New Derivation of Anomaly-Mediated Gaugino Mass in the Higher Derivative Regularization Method

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Abstract

The technique of higher derivative regularization is applied to the Super-Weyl-Kähler anomaly in supergravity coupled with chiral matters and gauge multiplet. Our method makes it possible to derive directly and diagrammatically the formulae for the gaugino mass (in the Abelian case). Our procedure is applied to derivation of not only the mass formulae known in the anomaly mediation scenario, but also additional terms that depend on the Kähler potential.

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§1. Introduction

One of the most important subjects in the present-day particle physics is to look for experimental signatures of supersymmetry (SUSY) and search for favorable sources of SUSY breaking mechanism. Among several others, extensive studies have been made to look for scenarios of SUSY breaking triggered by gravity\textsuperscript{11} and gauge interactions\textsuperscript{2}. Several years ago, a novel type of SUSY breaking mechanism was proposed which is often referred to as anomaly mediation mechanism. It has been argued in Refs.\textsuperscript{3} and\textsuperscript{4} that the SUSY breaking is communicated through the super-Weyl anomaly via supergravity.

The most appealing aspects of the anomaly mediation is its unique predictability of soft breaking terms. For example gaugino masses, scalar masses and triple coupling of matter multiplets are given respectively by

\begin{align}
    m_\lambda &= \frac{1}{3} \beta_g M^*, \\
    m_i^2 &= -|M|^2 \left( \frac{\partial \gamma_i}{\partial g} \beta_g + \frac{\partial \gamma_i}{\partial y} \beta_y \right), \\
    h^{ijk} &= \frac{1}{6} (\gamma_i + \gamma_j + \gamma_k) y^{ijk} M^*.
\end{align}

Here indices $i$, $j$ and $k$ label the chiral matter multiplets. These parameters are all determined in terms of the beta functions of the gauge coupling $\beta_g$ and the Yukawa coupling $\beta_y$ and the anomalous dimensions $\gamma_i$ of the $i$-th chiral multiplet field. $M$ is the auxiliary field of the gravity multiplet whose vacuum expectation value is the source of SUSY breaking.

In spite of its high predictability, the anomaly mediation scenario entails drawbacks of its own. Namely, scalar partners of leptons are given negative mass squared. Various ideas have been proposed to solve this tachyonic slepton mass problem\textsuperscript{4}–\textsuperscript{12}. They are, however, rather sophisticated and spoil the simplicity of the original proposal.

We would also like to note that there could exist additional terms in the formula (1.1)-(1.3) which depends on the matter Kähler potential.\textsuperscript{13} The minimal supergravity coupled with matter and gauge multiplets is invariant under the super-Weyl-Kähler transformation on the classical level. The quantum anomaly associated with this symmetry contains the Kähler potential on the order of $\kappa^2 = 8\pi G_N$. Such additional terms could hopefully change the nature of the slepton tachyonic problem.

Bearing these considerations in mind, we now examine diagrammatically the gaugino mass formula due to the super-Weyl-Kähler anomaly. We will see that the most suitable way to derive the gaugino mass directly is the use of the higher derivative regularization method. It has been known for some time that this method alone is not always very helpful at one-loop.\textsuperscript{14,15} It becomes useful only when we make a combined use of other regularizations,
typically such as Pauli-Villars (PV) method. Admitting such shortcoming, we would still like to show that the higher derivative regularization combined with the PV method turns out to be instrumental to uncover diagrammatical structures of the super-Weyl-Kähler anomaly and to produce formulae known in anomaly mediation.

The present paper is organized as follows. First we recapitulate the super-Weyl-Kähler anomaly in §2 and the higher derivative regularization in §3. We reproduce in §4 the mass formula (1.1) for the case of the Abelian gauge group. (To avoid unessential complications, we consider only the Abelian case throughout.) In §5 our method is applied further to derive the terms to be added to (1.1), which depends on the Kähler potential. §6 is devoted to conclusions.

