, where the potential vanishes due to the $\mathcal{N} = 1$ supersymmetry. The $SU(2)$ gauge symmetry is broken to $U(1)$ for $z_s > z_g$, while the gauge symmetry is not broken for $z_s < z_g$. If $N_f \geq 2$ and $z_g = z_s$, then, the $SU(2)$ gauge symmetry is always broken to $U(1)$ as long as $z_c (\equiv z_g = z_s) \neq 0$. In this case, the supersymmetry is broken by a unique bare mass z_c . The result is the same as the one obtained by the Scherk–Schwarz mechanism of supersymmetry breaking, in which the supersymmetry breaking parameter is also an unique and the gauge symmetry is always broken to $U(1)$. In these cases the adjoint Higgs scalar cannot be massless except that the models have the accidental $\mathcal{N} = 1$ supersymmetry.

We have considered the case $z_g \neq z_s$ for $N_f = 2$. We have shown the possible effect of the bare masses on the gauge symmetry breaking patterns through the Hosotani mechanism. By choosing the certain values of z_g , we have investigated the configuration that minimizes the effective potential according to the change of the values of z_s . We have found the critical values of z_{s1} , above (below) which the gauge symmetry is broken (unbroken). The phase transition in the supersymmetric model is the second order unlike the case of the nonsupersymmetric model. We have also found that the massless state of the adjoint Higgs scalar appears for the fine tuned values of z_s in this case.

There are many issues that are not discussed in this Letter. Let us comment on a few of them. In the supersymmetric gauge model discussed in Section 3, it is important to determine the behavior of the order parameter θ with respect to z_s precisely in the narrow region between z_{s1} and z_{s2} . As mentioned in the section, it is possible that the magnitude of the order parameter θ can be very small for (fine tuned) values of z_s . Then, if particle does not have a bare mass term, the mass square of $n = 0$ mode in the Kaluza–Klein modes behaves like $(\theta/L)^2$, so that the order of the mass is highly reduced compared with the compactification scale $1/L$ at the tree level. Therefore, we expect the light particle in four dimensions through the Hosotani mechanism.⁷

It may be interesting to study the case of massive fundamental fermion instead of the adjoint one. It has been known that the $SU(N)$ gauge symmetry is not broken for (supersymmetric) gauge model (with the Scherk–Schwarz mechanism) with massless fundamental fermion (matter) [6,7]. If we take the limit of the heavy bare mass, the fundamental fermion decouples from the effective potential and the potential is dominated by the gauge sector alone. Then, there are N physically equivalent vacua. On the other hand, in the massless limit, we expect that there

⁷ Let us note that the mass square of the particle is usually of order $(\theta/L)^2 \sim (O(1)/L)^2$ through compactification.

is a single $SU(N)$ symmetric vacuum for $N = \text{even}$ and a doubly degenerate $SU(N)$ vacuum for $N = \text{odd}$. Hence, we expect from these observations that there exist critical values of mL , at which a sort of phase transition, in which the number of the vacuum changes, occurs.

It may be interesting to consider higher rank gauge group and study the massive particle effect on the gauge symmetry breaking patterns. In particular, if we introduce the hierarchy among the bare masses, as we have done in the supersymmetric case, it may be expected to occur rich gauge symmetry breaking patterns. It is also interesting to study the mass spectrum in four dimensions, taking the smallness of θ into account, as discussed above.

One can also expect the same phenomena in other extra dimensions such as the orbifold S^1/Z_2 for example. According to the lessons obtained in this Letter, the gauge symmetry breaking patterns change even in the case of the orbifold if particles possess bare mass terms. It is expected that degeneracy of equivalent classes of boundary conditions, which has been discovered and discussed recently in [8], may be lifted due to the effect of the bare mass. These problems are under investigation and will be reported elsewhere.

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