Lattice chiral symmetry, Yukawa couplings and the Majorana condition

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Abstract

It is shown that the conflict between lattice chiral symmetry and the Majorana condition in the presence of Yukawa couplings, which was noted in our previous paper, is related in an essential way to the basic properties of Ginsparg-Wilson operators, namely, locality and species doubling.

1 Introduction

Recent developments in the treatment of lattice fermions paved a way to deal with lattice chiral symmetry in a unified manner\[1]-[6]. In a recent paper, we pointed out that the otherwise successful lattice chiral symmetry has a certain conflict with the definition of the Majorana fermion in the presence of Yukawa couplings\[7]. This issue was discussed in connection with the lattice regularization of the simplest supersymmetric theory\[8]-[12], namely, the Wess-Zumino model and its non-renormalization theorem\[13]-[15]. We consider that a consistent formulation of Majorana fermions is a prerequisite for a precise analysis of supersymmetry and its breaking, if one adopts the Ginsparg-Wilson operator as a basic building block.

In this letter, we further clarify this conflict of lattice chiral symmetry and the Majorana condition in the presence of Yukawa couplings. Our basic observation is the transformation of a chiral symmetric lattice theory with Yukawa couplings, which is defined in the manner of Niedermayer\[3], Narayanan\[16] and Chandrasekharan\[17], to a theory which is a generalization of the model noted by Lüscher\[4] by a singular field re-definition. By this way we can understand the origin of the conflict between the lattice chiral symmetry and the Majorana condition from a different view point. Our analysis indicates that the above conflict is related in an essential and subtle way to the basic issues of lattice chiral symmetry, namely, locality and species doubling.

In our analysis, we use a general class of Ginsparg-Wilson operators and our analysis below is valid for all these operators. The lattice Dirac operator $D$ is defined by the algebraic relation\[18]

$$\gamma_5 (\gamma_5 D) + (\gamma_5 D) \gamma_5 = 2a^{2k+1}(\gamma_5 D)^{2k+2}$$

(1.1)
where the parameter $a$ is the lattice spacing; $k$ stands for non-negative integers, and $k = 0$ corresponds to the conventional Ginsparg-Wilson relation\[6\]. When one defines a hermitian operator $H$ by
\[ H = a \gamma_5 D = H^\dagger = aD^\dagger \gamma_5 \] (1.2)
the above algebraic relation is written as
\[ \gamma_5 H + H \gamma_5 = 2H^{2k+2}. \] (1.3)
We can also show
\[ \gamma_5 H^2 = (\gamma_5 H + H \gamma_5)H - H(\gamma_5 H + H \gamma_5) + H^2 \gamma_5 = H^2 \gamma_5 \] (1.4)
which implies
\[ H^2 = a^2 D^\dagger D = \gamma_5 H^2 \gamma_5 = a^2 DD^\dagger. \] (1.5)
When we define
\[ \Gamma_5 \equiv \gamma_5 - H^{2k+1}, \] \[ \hat{\gamma}_5 \equiv \gamma_5 - 2H^{2k+1}, \] (1.6)
the defining algebra (1.1) is written as
\[ \gamma_5 H + H \hat{\gamma}_5 = 0 \] (1.7)
or $\Gamma_5 H + H \Gamma_5 = 0$, and $(\hat{\gamma}_5)^2 = 1$. We can also show the relation
\[ \gamma_5 \Gamma_5 + \Gamma_5 \gamma_5 = 2\Gamma_5^2 = 2(1 - H^{4k+2}) \] (1.8)
which implies $H^2 \leq 1$. We next note\[6\]
\[ D = P_+ \hat{D}_+ + P_- \hat{D}_- . \] (1.9)
Here we defined two projection operators
\[ P_{\pm} = \frac{1}{2}(1 \pm \gamma_5), \] \[ \hat{P}_{\pm} = \frac{1}{2}(1 \pm \hat{\gamma}_5) \] (1.10)
which satisfy the relations
\[ P_+ \hat{P}_+ = P_+ \gamma_5 \Gamma_5, \] \[ P_- \hat{P}_- = P_- \gamma_5 \Gamma_5. \] (1.11)
We then define the chiral components\[6][10]\[16\]
\[ \bar{\psi}_{L,R} = \bar{\psi} P_{\pm}, \quad \psi_{R,L} = \hat{P}_{\pm} \psi \] (1.12)
and the scalar and pseudoscalar densities by\[17\]
\[ S(x) = \bar{\psi}_L \psi_R + \bar{\psi}_R \psi_L = \bar{\psi}_5 \gamma_5 \psi, \] \[ P(x) = \bar{\psi}_L \psi_R - \bar{\psi}_R \psi_L = \bar{\psi}_5 \psi. \] (1.13)
2 Yukawa couplings and the Majorana condition