§2. Super-Weyl-Kähler Anomaly

As we mentioned in §1 the minimal supergravity coupled with matter and gauge fields is invariant under the simultaneous transformation of super-Weyl and Kähler transformations on the classical level. The anomaly on the quantum level of this symmetry has been discussed in literatures to a considerable extent.\(^{15,16,17}\)

The fermionic contribution due to the super-Weyl-Kähler anomaly to the effective action has been evaluated as

\[
\frac{g^2}{96\pi^2} b_0 \text{Tr} F_{mn} \tilde{F}^{mn} \frac{1}{\Box} \partial c^\ell. \tag{2.1}
\]

Here the connection \(c^\ell = b^\ell + 4i\kappa^2 a^\ell\) consists of the auxiliary field \(b^\ell\) in the gravitational superfield and the Kähler connection \(a^\ell\). The first coefficient of the beta-function is denoted by \(b_0 = 3T(G) - T(R)\). This anomaly would be lifted to the unique supersymmetric expression

\[
\frac{g^2}{256\pi^2} \int d^4 \Theta 2\varepsilon W^\alpha W_\alpha \frac{1}{\Box} (\mathcal{D}^2 - 8R) \left\{4b_0 R^\dagger + \frac{\kappa^2}{3} T_R \mathcal{D}^2 K + \cdots \right\} + \text{h.c.}, \tag{2.2}
\]

if supersymmetry could be respected throughout. Here \(K\) is the Kähler potential and \(R\) is the gravity superfield expanded as

\[
R = -\frac{1}{6} \left\{ M + \cdots + \Theta^2 \left( -\frac{1}{2} \mathcal{R} - ie_\ell^a m \mathcal{D}_m b^a + \cdots \right) \right\}. \tag{2.3}
\]

The Einstein scalar curvature is denoted by \(\mathcal{R}\). (The ellipses in (2.2) correspond to the sigma-model anomaly, into which we do not delve throughout the present paper.)

It is easy to see that if the the auxiliary field \(M\) in the superfield \(R\) in (2.2) acquires a vacuum expectation value, the formula (2.2) gives a mass term for the gaugino in accordance
with the formula (1.1). A natural question arises, however: Although nothing is wrong with the formula (2.2), one would wonder if we could derive the gaugino mass formula in a more direct diagrammatic way which could get a more insight into the regularization. It is somewhat puzzling that a naive look at Feynman rules does not provide us with Feynman diagrams producing such a gaugino mass term. This puzzling situation is cleared only when we give details of the regularization. The present work is an outcome of our efforts to visualize the derivation of the gaugino mass in a more familiar way. In the following sections, we are going to set up and apply the method of higher derivative regularization combined with the PV-type method. This allows us to derive the gaugino mass in a direct way in accordance with the formula (1.1). (See also Ref. [18] for a diagrammatic approach to the SUSY breaking.)

§3. Higher Derivative Regularization

Here we give a general consideration of the supersymmetric generalization of the higher derivative regularization. We confine ourselves to the case of rigid supersymmetry and postpone the description of the supergravity case in §4 and §5.

In conventional field theories, the higher derivative regularization method tells us to modify the Lagrangian, for example, of the scalar field \( A \), in the following way:

\[
A^\dagger D^2 A \rightarrow A^\dagger D^2 A - \frac{1}{A^2} A^\dagger (D^2)^2 A. \tag{3.1}
\]

Here \( D^2 = D_n D^n \) and \( D_n \) is the gauge covariant derivative. If we take the limit \( A \rightarrow \infty \), (3.1) reduces to the conventional kinetic term of the scalar field. The supersymmetric extension of (3.1) is straightforward, and the Lagrangian of a superfield \( Q \) is modified as Ref. [15]

\[
\mathcal{L} = \int d^2\theta d^2\bar{\theta} \left( Q^\dagger e^{2gV} Q - \frac{1}{16A^2} Q^\dagger e^{2gV} D^2 e^{-2gV} D^2 e^{2gV} Q \right) + \left( \int d^2\theta \frac{1}{2} mQ^2 + \text{h.c.} \right). \tag{3.2}
\]

Here \( V \) is a vector multiplet. One can see that (3.2) is gauge invariant for both Abelian and non-Abelian cases. (Here and hereafter we follow the convention of Ref. [20]. As for the higher derivative method in SUSY case we refer the reader to Ref. [21] which contains many useful formulas.)

