Locality of the fourth root of the Staggered-fermion
determinant: renormalization-group approach

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ABSTRACT

Consistency of present-day lattice QCD simulations with dynamical
("sea") staggered fermions requires that the determinant of the staggered-
fermion Dirac operator, det(D), be equal to det^4(D_{rg}) det(T) where D_{rg}
is a local one-flavor lattice Dirac operator, and T is a local operator con-
taining only excitations with masses of the order of the cutoff. Using
renormalization-group (RG) block transformations I show that, in the
limit of infinitely many RG steps, the required decomposition exists for
the free staggered operator in the “flavor representation.” The resulting
one-flavor Dirac operator D_{rg} satisfies the Ginsparg-Wilson relation in the
massless case. I discuss the generalization of this result to the interacting
theory.
1. Introduction

Lattice QCD simulations with dynamical (“sea”) staggered fermions [1] are providing predictions for hadronic observables with unprecedented accuracy [2]. In these numerical calculations, it is crucial that all sources of error be under systematic control. This raises the question of the validity of the “fourth root trick” used in these simulations.

Let me briefly explain the problem (more details may be found e.g. in ref. [3]). In four dimensions, the staggered Dirac operator $D$ is a one-component lattice operator which, in the free-field case, has sixteen poles in the Brillouin zone. These poles combine into four Dirac fermions (with a total of sixteen degrees of freedom) in the continuum limit. To account for three dynamical quarks – up, down and strange – the Boltzmann weight used for generating the dynamical configurations involves the factor

$$
\det^{1/4}(D(m_u)) \det^{1/4}(D(m_d)) \det^{1/4}(D(m_s)) .
$$

Taking the fourth root of each staggered-fermion determinant ensures that the lattice theory describes three (and not twelve) quarks in the continuum limit. While the “fourth root trick” is necessary in practice in order to reach the desired continuum theory, it is not obvious that this trick is consistent. The question is whether the gauge-field configurations generated with this Boltzmann weight correspond to a local lattice theory. If one could show that $\det(D) = \det^4(D_{rg})$, this would provide a positive answer.

In fact, a far weaker condition is sufficient to guarantee locality, and, hence, consistency.\(^3\) Suppose one can show that

$$
\det(D) = \det^4(D_{rg}) \det(\mathcal{T}) .
$$

Here $D_{rg}$ is a local one-flavor lattice Dirac operator, and $\mathcal{T}$ is a local operator containing only excitations with masses of the order of the cutoff. We may now write $\det(\mathcal{T}) = \exp(-S_{\text{eff}}(U))$. Since $\mathcal{T}$ contains only excitations with cutoff masses, we expect that the effective action $S_{\text{eff}}(U)$ is local; trivially, the same is true for $(1/4)S_{\text{eff}}(U)$.

By eq. (2), the fourth-root trick then amounts to using dynamical fermions with the local Dirac operator $D_{rg}$, together with the modification of the gauge-field action by $(1/4)S_{\text{eff}}(U)$.

The natural framework to realize relation (2) is through RG block transformations. After introducing the relevant concepts in Sect. 2, the free staggered-fermion

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\(^1\)In practice, $m_u = m_d$ in the simulations. Note also that the simulations correspond to a “hybrid” and/or partially-quenched theory, since the “sea” and the “valence” quarks may differ in their masses, as well as in the details of the discretization used for each.

\(^2\)In this paper, a local operator means an operator whose kernel decays exponentially with the separation $|x - y|$, with a decay rate which is $O(1)$ in lattice units. A similar notion of locality applies to the effective action $S_{\text{eff}}(U)$ discussed below Eq. (2).

\(^3\)This observation was recently made in ref. [4].

