Gravitational and Schwinger model anomalies: how far can the analogy go?

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ABSTRACT

We describe the most general treatment of all anomalies both for chiral and massless Dirac fermions, in two-dimensional gravity. It is shown that for this purpose two regularization dependent parameters are present in the effective action. Analogy to the Schwinger model is displayed corresponding to a specific choice of the second parameter, thus showing that the gravitational model contains anomaly relations having no analogy in the Schwinger model.
Two dimensional models of massless fermions coupled to gravity have been extensively used as toy-models for the understanding of gauge and gravitational anomalies. A variety of method of calculations has been developed and various specific models have been studied. Apparently, two distinct situations are present: models with either Dirac or Weyl fermions. The former presents the freedom of preserving one of the classical symmetries at the expense of the other, while in the latter anomalies are always present. This apparent discrepancy has been explained, and a unified description of anomalies in both cases has been given in the case of the Schwinger model \cite{1}. A preliminary discussion according to such a vantage viewpoint, in the case of gravity, has been given as well \cite{2}. However, the resulting effective action contained a multitude of arbitrary parameters since the local parts have been constructed in all possible ways in terms of zweibein, and, though giving correct results, an analogy to the Schwinger model was not evident. Furthermore, one has to bear in mind that in the case of fermions coupled to gravity one has three classical symmetries to start with, while in the case of the Schwinger model one encounters only two gauge symmetries. Thus, it is reasonable to expect, but necessary to verify explicitly, that it will not be possible to parametrize all the regularization ambiguities in terms of a single parameter, as it results from ref.(2). Notwithstanding, the classical analogy between the Schwinger model and fermions in a two-dimensional curved spacetime is compelling. The results of ref.(1) show that the zweibein is not a proper choice as a gauge field bearing similarity to the gauge potential \( A_\mu \). Another possible choice is the spin-connection \( \omega^{ab}_\mu \) that in two dimensions can be always written as

\[
\omega^{ab}_\mu = \epsilon^{ab} \omega_\mu .
\]  

(1)
On the other hand, the spin-connection is expressed in terms of the zweibein through the vanishing torsion constraint. In two-dimensions its *linearized* form, 
\( e^a_\mu = \delta^a_\mu + h^a_\mu \), is given by *

\[
\omega_\mu = \epsilon^{ab} \partial_b \bar{h}_{a\mu} + \frac{1}{2} \epsilon_{a\mu} \partial^a h + \frac{1}{2} \partial_\mu L \tag{2}
\]

where we decompose the “graviton” field in traceless symmetric part \( \bar{h}_{a\mu} \), totally anti-symmetric part \( h_{[a\mu]} \), and trace \( h \), according to

\[
h_{a\mu} = \bar{h}_{a\mu} + \frac{1}{2} \delta_{a\mu} h + h_{[a\mu]} \tag{3}
\]

finally, \( L = \epsilon^{ab} h_{ab} \) and \( h = h_{a\mu} \delta^{a\mu} \) are the Lorentz and Weyl degrees of freedom respectively.

From eq.(2) we learn that:

i) in two dimensions \( \omega_\mu \) should have only two components, therefore the decomposition in (2) shows that some of the components must be gauge degrees of freedom. Actually, at the classical level all of them can be gauged away since gravity has no dynamics. On the quantum level this indicates that, if anomalies are present, *they combine together in such a way that \( \omega_\mu \) has at most two dynamical components.*

ii) Comparing (2) to the decomposition of the gauge field \( A_\mu \) for the Schwinger

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* We maintain at the linearized level too the mixed notation with latin and greek indices, even if it is not strictly necessary because in this approximation all the indices are flat. This is to avoid confusion with the “standard” notation, where the graviton field \( h_{\mu\nu} = g_{\mu\nu} - \eta_{\mu\nu} \) is symmetric by definition, while in our case the linearized zweibein \( h_{a\mu} \) has no definite symmetry property under \( a \leftrightarrow \mu \) exchange.
model

\[ A_\mu = \partial_\mu \phi + \epsilon_{\mu\nu} \partial_\nu \rho \]  

(4)

we expect a close analogy between gravitational and Schwinger model only in the case \( \bar{h}_{a\mu} = 0 \), which amounts to imposing general coordinate invariance at the quantum level.