By decomposing (3.2) into component fields,

\[
Q = A + \sqrt{2} \theta \chi + F, \tag{3.3}
\]
we can easily derive propagators of scalar ($A$), spinor ($\chi$) and auxiliary ($F$) fields. We just list them up in order:

\[
<T(A(x)A(y))> = \frac{i}{\Box - m^2 - \Box^2/A^2}\delta^4(x-y),
\]  
(3.4)

\[
<T(A(x)F(y))> = <T(A^\dagger(x)F^\dagger(y))> = \frac{-im}{\Box - m^2 - \Box^2/A^2}\delta^4(x-y),
\]  
(3.5)

\[
<T(F(x)F^\dagger(y))> = \frac{i}{\Box - m^2 - \Box^2/A^2}\delta^4(x-y),
\]  
(3.6)

\[
<T(\chi_\alpha(x)\chi_\beta(y))> = i\delta_\alpha^\beta \frac{m}{\Box - m^2 - \Box^2/A^2}\delta^4(x-y),
\]  
(3.7)

\[
<T(\bar{\chi}_\alpha(x)\bar{\chi}_\beta(y))> = i\delta_\alpha^\beta \frac{m}{\Box - m^2 - \Box^2/A^2}\delta^4(x-y),
\]  
(3.8)

\[
<T(\chi_\alpha(x)\bar{\chi}_\beta(y))> = \sigma^m_{\alpha\beta}\partial_m \frac{1}{\Box - m^2 - \Box^2/A^2}\delta^4(x-y).
\]  
(3.9)

\[
<T(\bar{\chi}_\alpha(x)\chi_\beta(y))> = \frac{m'}{2}\phi^2 + h.c.
\]  
(3.10)

It is easy to confirm that these propagators reduce to the usual ones if we take the limit $\Lambda \to \infty$.

§4. The Gaugino Mass Formula (1.1) Revisited

Now let us extend the higher derivative method to the supergravity case and launch the analysis of the super-Weyl-Kähler anomaly. We separate our calculation into two parts: First we consider the anomaly term corresponding to $4b_0R^\dagger$ in the curly brackets of (2.2), and then proceed in §5 to the next term of order $\kappa^2$ in (2.2).

We are interested in the evaluation of the gaugino mass by using the following Lagrangian

\[
\mathcal{L} = \mathcal{L}_{\text{matter}} + \mathcal{L}_{\text{PV}},
\]  
(4.1)

which consists of a massless ($m = 0$) chiral multiplet $Q$ in $\mathcal{L}_{\text{matter}}$ and of massive PV regulator chiral field $\Phi$ (with mass $m'$) in $\mathcal{L}_{\text{PV}}$. More explicitly each term in (4.1) is expressed as

\[
\mathcal{L}_{\text{matter}} = \int d^2\Theta 2\Sigma \left[ -\frac{1}{8}(\bar{D}^2 - 8R) \left\{ Q^\dagger e^{2gV}Q - \frac{1}{16A^2}Q^\dagger e^{2gV}(\bar{D}^2 - 8R)e^{-2gV}D^2e^{2gV}Q \right\} \right] + h.c.,
\]  
(4.2)

\[
\mathcal{L}_{\text{PV}} = \int d^2\Theta 2\Sigma \left[ -\frac{1}{8}(\bar{D}^2 - 8R) \left\{ \phi^\dagger e^{2gV}\phi - \frac{1}{16A^2}\phi^\dagger e^{2gV}(\bar{D}^2 - 8R)e^{-2gV}D^2e^{2gV}\phi \right\} + m'^2\phi^2 \right] + h.c.,
\]  
(4.3)
This Lagrangian gives many additional terms containing $1/\Lambda^2$, and our text would be too much cluttered if we would write down all the Feynman rules. Since we are interested in the gaugino mass and therefore in the single insertion of the auxiliary field $M$ in (2.3), we pay our attention to the vertex containing $M$ singly. Some of the Feynman rules relevant to our calculation are illustrated in Fig. 1, where those containing the matter field $Q$ are given: the Feynman rules of PV-field $\Phi$ are exactly the same.