\(^4\)Locality of $\det^{1/4}(\mathcal{T})$ in the free theory is addressed in Sect. 3.
operator is dealt with in Sect. 3, which constitutes the main part of this work. Using
the “flavor representation” of staggered fermions [5] it is shown that the decompo-
sition (2) is realized in the limit of infinitely many RG blocking steps. Central to
this discussion is a theorem on the locality of RG-blocked Wilson fermions proved
in ref. [6]; only a trivial amendment is needed in order to generalize the theorem to
the flavor representation of staggered fermions. The limiting operator is constructed
explicitly. (For Wilson fermions, see ref. [7]. See also ref. [8] for an RG treatment of
staggered fermions within the one-component formalism.) Like any fixed-point oper-
ator [9], the limiting operator satisfies the Ginsparg-Wilson (GW) relation [10, 11].

Sect. 4 contains a discussion of some of the issues (both theoretic al and practical)
that arise in the interacting theory.

2. Renormalization-group transformations

We first consider a general setup, following refs. [12, 6]. Starting from a bilinear
fermion action with Dirac operator \( D_0 \), an RG blocking is introduced via the identities

\[
Z = \int d\psi d\bar{\psi} \exp(-\bar{\psi} D_0 \psi) \\
= \int d\psi d\bar{\psi} d\chi d\bar{\chi} \exp(-\bar{\psi} D_0 \psi - \alpha(\bar{\chi} - \bar{\psi} Q^\dagger)(\chi - Q \psi)) \\
= \int d\psi d\bar{\psi} d\chi d\bar{\chi} d\eta d\bar{\eta} \exp(-\bar{\psi} D_0 \psi + \alpha^{-1} \bar{\eta} \eta + (\bar{\chi} - \bar{\psi} Q^\dagger) \eta + \bar{\eta}(\chi - Q \psi)) \\
= \det^{-1}(G_1) \int d\chi d\bar{\chi} \exp(-\bar{\chi} D_1 \chi).
\]

Here \( \psi, \bar{\psi} \) live on the “fine” lattice with spacing \( a_0 \), whereas \( \chi, \bar{\chi} \) (and the auxiliar y
field \( \eta, \bar{\eta} \)) live on the “coarse” lattice with spacing \( a_1 \) equal to an integer
multiple of \( a_0 \). In this paper we set \( a_1 = 2a_0 \). The RG blocking kernel \( Q \) is a rectangular matrix
satisfying

\[
QQ^\dagger = cI,
\]

where \( I \) is the identity matrix on the coarse lattice. Explicitly,

\[
G_1 = \Gamma_1 = (D_0 + \alpha Q^\dagger Q)^{-1}, \\
D_1 = \alpha - \alpha^2 G_1 Q^\dagger, \\
D_1^{-1} = \alpha^{-1} + QD_0^{-1} Q^\dagger.
\]

Eqs. (5a) and (5b) are obtained by integrating over \( \psi, \bar{\psi} \) in Eq. (3b), and Eq. (5c)
by first integrating over \( \psi, \bar{\psi} \) and then over \( \eta, \bar{\eta} \) in Eq. (3c). Iterating the blocking
transformation we have

\[
\Gamma_j = (D_{j-1} + \alpha Q^{(j)\dagger} Q^{(j)})^{-1}, \\
D_j = \alpha - \alpha^2 Q^{(j)} \Gamma_j Q^{(j)\dagger}, \\
D_j^{-1} = \alpha^{-1} + Q^{(j)} D_{j-1}^{-1} Q^{(j)\dagger}.
\]
where $Q^{(j)}$ denotes the blocking kernel from the $(j-1)$-th lattice (with spacing $2^{j-1}a_0$) to the $j$-th lattice (with spacing $2^ja_0$). We may also go directly from the finest to the coarsest lattice:

$$G_n = (D_0 + \alpha_n Q_n^1 Q_n)^{-1}, \quad (7a)$$
$$D_n = \alpha_n - \alpha_n^2 Q_n G_n Q_n^1, \quad \alpha_n = \frac{1 - c}{1 - c^n}; \quad (7b)$$
$$D_n^{-1} = \alpha_n^{-1} + Q_n D_0^{-1} Q_n^1, \quad (7c)$$