With these general arguments in mind, we are going to perform an explicit anomaly calculation starting from the action

\[ W_F = \frac{i}{2} \int d^2 x e e_a^\mu \bar{\psi}_a \gamma^a \left( \partial_\mu - \omega_{\mu ab} \Sigma^{ab} \right) \left( \frac{1 - \beta \gamma^5}{2} \right) \psi + \text{H.C.} , \]  

(5)

where \( e \equiv \det e_a^\mu \), and the parameter \( \beta \) has been introduced to be able to consider simultaneously both Dirac(\( \beta = 0 \)) and Weyl fermions(\( \beta = \pm 1 \)).

To clarify our arguments we shall first compute the effective action in a manifestly general covariant scheme and introduce the regularization parameter not “ by hand ”, i.e. as the coefficient of a local counter-term, but rather through the formal properties of the two-dimensional Dirac operator\(^3\). Then, we shall extend the above scheme to the case where all the three symmetries are broken at the quantum level, and introduce a second arbitrary parameter.

Various method of calculation are at our disposal. One would naturally tend to perform perturbative calculations. However, if one is to keep \( \omega_\mu \) as a genuine gauge field throughout the calculations one is faced with the problem of vanishing spin current in two dimensions, leading to the lack of a “ current-gauge field ” vertex. Then, either one continues \( \gamma^5 \) out of two dimensions\(^4\), which has created some confusion in the literature on how useful these prescription are for an explicit
calculation, or works indirectly in terms of “zweibein-energy momentum tensor” vertices\textsuperscript{[5]} and reconstructs the spin-connection in the final results. To make the long story short and transparent we opt for the Fujikawa approach\textsuperscript{[6]}. The basis of this algorithm is a proper choice of the regularizing operator of the Jacobian arising from the symmetry transformations of the functional measure. As an appropriate choice we choose as a regularizing operator in Minkowski space\textsuperscript{*}

\begin{equation}
D = ie_m \gamma^m \partial_\mu + e_m \gamma^{5}_m \omega_\mu \left( \frac{1 - \beta \gamma^5}{2} \right) .
\end{equation}

Following the usual procedure, we obtain, from eq.(6), the euclidean form of the Fujikawa regulator as

\begin{equation}
D_E = ie_m \gamma^m \partial_\mu + \frac{1}{2} e_m \gamma^{5}_m \Omega^5_\mu + \frac{i}{2} e_m \gamma^{5}_m \Omega_\mu ,
\end{equation}

where

\begin{align}
\Omega^5_\mu & \equiv \left( a g_{\mu \nu} + \frac{\beta}{2} \sqrt{-g} e_{\mu \nu} \right) \omega^{5 \nu} , \\
\Omega_\mu & \equiv \left( (1 - a) g_{\mu \nu} + \frac{\beta}{2} \sqrt{-g} e_{\mu \nu} \right) \omega^\nu ,
\end{align}

and we have introduced an arbitrary parameter $a$ exploiting the two-dimensional identity $\gamma^\mu \gamma^5 = \epsilon^{\mu \nu} \gamma^\nu$. The final result of such a procedure, after analytic continuation to imaginary time of the Dirac operator, is a splitting of $D_E$ into the sum of an hermitian and a non-hermitian part. Although such a splitting may

\textsuperscript{*} We shall adopt the following conventions for the (minkowskian) Dirac matrices:

\begin{equation}
\{ \gamma^a, \gamma^b \} = 2\eta^{ab}, \eta_{00} = +1, \eta_{11} = -1 ; \quad \gamma^5 = \frac{1}{2} \epsilon^{ab} \gamma_a \gamma_b .
\end{equation}
seem straightforward it may be done in different ways (at least in the case of chiral fermions) and one has to be careful in order to find the correct anomalies. The introduction of an arbitrary parameter $a$ represents, in the Fujikawa approach, the counterpart of perturbative calculation ambiguities corresponding to different choices of the regularization method.