The most natural Lagrangian consistent with lattice chiral symmetry $\delta \psi = i\epsilon \gamma_5 \psi$, $\delta \bar{\psi} = \bar{\psi} i \gamma_5$ and $\delta \phi = -2i\phi$, which is softly broken by the mass term, is defined by\[1 9\][20][21]

\[
\mathcal{L} = \bar{\psi} D\psi + m\bar{\psi} \gamma_5 \Gamma_5 \psi + 2g \bar{\psi}(P_+\phi \hat{\phi}^+ + P_-\phi^\dagger \hat{\phi}^-)\psi
= \bar{\psi}_R D\psi_R + \bar{\psi}_L D\psi_L + m[\bar{\psi}_R \psi_L + \bar{\psi}_L \psi_R]
+ 2g[\bar{\psi}_L \phi \psi_R + \bar{\psi}_R \phi^\dagger \psi_L].
\] (2.1)

We fixed the mass term in such a way that it is generated by a shift $\phi(x) \rightarrow \phi(x) + m/(2g)$ in $\phi(x) = (A(x) + iB(x))/\sqrt{2}$ in the interaction terms; we adopt this procedure in the following. The fermion mass term is then defined by the scalar density formed of a fermion bi-linear (1.13).

It has been shown elsewhere\[7\] that the above Lagrangian (2.1) has a difficulty in performing the Majorana reduction, and thus the Majorana fermion is not defined in a manner consistent with lattice chiral symmetry: When one defines\[19\][20][21]

\[
\psi = (\chi + i\eta)/\sqrt{2},
\bar{\psi} = (\chi^T C - i\eta^T C)/\sqrt{2}
\] (2.2)
in $\mathcal{L}$ (2.1), one naively expects\[5\]by noting $P_+\phi \hat{\phi}^+ + P_-\phi^\dagger \hat{\phi}^- = \frac{1}{\sqrt{2}}(A\gamma_5 \Gamma_5 + iB\Gamma_5)$,

\[
\mathcal{L} = \frac{1}{2} \chi^T C D\chi + \frac{1}{2} m\chi^T C \gamma_5 \Gamma_5 \chi + g \frac{1}{\sqrt{2}} \chi^T C (A\gamma_5 \Gamma_5 + iB\Gamma_5) \chi
+ \frac{1}{2} \eta^T C D\eta + \frac{1}{2} m\eta^T C \gamma_5 \Gamma_5 \eta + g \frac{1}{\sqrt{2}} \eta^T C (A\gamma_5 \Gamma_5 + iB\Gamma_5) \eta
\] (2.3)

but this actually fails, since $(C\Gamma_5)^T \neq -C\Gamma_5$, where $C$ stands for the charge conjugation matrix, and the non-commuting property of the difference operator $[\gamma_5 \Gamma_5, A(x)] \neq 0$ though $(C\gamma_5 \Gamma_5)^T = -C\gamma_5 \Gamma_5$ or equivalently $C\gamma_5 \Gamma_5 C^{-1} = (\gamma_5 \Gamma_5)^T$ in the Yukawa couplings. To be precise, we need to perform a charge conjugation operation of the gauge field to satisfy $C\gamma_5 \Gamma_5 C^{-1} = (\gamma_5 \Gamma_5)^T$, for example, in the presence of the background gauge field; this extra operation of charge conjugation is implicitly assumed in the following.