where $Q_n = Q^{(n)}Q^{(n-1)}\cdots Q^{(1)}$. Eq. 7c follows by iterating Eq. 6c while using Eq. 3. The other equations follow by noting that the product of $n$ blocking transformations can also be represented as a single “big” blocking transformation as in Eq. 4, provided we let $\chi, \overline{\chi}$ live on the coarsest lattice, and we make the replacements $\alpha \rightarrow \alpha_n, Q \rightarrow Q_n$. Hence each relation in Eq. 7 must match the corresponding one in Eq. 4.

3. Free staggered fermions

We now turn to free staggered fermions. In the flavor representation the staggered Dirac operator has four-component spin and flavor indices, and is given explicitly by

$$D_0 = a^{-1} \sum_\mu ((\gamma_\mu \otimes I) \nabla_\mu + (\gamma_5 \otimes \tau_\mu) \Delta_\mu) + m, \quad (8)$$

where $a$ is the lattice spacing, $\nabla_\mu f(x) = (f(x + \hat{\mu}) - f(x - \hat{\mu}))/2$, and $\Delta_\mu f(x) = (2f(x) - f(x + \hat{\mu}) - f(x - \hat{\mu}))/2$. The usual Dirac matrices are $\gamma_\mu$, while the $\tau_\mu$ constitute another set of Dirac matrices acting on the flavor index. (Taken together, $(\gamma_\mu \otimes I)$ and $(\gamma_5 \otimes \tau_\mu)$ form a representation of the eight-dimensional Dirac algebra.) For $m = 0$, $D_0$ is anti-hermitian. Going to momentum space one has $D_0^{-1} = \Omega^{-1}D_0^\dagger$ where

$$D_0 = a^{-1} \sum_\mu ((\gamma_\mu \otimes I) i \sin(p_\mu a) + (\gamma_5 \otimes \tau_\mu)(1 - \cos(p_\mu a))) + m,$$
$$\Omega = D_0^\dagger D_0 = a^{-2} \sum_\mu (\sin^2(p_\mu a) + (1 - \cos(p_\mu a))^2) + m^2. \quad (9)$$

We will apply $n$ block transformations to the Dirac operator 8. We set $m = 0$, hence we may write

$$D_0^{-1} = D_0^{-1}(p; a) = -\sum_\mu (i(\gamma_\mu \otimes I) A_{\mu}^0(p; a) + (\gamma_5 \otimes \tau_\mu) B_{\mu}^0(p; a)). \quad (10)$$

We will hold fixed the lattice spacing obtained in the $n$-th step. We thus set $2^na_0 = 1$, or $a_0 = 2^{-n}$. The blocking kernel $Q^{(j)}$ is defined as follows. We label the sites of the $(j-1)$-th lattice by four integers $l = (l_1, l_2, l_3, l_4)$, $l_\mu \in Z$. A site $l' = (l'_1, l'_2, l'_3, l'_4)$
of the $j$-th lattice is identified with the site $2l' = (2l_1, 2l_2, 2l_3, 2l_4)$ on the $(j - 1)$-th lattice. The blocking transformation assigns to a field variable on the $j$-th lattice its arithmetic mean over a $2^4$ hypercube on the $(j - 1)$-th lattice. Explicitly, \( (Qf)(l') = \sum_{l} Q(l', l) f(l) = 2^{-4} \sum_{r=0,1} f(2l' + r) \). This definition implies $c = 2^{-4}$ in Eq. (11). Using eq. (7c) we obtain

$$D_n^{-1}(p) = \alpha_n^{-1} - \sum_{\mu} \left( i(\gamma_{\mu} \otimes I)A_n^{\mu}(p) + (\gamma_5 \otimes \tau_3 \tau_\mu)B_n^{\mu}(p) \right),$$