Using known formulae\textsuperscript{[6]} we can obtain the anomalous Ward identities for Weyl and Lorentz symmetries as

\[ T = i \text{tr} \left( A_1 \left[ D_E^\dagger D_E \right] - A_1 \left[ D_E D_E^\dagger \right] \right), \] (9a)

\[ T_{[\mu \nu]} = -\frac{i}{2} \sqrt{-g} \epsilon_{\mu \nu} \text{tr} \left( \gamma^5 A_1 \left[ D_E^\dagger D_E \right] + \gamma^5 A_1 \left[ D_E D_E^\dagger \right] \right), \] (9b)

where $A_1 [\Delta]$ is the second coefficient of the heat-kernel expansion for the elliptic, second order, differential operator $\Delta$. The above formulae show the need to consider non-hermitian operators in the Fujikawa approach in order to obtain all the anomalies. The choice of the hermitian part of the operator in eq.(7) corresponds to the Weyl invariant regularization, while the choice of an anti-hermitian Dirac operator corresponds to the Lorentz invariant regularization. An arbitrarily weighted combination of the two gives both anomalies.

Rotating back to Minkowski spacetime we obtain\textsuperscript{*}

\[ T = \frac{1}{\sqrt{-g}} \epsilon^{\mu \nu} \partial_\mu \Omega_\nu, \] (10a)

\[ T_{[\mu \nu]} = -\epsilon_{\mu \nu} \epsilon^{\rho \sigma} \partial_\rho \Omega_5^{\sigma}. \] (10b)

\textsuperscript{*} For simplicity we have rescaled all the following formulae by the global numerical coefficient appearing in the effective action as $1/192\pi$. 

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With the help of eq.(8), eq.(10) can be written in a more familiar and Schwinger-looking form

\[ T = 2(1 - a)R + \beta \partial_\mu \omega^\mu \]  \hspace{2cm} (11a)

\[ T_{[\mu\nu]} = -\frac{1}{2}\epsilon_{\mu\nu} (\beta R + 2a\partial_\mu \omega^\mu) \]  \hspace{2cm} (11b)

From the explicit expression of the anomalies one can construct the corresponding effective action. The easiest way to do it, is to notice that the trace and the anti-symmetric part of the energy-momentum tensor couple to the trace and to the anti-symmetric part of the zweibein and then expressing the later two in terms of spin-connection. In this way, one finds the effective action with the help of eqs.(2) and (11):

\[ I^{\text{eff}} = \int d^2x \sqrt{-g} \left[ R \frac{1}{\nabla^2} (R + \beta \partial_\mu \omega^\mu) + a g^{\mu\nu} \omega_\mu \omega_\nu \right] , \hspace{2cm} (12)\]

where \( \nabla^2 = \nabla_\mu \nabla^\mu \) is the scalar covariant D’Alambertian. Before we proceed let us make a few comments on the obtained results.

First of all, as previously explained, we have considered a particular situation in which the anomalies of only Lorentz and Weyl symmetry are considered, while general coordinate invariance is assumed to be valid at the quantum level. Clearly, this is not the most general situation, but certainly the one that exhibits analogy to the Schwinger model.