Being equipped with the Feynman rules, we are now in a position to compute the triangle diagrams of the type Fig. 2 giving rise to the gaugino mass term. There are several combinations of propagators running through the triangle and they are listed in Table 1. There are seven diagrams in which massive PV-field are circulating. Its component fields are all primed, i.e.,

$$\Phi = A' + \sqrt{2} \theta \chi' + F'. \quad (4.4)$$

The contribution of the massless chiral matter $Q$, the first entry in Table 1, is divergent, and gives an effective action

$$\mathcal{L}_{\text{massless}} \sim \frac{g^2}{16\pi^2} \frac{M^*}{3} \lambda \left\{ -\frac{1}{2} + 2\log \left( \frac{A^2}{\Lambda^2} \right) \right\}, \quad (4.5)$$

![Fig. 1. Feynman rules for vertices containing either the auxiliary field $M^*$ or matter field. Those of the PV-fields are the same and are omitted here. The gaugino field is denoted by $\lambda$.](image-url)
The limit $\Lambda$ is also divergent. All the remaining terms become vanishing when we take the ultra-violet cutoff, i.e., (4.6), (4.7) and (4.8) in the following way:

\[
\mathcal{L}_{PV} = \frac{g^2}{3i} M^* \lambda \left[ \frac{d^4p}{(2\pi)^4} \left( -p^2 - m^2 - (p^2)^2 / A^2 \right) \frac{1}{A^2} \left( 1 - \frac{p^2}{2A^2} \right) \left( 1 + \frac{2p^2}{A^2} \right) \right] \sim \frac{g^2}{16\pi^2} \frac{M^*}{3} \lambda \left\{ -\frac{1}{2} + 2\log \left( \frac{A^2}{\Lambda^2} \right) + \mathcal{O}(m^2 / A^2) \right\},
\]

where $A_c$ is the ultra-violet cutoff. The gaugino field is denoted by $\lambda$. The other contributions, $PV(1)$-(3), to the effective action are given, respectively, by

\[
\mathcal{L}_{PV(1)} = \frac{g^2}{3i} M^* \lambda \left[ \frac{d^4p}{(2\pi)^4} \left( -p^2 - m^2 - (p^2)^2 / A^2 \right) \frac{1}{A^2} \left( 1 - \frac{p^2}{2A^2} \right) \left( 1 + \frac{2p^2}{A^2} \right) \right] \sim \frac{g^2}{16\pi^2} \frac{M^*}{3} \lambda \left\{ -2 + \mathcal{O}(m^2 / A^2) \right\},
\]

\[
\mathcal{L}_{PV(2)} = \frac{g^2}{3i} m^2 \lambda \left[ \frac{d^4p}{(2\pi)^4} \left( -p^2 - m^2 - (p^2)^2 / A^2 \right) \frac{1}{A^2} \left( 1 - \frac{p^2}{2A^2} \right) \left( 1 + \frac{2p^2}{A^2} \right) \right] \sim \frac{g^2}{16\pi^2} \frac{M^*}{3} \lambda \left\{ -2 + \mathcal{O}(m^2 / A^2) \right\}.
\]

Note that (4.6) is also divergent. All the remaining terms become vanishing when we take the limit $\Lambda \to \infty$, i.e.,

\[
\mathcal{L}_{PV(4)} = \mathcal{L}_{PV(5)} = \mathcal{L}_{PV(6)} = \mathcal{L}_{PV(7)} = \mathcal{O}(m^2 / A^2).
\]

Although we have been dealing with a single PV-field, we can easily generalize our argument for the case of several PV-fields. In such a case we just superpose the sum of PV-contribution, i.e., (4.6), (4.7) and (4.8) in the following way:

\[
\frac{g^2}{16\pi^2} \frac{M^*}{3} \left[ \left\{ -\frac{1}{2} + 2\log \left( \frac{A^2}{\Lambda^2} \right) \right\} + \sum_i C_i \left\{ \left( -\frac{1}{2} - 2 + 3 \right) + 2\log \left( \frac{A^2}{\Lambda^2} \right) \right\} \right].
\]
Here $C_i$ is the weight factor of each PV-fields, the index $i$ numbering the PV-fields. It is now obvious that the cancellation of the ultraviolet divergences between the massless chiral matter and PV-fields is fulfilled by

$$1 + \sum_i C_i = 0. \quad (4.11)$$

Thus the generated gaugino mass is found to be

$$m_\lambda = \frac{g^2}{16\pi^2} \frac{M^*}{3} \left\{ -\frac{1}{2} - \left( -\frac{1}{2} - 2 + 3 \right) \right\} = -\frac{g^2}{16\pi^2} \times \frac{M^*}{3}, \quad (4.12)$$

which agrees with previous results for $U(1)$ case.