(11a)

$$A_n^{\mu}(p) = \sum_{k^{(n)}_{\mu}} A_0^{\mu}(p + 2\pi k^{(n)}; 2^{-n}) |Q_n(p, k^{(n)})|^2,$$

(11b)

$$B_n^{\mu}(p) = \sum_{k^{(n)}_{\mu}} B_0^{\mu}(p + 2\pi k^{(n)}; 2^{-n}) |Q_n(p, k^{(n)})|^2,$$

(11c)

$$|Q_n(p, k^{(n)})|^2 = \prod_{\nu} \left( \frac{\sin(p_{\nu}/2)}{2^n \sin((p_{\nu} + 2\pi k^{(n)}_{\nu})/2^{n+1}))} \right)^2,$$

(11d)

where $-\pi \leq p_{\mu} \leq \pi$ and for each $\mu$, $k^{(n)}_{\mu} = -2^{-n-1}, -2^{-n-1} + 1, \ldots, 2^{-n-1} - 1$. To arrive at Eq. (11d), observe that we have set $a = 1$ for the coarse-lattice spacing, so that a single-step blocking kernel is a mapping into this lattice from a lattice with spacing equal to $1/2$. For each $\mu$, this kernel has the momentum representation $(1/2)(\exp(i\mu/2) + 1) = \exp(i\mu/4) \cos(p_{\mu}/4) = \exp(i\mu/4) \sin(p_{\mu}/2)/(2 \sin(p_{\mu}/4))$. Eq. (11d) follows by going “backwards” down to the fine lattice with spacing $a_0 = 2^{-n}$.

The massless staggered operator satisfies $\{D_0, (\gamma_5 \otimes \tau_5)\} = 0$, and therefore the staggered action is invariant under a $U(1)$ chiral symmetry.\footnote{This symmetry correspond to the $U(1)_c$ symmetry of the one-component formalism [13]. Note that the interpretation of this symmetry in the continuum limit = axial, vector or some combination of them = depends in general on the choice of the staggered mass term [13]. For the simple mass term of Eq. it is a chiral symmetry.} For the RG-blocked operator we have Eq. (11a), which implies that $D_n^{-1}(x, y)$ anti-commutes with $(\gamma_5 \otimes \tau_5)$ except for $x = y$. In fact,

$$\{D_n, (\gamma_5 \otimes \tau_5)\} = 2\alpha_n^{-1} D_n (\gamma_5 \otimes \tau_5) D_n, \quad n \geq 1.$$

(12)

This is recognized as a generalization of the GW relation [10]. Thus, after the first blocking transformation the $U(1)$ chiral symmetry gets modified to a Ginsparg-Wilson-Lüscher (GWL) chiral symmetry [11].

For Wilson fermions, it was proved in ref. [9] that the RG-blocked operator

$$D_n(x - y) = \int_{BZ} dp \ e^{ip(x - y)} D_n(p),$$

(13)

is local. Here $\int_{BZ} dp \equiv \int_0^{2\pi} \frac{d^3p}{(2\pi)^3}$ denotes the integration over the Brillouin zone of the coarse lattice. As explained earlier, this means that $D_n(x - y)$ decays exponentially with $|x - y|$, and the decay rate is $O(1)$ in units of the coarse-lattice spacing. The bounds established in the course of the proof are uniform in $n$, and hold for $\alpha \leq \hat{\alpha}$.
where \( \hat{\alpha} > 0 \) is an \( O(1) \) constant whose actual value can be worked out by keeping track of the details of the proof.