Secondly, comparison to Schwinger model effective action shows an asymmetric form of the non-local part in eq.(12). It is our believe that this has created a certain confusion in the literature \(^7\) and we would like to explain this apparent difference.
It is possible to show that there is a relation between the symmetric and the asymmetric form of the effective actions. This is due to the relation

\[ R^2 = (\nabla_\mu \omega^\mu)^2 - \nabla_\mu \nabla^\mu \omega_\nu \omega^\nu \]  

which leads to

\[ R \frac{1}{\nabla^2} \left( R + \frac{2\beta}{1 + \beta^2} \nabla_\mu \omega^\mu \right) + b \omega_\mu \omega^\mu \equiv \left( \frac{1}{1 + \beta^2} \right) \left[ (R + \beta \nabla_\mu \omega^\mu)^2 (R + \nabla_\mu \omega^\mu) + a \omega_\mu \omega^\mu \right] \]  

where the two arbitrary parameters \( a \) and \( b \) are related by \( b = \frac{a + \beta^2}{1 + \beta^2} \). Both (12) and (15) lead to the same anomalies once the above relation between the parameters is taken into account. Then, both form of the effective action can be used on the same footing, and (12) can be rewritten, in complete analogy to the Schwinger model, as

\[ I_{\text{eff}} = \int d^2 x \sqrt{-g} \omega_\mu \left[ a g^{\mu\nu} - \left( \frac{1}{\sqrt{-g}} \epsilon^{\mu\rho} + \beta g^{\mu\rho} \right) \nabla_\rho \frac{1}{\nabla^2} \nabla_\sigma \left( \frac{1}{\sqrt{-g}} \epsilon^{\sigma\nu} + \beta g^{\sigma\nu} \right) \right] \omega_\nu . \]  

As we hope to have explained at the beginning, and confirmed through our explicit calculation, the analogy to the Schwinger model (i.e. arbitrariness in terms of a single parameter) is possible as long as we consider only Lorentz and Weyl symmetries, assuming quantum general covariance. We would like to go further and consider all three symmetries on the same footing, i.e. without making any assumptions on their validity at the quantum level. To do so we need to devise a way of breaking general covariance. Since all the components of the zweibein transform under
general coordinate transformations, contrary to Weyl and Lorentz symmetry, the presence of all the components in covariant expressions is a necessary condition to guarantee covariance, and the absence of one of them will spoil it. Hence, it seems the most convenient to introduce an additional parameter $c$ in the decomposition (3) as

$$h^{(c)}_{a\mu} = \bar{h}_{a\mu} + \frac{c}{2}(\delta_{a\mu}h + 2h_{[a\mu]}).$$  

(16)

We shall not proceed by any explicit choice of regularization but rather consider the problem in its full generality.

One finds linearized transformations of the zweibein (16) as

$$\delta h^{(c)}_{a\mu} = \frac{c+1}{2}\partial_\mu \xi_a + \frac{1-c}{2}\partial_a \xi_\mu + c\frac{1-1}{2}\delta_{a\mu} \partial_d \xi^d + c\Lambda \delta a_\mu - c\theta \epsilon a_\mu.$$  

(17)

where $\xi^\mu$, $\Lambda$ and $\theta$ are covariant, Weyl and Lorentz transformation parameters. Further, using eq.(2) in its linearized form, we obtain the following variations

$$\delta \omega^{(c)}_{\mu} = 1 - c^2 \epsilon de \partial_e \partial_d \delta_\mu \xi_a + \frac{c-1}{2} \epsilon \mu d \partial_d \delta_\tau \partial_a \xi^a + c \epsilon_{\mu a} \partial^a \Lambda - c \partial_\mu \theta$$  

(18a)

$$\delta \partial^\mu \omega^{(c)}_{\mu} = \frac{1-c}{2} \epsilon e de \partial^2 \partial_\mu \xi_d - c \partial^2 \theta$$  

(18b)

$$\delta \epsilon^\nu \partial_\nu \omega^{(c)}_{\mu} = \frac{c-1}{2} \partial^2 \partial_a \delta^d + c \partial^2 \Lambda \equiv -\delta R^{(c)}.$$  

(18c)