§5. Kähler Potential and the Gaugino Mass

We now move on to the anomalous term associated with the Kähler connection, the second term in (2.2). This term appears on the order of $O(\kappa^2)$. For the purpose of evaluating it, we have to extend the higher derivative regularization by including those of $O(\kappa^2)$. To explain the higher derivative terms of $O(\kappa^2)$ with reference to those of §4, we restart from the Kähler potential of the following form:

$$K = Q^\dagger Q + \Phi^\dagger \Phi, \quad (5.1)$$

where the higher derivative terms are not yet included. The massless chiral matter field is again denoted by $Q$ and $\Phi$ is a massive generic PV-chiral field. The Lagrangian may be expanded in the power series of the gravitational constant $\kappa^2$

$$\mathcal{L} = \frac{1}{\kappa^2} \int d^2\Theta 2\mathcal{E} \left\{ \frac{3}{8} \left( \mathcal{D}^2 - 8R \right) e^{-\kappa^2 K/3} \right\} + \text{h.c.}$$

$$= -\frac{6}{\kappa^2} \int d^2\Theta \mathcal{E} R + \mathcal{L}_1 + \mathcal{L}_2 + O(k^4) + \text{h.c.}, \quad (5.2)$$

where

$$\mathcal{L}_1 = \int d^2\Theta 2\mathcal{E} \left\{ -\frac{1}{8} \left( \mathcal{D}^2 - 8R \right) K \right\}, \quad (5.3)$$

$$\mathcal{L}_2 = \kappa^2 \int d^2\Theta 2\mathcal{E} \left\{ \frac{1}{48} \left( \mathcal{D}^2 - 8R \right) K^2 \right\}. \quad (5.4)$$

Comparing (4.2), (4.3) and (5.3), we immediately realize that inclusion of the higher derivative terms, in the absence of the gauge interaction, is achieved by replacing the Kähler potential (5.1) by

$$K = Q^\dagger \left\{ 1 - \frac{1}{16A^2} \left( \mathcal{D}^2 - 8R \right) \mathcal{D}^2 \right\} Q + \Phi^\dagger \left\{ 1 - \frac{1}{16A^2} \left( \mathcal{D}^2 - 8R \right) \mathcal{D}^2 \right\} \Phi. \quad (5.5)$$
At the $\mathcal{O}(\kappa^2)$ level, our higher derivative terms are obtained by putting (5.5) into (5.4). Among many terms, those relevant to our anomaly calculation are

\[
\mathcal{L}_2 = -\frac{\kappa^2}{6} \left\{ K_i \left( 1 - \frac{\Box}{A^2} \right) A^i \right\} \left\{ F^\dagger \left( 1 - \frac{\Box}{A^2} \right) F + F'^\dagger \left( 1 - \frac{\Box}{A^2} \right) F' \right\} \\
- \frac{\kappa^2}{6} \left\{ K_i \left( 1 - \frac{\Box}{A^2} \right) F^i \right\} \left\{ F'^\dagger \left( 1 - \frac{\Box}{A^2} \right) K_j \right\} + \text{h.c.} + \cdots. \tag{5.6}
\]

Here our notations are

\[
K_i \left( 1 - \frac{\Box}{A^2} \right) A^i = A^\dagger \left( 1 - \frac{\Box}{A^2} \right) A + A'^\dagger \left( 1 - \frac{\Box}{A^2} \right) A', \tag{5.7}
\]

\[
K_i \left( 1 - \frac{\Box}{A^2} \right) F^i = A^\dagger \left( 1 - \frac{\Box}{A^2} \right) F + A'^\dagger \left( 1 - \frac{\Box}{A^2} \right) F', \tag{5.8}
\]

\[
F'^\dagger \left( 1 - \frac{\Box}{A^2} \right) K_j = F^\dagger \left( 1 - \frac{\Box}{A^2} \right) A + F'^\dagger \left( 1 - \frac{\Box}{A^2} \right) A'. \tag{5.9}
\]

Thus we arrive at the Feynman rule depicted in Fig. 3 for the vertex that contains $K_i F^i$. Note that the vertex in Fig. 3 is what we are interested in to derive the second term in (2.2), i.e., it is the $\theta$-independent term in $\mathcal{D}^2 K$

\[
\mathcal{D}^2 K = -4K_i F^i + \cdots. \tag{5.10}
\]