We now argue that the proof of locality continues to hold if, at the starting point, we replace the Wilson operator by the staggered Dirac operator \([\mathbf{8}]\). This amounts to replacing the Wilson term \( W \) by a “skewed” Wilson term \( W_{st} \), where

\[
W = \sum_{\mu} (1 - \cos(p_\mu a)),
\]
\[
W_{st} = \sum_{\mu} (\gamma_5 \otimes \tau_5 \tau_\mu) (1 - \cos(p_\mu a)).
\]

The proof requires lower and upper bounds on \( W \) as a function of \( p_\mu a \). Introducing the vector-space norms \( |x_\mu|_\gamma = (\sum_\mu |x_\mu|_\gamma)^{1/\gamma} \) we observe that \( W = |W| = |1 - \cos(p_\mu a)|_1 \).

In the staggered case, we have the operator norm \( \| W_{st} \|_2 = W_{st} W_{st}^\dagger = \sum_\mu (1 - \cos(p_\mu a))^2 \), or equivalently \( \| W_{st} \| = |1 - \cos(p_\mu a)|_2 \). Since the following equivalence-of-norms inequalities hold in \( d \) dimensions

\[
d^{1/2} |x_\mu|_2 \leq |x_\mu|_1 \leq d^{1/2} |x_\mu|_2,
\]

it follows that every lower or upper bound on \( |W| \) entails a corresponding bound on \( \| W_{st} \| \), and \textit{vice versa}. This simple argument shows that, indeed, the proof given in ref. \([\mathbf{11}]\) generalizes to the RG-blocked staggered Dirac operator in its flavor representation.

The key physical input that goes into the proof \([\mathbf{11}]\) is that \( D_0^{-1}(p) \) and \( D_n^{-1}(p) \) share the same singularity as \( p \to 0 \), namely, \( -i \sum_\mu \gamma_\mu p_\mu / p^2 \). Indeed, the singularity of \( D_n^{-1}(p) \) arises only from the \( k_\mu^{(n)} = 0 \) term on the right-hand side of Eq. \([11b]\).

Factoring out this singularity by writing \( D_0(p) = R_n(p)D_n(p) \), one can prove that the operator \( R_n(p) \) is analytic in \( p_\mu \) and that both \( R_n(p) \) and \( R_n^{-1}(p) \) are bounded. This is then used to prove that \( D_n(p) \) and \( G_n(p) \) are analytic, and that \( D_n(p), G_n(p) \) and \( G_n^{-1}(p) \) are bounded. Exponential localization of the corresponding coordinate-space kernels follows from general theorems.

We will next show that, in the limit of infinitely many RG steps, \( D_n \) becomes diagonal in flavor space. The flavor-mixing part of \( D_n^{-1}(p) \) is given by Eq. \([11c]\), where explicitly

\[
\mathcal{B}_\mu^0(p + 2\pi k^{(n)}; 2^{-n}) = \frac{2 \sin^2((p_\mu + 2\pi k_\mu^{(n)})/2^{n+1})}{2^n \sum_\nu \left( \sin^2((p_\nu + 2\pi k_\nu^{(n)})/2^n) + 4 \sin^4((p_\nu + 2\pi k_\nu^{(n)})/2^{n+1}) \right)}.
\]

