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The Wess-Zumino-Witten term in non-commutative two-dimensional fermion models

Enrique F. Moreno* **and Fidel A. Schaposnik**†

Departamento de Física, Universidad Nacional de La Plata
C.C. 67, (1900) La Plata, Argentina
E-mail: emoreno@venus.fisica.unlp.edu.ar,
fidel@venus.fisica.unlp.edu.ar

Abstract: We study the effective action associated to the Dirac operator in two dimensional non-commutative Field Theory. Starting from the axial anomaly, we compute the determinant of the Dirac operator and we find that even in the U(1) theory, a Wess-Zumino-Witten like term arises.

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*Investigador CONICET
†Investigador CICBA

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1. Introduction

Interest in non-commutative spaces has been renewed after the discovery that non-commutative gauge theories naturally arise when D-branes with constant B fields are considered [1, 2]. These works as well as that in [3] prompted many investigations both in field theory and in string theory (see references in [3]). Concerning gauge field theories, recent results on chiral and gauge anomalies [4]–[6] have shown that well-known results on "ordinary" models extend naturally and interestingly to the case in which non-commutative spaces are considered. In this work we consider a problem which can be seen as closely related to that of anomalies, namely the evaluation of the two-dimensional fermion determinant in non-commutative space-time. This problem is of interest not only for the analysis of two-dimensional QED and QCD in non-commutative space, but also in connection with abelian and non-abelian bosonization since, as it is well-known, the knowledge of the fermion determinant leads more or less directly to the bosonization rules.

We start by evaluating in section 2 the chiral anomaly in two-dimensional non-commutative space-time in a way adapted to the calculation of fermion determinants through integration of the anomaly. This last is done in section 3 where both the abelian and ($\mathrm{U}(N)$) non-abelian fermion determinant is calculated exactly. In both cases we obtain for the determinant a Wess-Zumino-Witten term. Consequences of our results and possible extensions are discussed in section 4.

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2. The chiral anomaly

2.1 Conventions

As usual, we define the $*$-product between a pair of functions $\phi(x)$, $\chi(x)$ as

$$\begin{aligned}\phi * \chi(x) &\equiv \exp\left(\frac{i}{2}\theta^{\mu\nu}\partial_{x_\mu}\partial_{y_\nu}\right)\phi(x)\chi(y)|_{x=y} \\ &= \phi(x)\chi(x) + \frac{i}{2}\theta^{\mu\nu}\partial_\mu\phi\partial_\nu\chi(x) + O(\theta^2)\,,\end{aligned} \tag{2.1}$$

and the (Moyal) bracket in the form

$$\{\phi,\chi\}(x) \equiv \phi(x) * \chi(x) - \chi(x) * \phi(x)\,, \tag{2.2}$$

so that, when applied to (euclidean) space-time coordinates x^μ, x^ν, one has

$$\{x^\mu, x^\nu\} = i\theta^{\mu\nu}\,, \tag{2.3}$$

which is why one refers to non-commutative spaces. Here $\theta^{\mu\nu}$ is a real, anti-symmetric constant tensor. Since we shall be interested in two dimensional space-time, one necessarily has $\theta^{\mu\nu} = \theta\varepsilon^{\mu\nu}$ with $\varepsilon^{\mu\nu}$ the completely anti-symmetric tensor and θ a real constant. In the context of string theory, non-commutative spaces are believed to be relevant to the quantization of D-branes in background Neveu-Schwarz constant B-field $B_{\mu\nu}$ [1]–[3]. In this context $\theta^{\mu\nu}$ is related to the inverse of $B^{\mu\nu}$. Afterwards, this original interest was extended to the analysis of field theories in non-commutative space and then, as signaled in [6] it becomes relevant to know to what extent old problems and solutions in standard field theory fit in the new non-commutative framework.