From eq.(18c) one can see that the linearized variation of the Ricci scalar with respect to a general coordinate transformation is zero for the choice $c = 1$ (which leads to the decomposition (3)). If $c \neq 1$ general covariance is spoiled as expected from the above discussion. With the generalized choice (16), starting from the
symmetric form of the effective action (15), and using eq.(18) one can find anomaly
relations for general covariance, Lorentz and Weyl symmetries as

\[ \nabla^\nu T_{\mu\nu} = (1 - c) \left\{ \nabla_\mu \left[ (1 - a) R + \beta \nabla_\rho \omega^\rho \right] + \right. \\
\left. \sqrt{-g} \epsilon_{\mu\rho} \nabla^\rho \left[ (a + \beta^2) \nabla_\sigma \omega^\sigma + \beta R \right] \right\} \]  
(19a)

\[ T = -c \left[ (1 - a) R + \beta \nabla_\mu \omega^\mu \right] \]  
(19b)

\[ T_{[\mu\nu]} = \frac{c}{2} \epsilon_{\mu\nu} \left[ \beta R + (a + \beta^2) \nabla_\rho \omega^\rho \right] . \]  
(19c)

It is evident from eqs.(19) that introduction of the second parameter was crucial to
put all the anomalies of the quantum gravity on the same footing and, therefore,
to describe the most general situation. It contains all the partial results present in
the literature as specific choices of the constants \( a, c, \beta \), and gives further insight
in the interplay among various anomalies in quantum gravity.

First of all, one notices the emergence of the general pattern conjectured at the
beginning, and based on the spin-connection decomposition, i.e. either one chooses
to preserve general covariance \( (c = 1) \) introducing Weyl and Lorentz anomalies,
or breaks general covariance preserving Weyl and Lorentz symmetry \[^{[8]}\] \( (c = 0) \).
In the former case one can see that eqs.(19) leads to the anomaly relations (11)
and the analogy to the Schwinger model is due to the absence of the covariant anomaly.

Within this general scheme there is still a further freedom related to the choice
of the parameters \( a \) and \( \beta \). It is generally believed that in models with Dirac

\[^{\star}\] Although we used linearized variations (18) at the end of calculations we have reconstructed
covariant results. One could have used the Fujikawa approach with the decomposition (17),
which gives the same anomalies (19).
fermions the interplay of anomalies is between general covariance and Weyl symmetry, preserving Lorentz symmetry.

This view needs some explanation based on the results in eqs.(19): as already said, parameter $c$ is relevant for absence of general covariant anomaly while parameter $a$ interpolates (once $c$ is chosen) between Lorentz and Weyl anomalies. Therefore, insisting on the absence of general covariant anomaly one can, in principle, still have Lorentz anomaly and no Weyl anomaly ($a = 0$). However, the opposite is usually preferred ($a = 1$) but not necessary, and surely, absence of covariant anomaly does not imply a priori absence of Lorentz anomaly.

The case of chiral fermions ($\beta = \pm 1$) is even more intriguing since no choice of the parameter $a$ can remove Lorentz or Weyl anomalies (similarly to the Schwinger model), which are, therefore, genuine quantum effects [2,9], and not regularization dependent pathologies.

Various possibilities exist in this case. One can always produce combinations of $T \pm e^{ab}T_{ab}$ which are proportional to either $1 + \beta$ or $1 - \beta$, and choosing, say $\beta = 1$, only one of them survives. The analogy with the chiral Schwinger model, where the anomalous current is the combination of the gauge and axial currents and no choice of the parameter can eliminate the anomaly of this combined current [1], is then evident. This remark is relevant as long as one wants to write the induced effective action in terms of only one light-cone component of the gauge field, i.e. the one that is classically coupled to chiral fermions. Otherwise, one can keep the local term ($a \neq 0$) in which both light-cone components of the gauge field are present [10].
In the case of the spin-connection, the situation is a bit more general than in the Schwinger model, and we feel it deserves a detailed explanation. Chiral models are best described in terms of light-cone components and from now on we shall stick to this notation *. The decomposition (2) can be written as