It is easy to see that \( |\mathcal{B}_\mu^0(p + 2\pi k^{(n)}; 2^{-n})| \leq c_1 2^{-n} \) where \( c_1 = O(1) \). Away from the singularity at \( p + 2\pi k^{(n)} = 0 \) this is evident. For \( |p + 2\pi k^{(n)}| \ll 2^n \), the same result follows using \( \sin(x) \sim x \) for \( x \ll 1 \). Thus, \( |\mathcal{B}_\mu^0(p + 2\pi k^{(n)}; 2^{-n})| = O(2^{-n}) \), uniformly in \( p \) and \( k^{(n)} \). In addition, \( |Q_n(p, k^{(n)})|^2 \leq c_2 \prod_\nu |p_\nu|^2 / |(p_\nu + 2\pi k_\nu^{(n)})|^2 \) where again \( c_2 = O(1) \). Hence the \( k_\mu^{(n)} \)-summation converges, and \( |\mathcal{B}_\mu^0(p)| = O(2^{-n}) \) for all \( p \). Taking the limit \( n \to \infty \) we conclude that \( \mathcal{B}_\mu^0(p) \to 0 \), uniformly in \( p \).
The inverse fixed-point operator obtained in the limit \( n \to \infty \) can be expressed as
\[
D_\infty^{-1}(p) = D_{rg}^{-1}(p) \otimes I = \left( \alpha_\infty^{-1} - i \sum_\mu \gamma_\mu A_\infty^\mu(p) \right) \otimes I ,
\]
where \( \alpha_\infty = (15/16)\alpha \), and \( A_\infty^\mu(p) \sim p_\mu/p^2 \) for \( p_\mu \ll 1 \). This shows that \( D_\infty \) is diagonal in flavor space and satisfies the (standard) GW relation.\(^6\) Letting \( G_\infty = \lim_{n \to \infty} G_n \), Eqs. (3) and (17) imply
\[
\lim_{n \to \infty} \det 1/4(D_0(a_0 = 2^{-n})) = \det(D_{rg}) \det 1/4(G_\infty^{-1}) .
\]
In view of the locality and boundedness properties established above, the desired decomposition (2) is achieved in this limit!

The essence of the RG blocking is that it distills the long-distance dynamics, extracting it out of the underlying short-distance theory. The long-distance dynamics is contained in \( D_{rg} \) which is manifestly diagonal in flavor space, while all the flavor-mixing effects are contained in \( G^{-1}_\infty \). Since \( G^{-1} \) is analytic in momentum space and has an \( O(1) \) gap, its fourth root shares similar properties. Hence \( G^{-1/4} \) is local, and has only cutoff-mass excitations, uniformly in \( n \).

Let us elaborate on this last statement. Observing that \( \alpha \) has mass dimension equal to one, and focusing e.g. on the first blocking step, an RG transformation works by first replacing \( D_0 \) with \( G_{\text{rg}}^{-1} = D_0 + \alpha Q^\dagger Q \) (cf. Eq. (3b)). Now, the massless operator \( D_0 \) has vanishingly small eigenvalues near \( p_\mu = 0 \). The contribution from the blocking kernel, \( \alpha Q^\dagger Q \), lifts these small eigenvalues and generates an \( O(\alpha) = O(1) \) gap in the spectrum of \( G^{-1} \). We may define the fourth-root operator for any finite \( n \) via
\[
M_n(\tilde{x}, \tilde{y}) = \int_{BZ}^{(n)} d\tilde{p} e^{i\tilde{p}(\tilde{x} - \tilde{y})} M_n(\tilde{p}) , \quad M_n(\tilde{p}) = (G_n(\tilde{p})G_n^\dagger(\tilde{p}))^{-1/8} .
\]
In this equation, \( \tilde{x}, \tilde{y} \) take values on the fine lattice, and \( \int_{BZ}^{(n)} d\tilde{p} \) denotes the integration over the fine-lattice Brillouin zone. The argument why \( M_n(\tilde{x}, \tilde{y}) \) is local is standard \[6\]. If we let one of the momentum components become complex, the singularity closest to the real axis will be at a distance which is \( O(\alpha_n) \). Deforming the contour of integration, this implies that \( M_n(\tilde{x}, \tilde{y}) \) decays exponentially with \( |\tilde{x} - \tilde{y}| \), with a decay rate which is \( O(\alpha_n) \), namely, \( O(1) \) in units of the coarse-lattice spacing.

It is interesting to compare this result to ref. \[4\] which attempts to find an operator \( \mathcal{N} \) such that \( \det(\mathcal{N}) = \det 1/4(D_0) \), without the help of RG transformations (see also ref. \[16\]). In this case, the gap is provided by the \textit{physical} mass, and the decay rate of the square-root kernel needed in the construction is found to be \( O(\sqrt{m/a}) \). Hence, the limit \( m \to 0 \) is problematic.