We now observe that the field re-definition\[6\]

\[
\psi' = \gamma_5 \Gamma_5 \bar{\psi},
\bar{\psi}' = \bar{\psi}
\] (2.4)
in the above Lagrangian gives rise to the Lagrangian

\[
\mathcal{L}' = \bar{\psi}' D\psi' \frac{1}{\gamma_5 \Gamma_5} \psi' + m\bar{\psi}' \psi' + 2g \bar{\psi}'(P_+\phi \hat{\phi}^+ + P_-\phi^\dagger \hat{\phi}^-)\psi'
\] (2.5)
where we used the relations (1.11). This shows that the theory defined by the Lagrangian invariant under the lattice chiral symmetry

\[ Z = \int D\psi D\bar{\psi} \exp[\int \mathcal{L} d^4x] \] (2.6)

is related to the theory defined by the transformed Lagrangian as

\[ Z = (\det \gamma_5 \Gamma_5) Z' \equiv (\det \gamma_5 \Gamma_5) \int D\psi' D\bar{\psi}' \exp[\int \mathcal{L}' d^4x]. \] (2.7)

This new Lagrangian (2.5) corresponds to a generalization of the Lagrangian considered by Lüscher if one eliminates the auxiliary field. To be specific, Lüscher considered the Lagrangian\[4]\n
\[ \mathcal{L}_L = \bar{\psi} D\psi - \frac{1}{a} \bar{\chi} \chi + 2g(\bar{\psi} + \bar{\chi})(P_+ \phi P_+ + P_- \phi^\dagger P_-)(\psi + \chi) \] (2.8)

which is shown to be invariant under a modified chiral transformation, if one assumes that \( D \) satisfies the Ginsparg-Wilson relation (1.1) with \( k = 0 \), namely, the overlap operator\[2]. By considering the re-definition of field variables

\[ \psi' = (\psi + \chi), \quad \bar{\psi}' = (\bar{\psi} + \bar{\chi}), \quad \Psi = (\psi - \chi), \quad \bar{\Psi} = (\bar{\psi} - \bar{\chi}) \] (2.9)

one obtains\[4] after the path integral over \( \Psi \) and \( \bar{\Psi} \)

\[ \int D\bar{\Psi} D\Psi D\bar{\chi} D\chi \exp[\int d^4x \mathcal{L}_L] = \det(1 - aD) \int D\bar{\psi}' D\psi' \exp[\int d^4x \mathcal{L}'_L] \] (2.10)

where

\[ \mathcal{L}'_L = \bar{\psi}' D(1 - aD) \psi' + 2g\bar{\psi}'(P_+ \phi P_+ + P_- \phi^\dagger P_-)\psi'. \] (2.11)

This final expression (2.11) corresponds to our Lagrangian (2.5) for \( k = 0 \) (i.e., the standard overlap operator for which \( \gamma_5 \Gamma_5 = 1 - aD \)), up to the chiral symmetry breaking mass term.

The above transformation (2.4) is singular if \( \gamma_5 \Gamma_5 = 1 - \gamma_5 HH^{2k} = 0 \). Since \( \Gamma_5^2 = 1 - H^{4k+2} \), the necessary condition for the appearance of the singularity is \( H^2 = 1 \). This condition is analyzed in more detail as follows: \( H^2 \) is found from an explicit form of \( H \), which is local and free of species doubling\[23][25],

\[ H(ap_\mu) = \frac{1}{2} \sqrt{\frac{H_W^2}{M_k}} \left( \frac{1}{\sqrt{H_W^2}} + M_k \right)^{\frac{k+1}{4}} \left( (\sqrt{H_W^2} - M_k) \frac{1}{\sqrt{H_W^2}} + \sqrt{H_W^2} - M_k \right) \]

where \( \gamma' = \gamma^\mu \sin ap_\mu \) with anti-hermitian \( \gamma^\mu \) and

\[ M_k(p) = \left( \frac{r}{a} \sum_\mu (1 - \cos ap_\mu) \right)^{2k+1} - \left( \frac{m_0}{a} \right)^{2k+1}, \]

\[ H_W^2 = \left( \sum_\mu \frac{1}{a^2} \sin^2 ap_\mu \right)^{2k+1} + (M_k(p))^2. \] (2.12)

\[3\] This was discussed in a different context in Ref.[\[23\]].
We thus obtain

\[ H^2(ap) = \left( \frac{1}{2 \sqrt{H_W^2}} \right)^{\frac{2k+2}{2k+1}} \left( (\sqrt{H_W^2} + M_k)^{\frac{2k+2}{2k+1}} + (\sqrt{H_W^2} - M_k)^{\frac{2k+2}{2k+1}} \right)^2. \] (2.13)

Using the relation \( s^2/a^2 = (H_W^2 - M_k^2)^{\frac{1}{2k+1}} \), we find that \( H^2(ap) = 1 \) implies \( \sqrt{H_W^2} = M_k \).