A "non-commutative gauge theory" is defined just by using the $*$-product each time the gauge fields have to be multiplied. Then, even in the U(1) abelian case, the curvature $F_{\mu\nu}$ has a non-linear term (with the same origin as the usual commutator in non-abelian gauge theories in ordinary space)

$$\begin{aligned}F_{\mu\nu} &= \partial_\mu A_\nu - \partial_\nu A_\mu - ie\left(A_\mu * A_\nu - A_\nu * A_\mu\right) \\ &= \partial_\mu A_\nu - \partial_\nu A_\mu - ie\{A_\mu, A_\nu\}\,.\end{aligned} \tag{2.4}$$

This field strength is gauge-covariant (not gauge-invariant, even in the abelian case) under gauge transformations which should be represented by U elements of the form

$$U(x) = \exp_*(i\lambda) \equiv 1 + i\lambda - \frac{1}{2}\lambda * \lambda + \cdots \tag{2.5}$$

The covariant derivative implementing infinitesimal gauge transformations takes the form

$$\mathcal{D}_\mu[A]\lambda = \partial_\mu\lambda + ie\left(\lambda * A_\mu - A_\mu * \lambda\right) \tag{2.6}$$

so that an infinitesimal gauge transformation on A_μ reads as usual

$$\delta A_\mu = \frac{1}{e}\mathcal{D}_\mu\lambda\,. \tag{2.7}$$

Concerning finite gauge transformations, one has

$$A_\mu^U = \frac{i}{e}U(x) * \partial_\mu U^{-1}(x) + U(x) * A_\mu * U^{-1}(x)\,. \tag{2.8}$$

Given a fermion field ψ, one can easily see that the combination

$$\gamma^\mu D_\mu[A]\psi = \gamma^\mu\partial_\mu\psi - ie\gamma^\mu A_\mu * \psi \tag{2.9}$$

transforms covariantly under gauge transformations (2.8),

$$\gamma^\mu D_\mu[A^U]\psi^U = U * \gamma^\mu D_\mu[A]\psi\,, \tag{2.10}$$

with

$$\psi^U = U(x) * \psi \tag{2.11}$$

and

$$U(x) * U^{-1}(x) = U^{-1} * U(x) = 1\,. \tag{2.12}$$

A gauge invariant Dirac action can be defined in the form

$$S_f = \int d^dx\ \bar\psi(x) * i\gamma^\mu D_\mu[A]\psi(x)\,. \tag{2.13}$$

2.2 The anomaly

Chiral transformations will be written as

$$\psi'(x) = U_5(x) * \psi\,, \tag{2.14}$$

with

$$U_5(x) = \exp_*(\gamma_5\alpha(x)) = 1 + \gamma_5\alpha + \frac{1}{2}\alpha(x) * \alpha(x) + \cdots \tag{2.15}$$

The chiral anomaly $\mathcal{A}_d$ in d-dimensional space can be calculated from the formula

$$\log J_d[\alpha] = -2\mathbf{A}_d\,, \tag{2.16}$$

$$\mathbf{A}_d = \mathrm{Tr}\,\gamma_5\delta\alpha(x)|_{\mathrm{reg}} \tag{2.17}$$

here $J_d[\alpha]$ is the Fujikawa jacobian associated with an infinitesimal chiral transformation $U = 1 + \gamma_5\delta\alpha$ and Tr includes a matrix and functional space trace.

Let us specialize to the two dimensional case. We shall use the heat-kernel regularization so that (2.17) will be understood as

$$\mathbf{A}_2 = \int d^2x\ \mathcal{A}_2(x) * \delta\alpha(x)\,, \tag{2.18}$$

$$\mathcal{A}_2(x) = \lim_{M\to\infty}\mathrm{Tr}\,\gamma_5\exp_*\left(\frac{\not{D} * \not{D}}{M^2}\right). \tag{2.19}$$

After some standard manipulations, (2.19) takes the form

$$\mathcal{A}_2(x) = \frac{1}{4\pi}\,\mathrm{tr}\,\gamma_5 \not{D} * \not{D} = \frac{1}{4\pi}\,\mathrm{tr}\,(\gamma_5\gamma^\mu\gamma^\nu)\,D_\mu * D_\nu\,. \tag{2.20}$$

Here tr is just the matrix trace. Using $\mathrm{tr}(\gamma_5\gamma^\mu\gamma^\nu) = 2i\,\varepsilon^{\mu\nu}$, eq. (2.20) can be written as

$$\mathcal{A}_2(x) = \frac{e}{2\pi}\varepsilon^{\mu\nu}(\partial_\mu A_\nu - ieA_\mu * A_\nu) = \frac{e}{4\pi}\varepsilon^{\mu\nu}F_{\mu\nu}\,. \tag{2.21}$$

This result coincides with that first obtained in [4].