\begin{align}
\omega_- &= \partial_+ \bar{h}_- + \frac{1}{2} \partial_-(L - h) \\
\omega_+ &= \partial_- \bar{h}_+ + \frac{1}{2} \partial_+(L + h)
\end{align}

and eqs. (20) can be compared to the light-cone decomposition of the gauge field $A_\mu$ for the Schwinger model

$$A_\pm = \partial_\pm (\phi \pm \rho).$$

The scheme we previously described for chiral models corresponds to having only $\omega_+$ or $\omega_-$ ($a = 0$). From eq. (20) it is possible to have only one component of the spin-connection, in the chiral models, by using residual symmetries to impose, say $\bar{h}_{\pm \pm} = 0$ and $L = -h$, thus obtaining the anomaly equation

$$T_{-+} = \partial_+ \omega_-.$$ 

Keeping in mind that $T_{-+} \equiv (\delta_{ab} - \epsilon_{ab})\partial^a S^b = \partial_- S_+$, with $S^a$ the quantum

* The scalar product is written in terms of light-cone components as

$$A_\mu B^\mu \equiv \frac{1}{2} (A_+ B_- + A_- B_+).$$
spin-current, the preceding relation can be written as

\[ \partial_- S_+ = \partial_+ \omega_- , \] (23)

which is apparently the Schwinger model anomaly with substitution \( S_+ \rightarrow J_+ \) and \( \omega_- \rightarrow A_- \) \[1\].

The other possibility corresponds to choose \( L = 0 \) and \( h = 0 \) independently (by preserving both Lorentz and Weyl symmetry) and \( \bar{h}_{++} = 0 \), (choice of particular chiral coupling \( a = 0 \), \( \beta = 1 \)) producing a gravitational anomaly \[5\]

\[ \partial_- T_{++} = \partial^2 \omega_- \equiv \partial^2 \bar{h}_{--} . \] (24)

This is an equation that has no analogue in the case of the Schwinger model. Both situations can be simply described in terms of an anomalous equation with \( \omega_- \) given by eq.(20a) and, in one case \( \bar{h}_{--} \) is the dynamical variable, while in the other case it is \( L - h \).

We have shown that the general treatment of all anomalies of two-dimensional quantum gravity requires \textit{two} arbitrary (regularization dependent) parameters and, therefore, analogy to Schwinger model is achieved only with the particular choice of the second parameter. Within a particular regularization scheme this parameters are fixed and, for example, \( c = 0 \) corresponds to the two-dimensional decomposition of the gravitational coupling in the Feynman rules \( ee_{a\mu} = \delta_{a\mu} + \bar{h}_{a\mu} \), leading to the covariant anomaly \[6\], while \( c = 1 \) is achieved by \( ee_{a\mu} = \delta_{a\mu} + \frac{2P}{2}\delta_{a\mu}h + \bar{h}_{a\mu} + h_{a[\mu]} \) \[11\]. The fixing of the parameter \( a \) is already known through the Schwinger model literature. Spin-connection turns out to be more appropriate as
a gauge field, not only because it bears similarity to the vector gauge potential of
the Schwinger model, but also because it leads to the quantum spin current and
nicely separate local from non-local pieces in the effective action, allowing ,in such
a way, the identification of arbitrary parameters (and reducing their number). It
is still worth reminding that we are working within the second order formulation
and treat spin-connection as a function of zweibein.

Within present formulation gravitational (in terms of $\omega_\mu$), vector an axial (in
terms of $A_\mu$) anomalies of the Schwinger model are treated separately in terms of
independent regularization parameters. It would be interesting to see what kind of
relation among parameters is induced in super-symmetric models of gravity, partic-
cularly an (2, 2) super-gravity where axial gauge field is a member of a gravitational
super-multiplet. Preliminary investigation in this direction, but with different mo-
tivations, has been done both for (0, 1) and (1, 1) super-gravity\textsuperscript{[12]} without giving
answer to the above question since the axial gauge field, in these models, is not yet
a member of the same super-multiplet as the graviton and gravitino. This point is
now under investigation.
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