This last condition is written explicitly as

\[ \sqrt{\left( \sum_{\mu} \frac{1}{a^2} \sin^2 ap_\mu \right)^{2k+1} + (M_k(p))^2} = M_k(p). \] (2.14)

Since \( m_0 \) is constrained by \( 0 < m_0 < 2r \) to avoid the appearance of species doublers, we have \( M_k < 0 \) for a physical mode (\( p_\mu = (0,0,0,0) \)) and \( M_k > 0 \) for doubler modes (\( p_\mu = (\pi/a, 0, 0, 0) \), etc). Therefore the above equation (2.14) holds only for would-be doubler modes. Just on top of the doubler modes, one can also confirm

\[ \gamma_5 \Gamma_5 = 1 - \gamma_5 H H^{2k} = 0. \] (2.15)

Thus the singularity of the transformation comes from the momentum regions of would-be species doublers. It is interesting to notice that when \( m_0 > 8r \) (where all the doubler modes appear as massless modes), the singularity of the above transformation disappears.

The above transformation (2.4) thus induces singularities inside the Brillouin zone and thus spoils the locality of the Dirac operator \( D' \equiv D/(\gamma_5 \Gamma_5) \). The ordinary formulation (2.1), which is local, and the model (2.8), which appears to be local but actually non-local, formally give rise to the same path integral as in (2.7) and (2.10), but this equivalence does not hold in a strict sense since one has to go through the singular Lagrangian such as (2.5) in the intermediate stage.

The interesting property of this field re-definition (2.4) in the context of the present analysis of the Majorana condition is that the transformed singular theory (2.5) is invariant under the naive continuum chiral symmetry and that it allows the Majorana reduction. This fact is understood by noting the relation

\[ \gamma_5 \Gamma_5 \hat{\gamma}_5 = \gamma_5 (\gamma_5 \Gamma_5 + \Gamma_5 \gamma_5) - \gamma_5 \Gamma_5 \gamma_5 = \gamma_5 (\gamma_5 \Gamma_5) \] (2.16)

where we used (1.8). This relation shows that the chiral transformation of \( \psi \) generated by \( \hat{\gamma}_5 \) of the regular theory is related to the chiral transformation of \( \psi' \) generated by the continuum \( \gamma_5 \) as

\[ \delta \psi' \equiv (\gamma_5 \Gamma_5) \delta \psi = \gamma_5 \Gamma_5 i \epsilon \hat{\gamma}_5 \psi = i \epsilon \gamma_5 (\gamma_5 \Gamma_5) \psi = i \epsilon \gamma_5 \psi' \] (2.17)

and of course \( \delta \bar{\psi}' = \delta \bar{\psi} = \bar{\psi}' i \epsilon \gamma_5 \). As for an analysis of the Majorana reduction of this transformed singular Lagrangian, we note that

\[ (CD \frac{1}{\gamma_5 \Gamma_5})^T = (CD \frac{1}{C \gamma_5 \Gamma_5} C)^T = C^T (\frac{1}{C \gamma_5 \Gamma_5})^T (CD)^T = -C \frac{1}{C \gamma_5 \Gamma_5} CD = -C \frac{1}{\gamma_5 \Gamma_5} D = -CD \frac{1}{\gamma_5 \Gamma_5} \] (2.18)
where $C$ is the charge conjugation matrix and we used the properties
\[ C^T = -C, \]
\[ (C\gamma_5\Gamma_5)^T = -C\gamma_5\Gamma_5, \]
\[ (CD)^T = -CD \] (2.19)
and the relation $D\gamma_5\Gamma_5 = D(1 - \gamma_5HHH^2k) = \gamma_5\Gamma_5D$ by noting $DH^2 = H^2D$ which follows from (1.5). Our operator thus satisfies the condition necessary for the Majorana reduction, if one ignores the singularity in $CD/(\gamma_5\Gamma_5)$.

When one performs the field transformation (2.2) (with $\psi$ replaced by $\psi'$) in $L'$ (2.5) by noting $(C\gamma_5\Gamma_5)^T = -C\gamma_5\Gamma_5$, one can thus define the Majorana fermion in a formal sense by
\[ L'_{\text{Majorana}} = \frac{1}{2} \chi^T CD \frac{1}{\gamma_5\Gamma_5} \chi + \frac{1}{2} m\chi^T C\chi + g\chi^T C(P_+\phi P_+ + P_-\phi^\dagger P_-)\chi. \] (2.20)
This formulation of the Majorana fermion and the resulting Pfaffian, if one does not care about the singularity, gives rise to the same result as in our previous paper[7] which utilized $\sqrt{Z}$, on the basis of the relation
\[ \sqrt{(\det \gamma_5\Gamma_5)Z'} = \sqrt{Z} \] (2.21)
in a formal perturbation theory, for example[4].