3. The two-dimensional fermion determinant

Let us write the gauge field in the two-dimensional case in the form

$$\not{A} = \frac{1}{e}\left(i\not{\partial}\exp_*(\gamma_5\phi + i\eta)\right) * \exp_*(-\gamma_5\phi - i\eta)\,. \tag{3.1}$$

Note that in the $\theta_{\mu\nu} \to 0$ limit, eq. (3.1) reduces to the usual decomposition of a two-dimensional gauge field in the form

$$eA_\mu = \varepsilon_{\mu\nu}\partial^\nu\phi + \partial_\mu\eta\,, \tag{3.2}$$

which allows to decouple fermions from the gauge-field and then obtain the fermion determinant as the jacobian associated to this decoupling [8]. Now, the form (3.1) was precisely proposed in [9] to achieve the decoupling in the case of non-abelian gauge field backgrounds, this leading to the calculation of the QCD_2 fermion determinant in a closed form. Afterwards [10], it was shown that writing a two dimensional gauge field as in eq. (3.1) (without the $*$-product but in the U(N) case) does correspond to the choice of a gauge condition . Eq. (3.1) is then the extension of this approach for a case in which non-commutativity arises from the use of the $*$-product.

At the classical level, the change of fermionic variables

$$\begin{aligned}\psi &= \exp_*(\gamma_5\phi + i\eta) * \chi\\ \bar\psi &= \bar\chi * \exp_*(\gamma_5\phi - i\eta)\end{aligned} \tag{3.3}$$

completely decouples the gauge field, written as in (3.1), leading to an action of free massless fermions,

$$S_f = \int d^2x\,\bar\chi * i\not{\partial}\chi \tag{3.4}$$

Of course, this is not the whole story: at the quantum level there is a Fujikawa jacobian J [7] associated to change (3.3). In order to compute this jacobian, we follow the method introduced in [8, 9]. Consider then the change of variables

$$\begin{aligned}\psi &= U_t * \chi_t\,,\\ \bar\psi &= \bar\chi_t * U_t^\dagger\,,\end{aligned} \tag{3.5}$$

where

$$U_t = \exp_* \left(t \left(\gamma_5 \phi + i\eta \right) \right) , \tag{3.6}$$

and t is a real parameter, $0 \leq t \leq 1$. Given the fermion determinant defined as

$$\det(\not\partial - ie \not A) = \int \mathcal{D}\bar{\psi}\mathcal{D}\psi \exp\left(-S_f[\bar{\psi}, \psi]\right) \tag{3.7}$$

we proceed to the change of variables (3.5) which leads to

$$\begin{aligned} \det(\not\partial - ie \not A) &= J[\phi, \eta; t] \int \mathcal{D}\bar{\chi}_t \mathcal{D}\chi_t \exp\left(-S_f[\bar{\chi}_t, \chi_t]\right) \\ &= J[\phi, \eta; t] \det D_t \,, \end{aligned} \tag{3.8}$$

where $J[\phi, \eta; t]$ stands for the jacobian

$$\mathcal{D}\bar{\psi}\mathcal{D}\psi = J[\phi, \eta; t] \mathcal{D}\bar{\chi}_t \mathcal{D}\chi_t \tag{3.9}$$

and we have defined

$$D_t = U_t^\dagger * (\not\partial - ie \not A *) \, U_t \,. \tag{3.10}$$

Now, since the l.h.s. in (3.8) does not depend on t we get, after differentiation,

$$\frac{d}{dt} \log \det D_t = -\frac{d}{dt} \log J[\phi, \eta; t] \tag{3.11}$$

or, after integrating on t and using that $D_0 = \not\partial - ie \not A$ and $D_1 = \not\partial$

$$\det(\not\partial - ie \not A) = \det \not\partial \exp\left(-2 \int_0^1 dt \mathbf{A}_2(t) \right) , \tag{3.12}$$

where we have used

$$\mathbf{A}_2(t) = \frac{d}{dt} \log J[\phi, \eta; t] \,. \tag{3.13}$$

Now, it is trivial to identify $\mathbf{A}_2(t)$ with the two-dimensional chiral anomaly as defined in eq. (2.17), just by writing $\delta\alpha = \phi dt$,

$$\mathbf{A}_2(t) = \left. \mathrm{Tr} \left(\gamma_5 * \phi \right) \right|_{\mathrm{reg}} \,. \tag{3.14}$$