Fig. 3. The Feynman rule for the vertex containing $K_i F^i$

Now the anomalous $(K_i F^i) \lambda \lambda$ vertex is produced by the Feynman diagrams in Fig. 4. Various combinations of the propagators in Fig. 4 are summarized in Table 2. From the Feynman rule in Fig. 3 it is almost obvious that we can borrow the loop calculations in §4 simply by replacing

\[
M^* \longrightarrow \kappa^2 K_i F^i \tag{5.11}
\]

To sum up, our loop calculations of Fig. 4 are as follows:

\[
\mathcal{L}_{\text{masless}} \sim \frac{g^2}{12\pi^2} \kappa^2 (K_i F^i) \lambda \lambda \left\{ \frac{1}{2} - \log \left( \frac{A^2}{\Lambda^2} \right) \right\}, \tag{5.12}
\]
Table 2. Various combinations of the Feynman diagrams in Fig. 4

| PV(i)   | (a)       | (b)       | (c)       |
|---------|-----------|-----------|-----------|
| massless| $< AA^\dagger >$ | $< \chi \bar{\chi} >$ | $< FF^\dagger >$ |
| PV(ii)  | $< A^\prime A^\dagger >$ | $< \chi^\prime \bar{\chi} >$ | $< F^\prime F^\dagger >$ |
| PV(iii) | $< A^\prime F^\dagger >$ | $< \chi^\prime \bar{\chi} >$ | $< A^\dagger F^\prime >$ |
| PV(iv)  | $< A^\prime F^\prime >$ | $< \chi^\prime \bar{\chi} >$ | $< F^\prime F^\dagger >$ |

Fig. 4. The triangle diagram producing $(K_i F^i) \lambda \lambda$ term. Various combinations of the propagators (a), (b) and (c) are given in Table 2.

Again we are able to increase the number of PV-regulator arbitrarily, and after summing up all regulators with the weight $C_i$, the coefficient of the gaugino mass term becomes

$$\mathcal{L}_{PV(i)} \sim \frac{g^2}{12\pi^2} \kappa^2 (K_i F^i) \lambda \lambda \left\{ \frac{1}{2} - \log \left( \frac{A^2}{\Lambda^2} \right) + \mathcal{O}(m^2/\Lambda^2) \right\},$$  
(5.13)

$$\mathcal{L}_{PV(ii)} \sim \frac{g^2}{12\pi^2} \kappa^2 (K_i F^i) \lambda \lambda \left\{ -\frac{1}{2} + \mathcal{O}(m^2/\Lambda^2) \right\},$$  
(5.14)

$$\mathcal{L}_{PV(iii)} = \mathcal{L}_{PV(iv)} = \mathcal{O}(m^2/\Lambda^2).$$  
(5.15)

The condition canceling the ultra-violet divergence is the same as before

$$1 + \sum_i C_i = 0,$$  
(5.17)

and we get the gaugino mass coming from the $K_i F^i$ as

$$\frac{g^2}{16\pi^2} \times \frac{2}{3} \kappa^2 (K_i F^i).$$  
(5.18)

This result is in agreement with the naive one obtained by using (2.2).

§6. Conclusion

In the present paper, we have examined the calculational basis of the super-Weyl-Kähler anomaly. We have pointed out that the higher derivative regularization method combined with the type of PV’s is useful to derive the gaugino mass formula in a direct and diagrammatic method. The formulae derived in the anomaly mediation scenario are thus put on as familiar footing as our old-day anomaly calculation.
There remain, however, several important problems. The most imminent is a generalization to the non-Abelian gauge group. In principle there does not exist a serious stumbling block in such a generalization, but the higher derivative counter terms become considerably involved. We will come to the non-Abelian generalization in our future publication.\textsuperscript{22}

Another problem is to reconsider the slepton mass problem, considering the additional contributions of order $\kappa^2$. We have to look for a model in which a natural vacuum expectation value is given to $K_i F^i$ and to give a positive definite mass squared to sleptons. In connection with this problem, we should also include the sigma-model anomaly. Although we did not touch on the sigma-model anomaly at all in this paper, we believe that our method is equally useful to analyze the sigma-model anomalous terms. All of these belong to our future works.

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