If one tentatively adopts the singular Lagrangian (2.20), a lattice version of the Wess-Zumino model in our previous paper[7] is re-written as (after a rescaling of the auxiliary field $F$)
\[ L_{WZ} = \frac{1}{2} \chi^T C \frac{1}{\Gamma_5} \gamma_5 D\chi + \frac{1}{2} m\chi^T C\chi + g\chi^T C(P_+\phi P_+ + P_-\phi^\dagger P_-)\chi \]
\[ -\phi^\dagger D^\dagger D\phi + F^\dagger \frac{1}{\Gamma_5^2} F + m[F\phi + (F\phi)^\dagger] + g[F\phi^2 + (F\phi^2)^\dagger]. \] (2.22)
Note that $((1/\Gamma_5)\gamma_5 D)^2 = -D^\dagger D(1/\Gamma_5^2)$, and thus the kinetic (Kähler) terms satisfy a necessary condition for supersymmetry provided that one ignores the $4 \times 4$ unit matrix in $D^\dagger D$ and $1/\Gamma_5^2$ in bosonic terms. The (super-)potential parts of this Lagrangian (2.22) are identical to those of the continuum theory. This representation of the Lagrangian, when treated with due qualifications, is thus useful to understand the symmetry aspects of the model. But the formulation of this Lagrangian is not satisfactory since it does not give a uniform wave function renormalization factor in the one-loop level of perturbation theory[7], besides the issues related to the Leibniz rule[8].

Alternatively, one may consider the symmetric definitions[26] of left and right components by $\psi_{R,L} = (1 \pm \Gamma_5/\Gamma)\psi/2$ and $\bar{\psi}_{R,L} = \bar{\psi}(1 \mp \gamma_5\Gamma_5\gamma_5/\Gamma)/2$, where $\Gamma = \sqrt{1 - H^{4k+2}}$.

4 In the non-perturbative sense, the relation (2.21) stands for something like $0 \times (1/0) = 1$ if one takes possible zero modes in $\gamma_5\Gamma_5$ into account.
respectively. In this case, the second expression of (2.1) is written as
\[
\mathcal{L} = \bar{\psi} D \psi + m \bar{\psi} \gamma_5 \Gamma_5 \psi + \frac{g}{\sqrt{2}} \bar{\psi} \left[ A + (\gamma_5 \Gamma_5 \gamma_5 / \Gamma) A(\Gamma_5 / \Gamma) + i(\gamma_5 \Gamma_5 \gamma_5 / \Gamma) B + iB(\Gamma_5 / \Gamma) \right] \psi \tag{2.23}
\]
and the Majorana fermion can be defined. But the modified chiral operators, $\Gamma_5 / \Gamma$ and $\gamma_5 \Gamma_5 \gamma_5 / \Gamma$, are ill-defined at $H^2 = 1$, namely
\[
\Gamma_5 / \Gamma \simeq \gamma_5 (\gamma^\mu \sin a p_\mu / a) / \sqrt{\sum_\mu \sin^2 a p_\mu / a^2} \tag{2.24}
\]
for $H^2 \simeq 1$ by noting (2.14), and Yukawa couplings become singular. Incidentally, an analysis of the Euclidean Majorana condition is related to that of CP symmetry.[7][26]

We summarize our analysis as follows:
1. The most natural formulation consistent with lattice chiral symmetry (2.1), which is successful in QCD, does not accommodate the Euclidean Majorana fermion.[7]
2. If one allows a non-local singular Lagrangian such as (2.5), (2.11) and (2.23) (or if one allows species doubling by choosing $m_0 > 8r$), one can accommodate the Euclidean Majorana fermion and the resulting Pfaffian. The non-locality is however expected to become serious in the presence of background gauge field[23].
3. Our analysis is valid for a general class of Ginsparg-Wilson operators (1.1) and thus exhibits generic properties of these lattice Dirac operators.

In conclusion, a deeper understanding of Ginsparg-Wilson operators is required to incorporate Majorana fermions with Yukawa couplings in a manner consistent with lattice chiral symmetry.

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