In order to have a gauge-invariant regularization ensuring that the η part of the transformation does not generate a jacobian, we adopt, in agreement with (2.18) and (2.19),

$$\mathbf{A}_2(t) = \lim_{M \to \infty} \mathrm{Tr} \left(\gamma_5 \exp\left(\frac{\not D_t * \not D_t}{M^2} \right) * \phi \right) \tag{3.15}$$

so that finally one has

$$\mathbf{A}_2(t) = \frac{e}{2\pi} \int d^2x \varepsilon^{\mu\nu} \left(\partial_\mu A^t_\nu - ie A^t_\mu * A^t_\nu \right) * \phi = \frac{e}{4\pi} \int d^2x \varepsilon^{\mu\nu} F^t_{\mu\nu} * \phi \,, \tag{3.16}$$

where we have introduced

$$\gamma_\mu A^t_\mu = -\frac{1}{e}\,(i\,\partial\!\!\!/ U_t) * U_t^{-1} \tag{3.17}$$

and analogously for $F^t_{\mu\nu}$. In summary, we can write for the U(1) fermion determinant

$$\det(\partial\!\!\!/ - ie\, A\!\!\!/) = \exp\left(-\frac{e}{2\pi}\int d^2x \int_0^1 dt\; \varepsilon^{\mu\nu}F^t_{\mu\nu} * \phi\right)\det\,\partial\!\!\!/\,. \tag{3.18}$$

It will be convenient to use the relation

$$\gamma^\mu\gamma_5 = -i\varepsilon^{\mu\nu}\gamma_\nu \tag{3.19}$$

to rewrite (3.18) in the form

$$\det(\partial\!\!\!/ - ie\, A\!\!\!/) = \exp\left(\frac{ie}{2\pi}\,\mathrm{tr}\int d^2x\int_0^1 dt\gamma_5\phi * \left(\partial\!\!\!/\, A\!\!\!/^t - ie\, A\!\!\!/^t * A\!\!\!/^t\right)\right)\det\,\partial\!\!\!/\,. \tag{3.20}$$

Then, one can exploit the identity

$$\begin{aligned}\mathrm{tr}\int d^2x\frac{1}{2}\frac{d}{dt}\, A\!\!\!/^t * A\!\!\!/^t &= \frac{1}{e}\,\mathrm{tr}\int d^2x\gamma_5 i\,\partial\!\!\!/\phi * A\!\!\!/^t + 2\,\mathrm{tr}\int d^2x\gamma_5\, A\!\!\!/^t * \phi * A\!\!\!/^t + \\ &+\frac{1}{e}\,\mathrm{tr}\int d^2x(D\!\!\!\!/\,\eta) * A\!\!\!/\end{aligned} \tag{3.21}$$

and find for (3.20)

$$\begin{aligned}\log\det(\partial\!\!\!/ - ie\, A\!\!\!/) &= -\frac{e^2}{4\pi}\,\mathrm{tr}\int d^2x\;\; A\!\!\!/ * A\!\!\!/ + \frac{e^2}{2\pi}\,\mathrm{tr}\int dt\int d^2x\;\gamma_5\phi * A\!\!\!/ * A\!\!\!/ + \\ &+\frac{e}{2\pi}\int dt\int d^2x\;(D\!\!\!\!/\,\eta) * A\!\!\!/ + \log\det\,\partial\!\!\!/\,.\end{aligned} \tag{3.22}$$

This is the final form for the fermion determinant in a U(1) gauge theory. In order to write it in a more suggestive way connecting it with the Wess-Zumino-Witten term, let us consider the light-cone gauge $A_+ = 0$, Then, one can see after some algebra that [11]

$$\begin{aligned}\log\left(\frac{\det(\partial\!\!\!/ - ie\, A\!\!\!/)}{\det\,\partial\!\!\!/}\right) &= -\frac{1}{8\pi}\int d^2x\left(\partial_\mu g(x)^{-1}\right) * \left(\partial_\mu g(x)\right) + \\ &+\frac{i}{12\pi}\epsilon_{ijk}\int_B d^3y\, g(x,t)^{-1} * \left(\partial_i g(x,t)\right) * \\ &\qquad * g(x,t)^{-1} * \left(\partial_j g(x,t)\right) g^{-1} * \left(\partial_k g(x,t)\right)\end{aligned} \tag{3.23}$$

here we have written $A_- = (i/e)g(x) * \partial_- g^{-1}(x)$ with $g(x) = \exp_*(2\phi(x))$, $g(x,t) = \exp_*(2\phi(x)t)$ and $d^3y = d^2xdt$ so that the integral in the second line of eq. (3.24) runs over the three dimensional manifold B, which in compactified euclidean space can be identified with a ball with boundary S^2. Index i runs from 1 to 3. As in the ordinary commutative case, because the determinant was computed in euclidean space, elements g should be considered as belonging to $\mathrm{U}(1)_C$ (the complexified U(1)) [11, 12].

So, we have found for the two-dimensional non-commutative fermion determinant that, even for a U(1) gauge field background, a Wess-Zumino-Witten term arises due to non-commutativity of the $*$-product. Of course, in the $\theta^{\mu\nu} \to 0$ limit in which the $*$-product becomes the ordinary one, the U(1) fermion determinant contribution to the gauge field effective action reduces to $(-1/2\pi) \int d^2x \phi \partial^\mu \partial_\mu \phi$ which is nothing but the Schwinger determinant result expressed in a gauge-invariant way.

The method we have employed has the advantage that it can be trivially generalized to the case of a U(N) gauge group. One has just to take into account that in (3.1) one has

$$\phi = \phi^a t^a \,, \qquad \eta = \eta^a t^a \,, \tag{3.24}$$

with t^a the U(N) generators. Then, as originally shown in [9] for the commutative case, the fermion determinant can be seen to be given by

$$\det(\partial\!\!\!/ - ie A\!\!\!/) = \exp\left(-\frac{e}{4\pi}\,\mathrm{tr}^c \int d^2x \int_0^1 dt \varepsilon^{\mu\nu} F^t_{\mu\nu} * \phi\right) \det \partial\!\!\!/ \,, \tag{3.25}$$

where tr^c is a trace over the U(N) algebra. Then, following the same steps leading to (3.24), one gets, in the U(N) case

$$\begin{aligned}\log\left(\frac{\det(\partial\!\!\!/ - ie A\!\!\!/)}{\det \partial\!\!\!/}\right) &= -\frac{1}{8\pi}\,\mathrm{tr}^c \int d^2x \left(\partial_\mu g^{-1}\right) * \left(\partial_\mu g\right) + \\ &\quad + \frac{i}{12\pi}\epsilon_{ijk}\,\mathrm{tr}^c \int_B d^3y g^{-1} * (\partial_i g) * g^{-1} * (\partial_j g)\, g^{-1} * (\partial_k g) \,,\end{aligned} \tag{3.26}$$

where again, in the light-cone gauge we have written

$$A_- = -\frac{i}{e} g * \partial_- g^{-1} \,, \qquad A_+ = 0 \tag{3.27}$$

$$g = \exp_*(2\phi^a t^a) \,. \tag{3.28}$$

Eq. (3.27) is the generalization of the expression given in [13] for the two-dimensional non-abelian fermion determinant to the case of non commutative space-time.

4. Conclusion

We studied in this article the effective action of the gauge degrees of freedom in a two dimensional non-commutative Field Theory of fermions coupled to a gauge field. Using Fujikawa's approach, we computed the chiral anomaly and, from it, the fermionic determinant of the non-commutative Dirac operator.

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As it was to be expected, the result for the fermion determinant corresponds to the $*$-deformation of the standard result. Now, the fact that a Moyal bracket enters in the field strength curvature even in the abelian case, has important consequences, some of which have already been signaled in [4]–[6] where chiral and gauge anomalies in non-commutative spaces have been analyzed.

In our framework, where the anomaly was integrated in order to obtain the fermion determinant, this reflects in the fact that a Wess-Zumino-Witten like term arises *both* in the abelian and in the non-abelian cases (eqs. (3.24) and (3.27) respectively). This should have, necessarily, implications in relevant aspects of two-dimensional theories since, as it is well-known, bosonization is closely related to the form of the fermion determinant [15]. Indeed, the bosonization rules for fermion currents as well as the resulting current algebra can be easily derived by differentiation of the Dirac operator determinant $\det(\not{d} - i\not{s})$ with respect to the source s_μ (see [16] for a review). Now, as one learns from ordinary non-abelian bosonization, where the Polyakov-Wiegmann identity plays a central rôle in the bosonization recipe, here one should have an analogous identity which will lead to non-trivial changes at the level of currents and, a fortiori, for the current algebra. In view of the relevance of these objects in connection with two-dimensional bosonic and fermionic models, it will be worthwhile to pursue the investigation initiated here in this direction.

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