With the help of RG blocking, the small-distance scale relevant for the fourth root is \( a_0 \), the spacing of the original fine lattice, and the relevant large-distance scale

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\(^6\)The relevance of the GW relation for establishing consistency of the fourth-root trick was recently pointed out in ref. \[15\].
is $a$, the spacing of the coarse lattice. In comparison, in refs. [16, 4] the relevant short- and long-distances scales are $a$ and $1/m$ respectively. The use of RG blocking effectively achieves the replacements $a \rightarrow a_0$, $1/m \rightarrow a$. While the mathematics is similar, the physical conclusion is different. The “fourth-root” kernel $\mathcal{M}_n(\tilde{x}, \tilde{y})$ is long-ranged with respect to the fine-lattice scale $a_0$ (in analogy with refs. [16, 4]); but the same kernel is short-ranged with respect to the coarse-lattice scale. This is true uniformly in $n$, hence also in the limit $n \rightarrow \infty$.

The RG analysis generalizes to the case that the mass term is switched back on in Eq. (8).

We find that Eq. (11a) still holds, except that the explicit expressions for $A^n_\mu(p)$, $B^n_\mu(p)$ and $\alpha^{-1}_n$ get modified (now $\alpha^{-1}_n = \alpha^{-1}_n(p)$ becomes a non-trivial function of $p$). The proof of locality generalizes to $m \neq 0$ for Wilson fermions [6], and the same should be true for the flavor-representation staggered fermions. Also, clearly $B^0_\mu(p; m \neq 0) \leq B^0_\mu(p; m = 0)$. Therefore, for $n \rightarrow \infty$, $B^n_\mu(p)$ tends to zero as before.

The limiting operator $D_{\text{rg}}$, defined by the first equality in Eq. (17), is again diagonal in flavor space. Of course, for $m \neq 0$, $D_{\text{rg}}$ will not satisfy the GW relation any more.

4. Interacting staggered fermions

It is unlikely that rigorous theorems such as those of ref. [6] will ever be generalized to an interacting lattice theory. Still, physical intuition suggests that similar statements on locality and boundedness may hold true in an interacting theory as well. In this section, I address some of the issues that arise when dealing with an interacting theory.

The question is what are the properties of an RG-blocked lattice theory, when the initial interacting theory involves one-component staggered fermions, and the fourth-root trick is applied. In an interacting theory RG transformations may be realized in numerous ways. A common feature is that RG transformations naturally give rise to a “two-cutoff” theory: the RG-blocked theory living on a coarse lattice with spacing $a$ is obtained after applying $n$ RG transformations to an initial theory defined on a fine lattice with spacing $a_0 \ll a$. For example, in the context of “perfect action” one applies RG transformations to fermion and gauge fields alike, and the limit $n \rightarrow \infty$ (and $a_0 \rightarrow 0$) is taken while keeping $a$ fixed [9].

Here I will limit the discussion to a simpler framework, where RG transformations are applied only to the fermion variables. Among other things, this has the advantage that some simple tests can be carried out on existing dynamical configurations.

The first problem that must be tackled is that, as discussed below, the interacting theory is defined using one-component staggered fermions for a reason [14, 17]. In the free-field case, there is a unitary operator $Q_0$ that maps the one-component staggered operator, defined on a lattice with spacing $a_0$, to the flavor-representation staggered operator on a lattice with spacing $a_1 = 2a_0$ [5]. In the interacting case, the mapping must preserve gauge invariance but there is no unique, obvious way to define it.

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7It is desirable (though not a necessary condition for $ma \ll 1$) to choose the same relative sign for $m$ and $\alpha$, such that $m + \alpha Q^1 Q$ is a strictly positive operator.
We propose to deal with this problem by a single RG blocking step which keeps the number of fermionic degrees of freedom unchanged. Specifically, in Eq. (3) we take $\psi, \bar{\psi}$ to be single-component fields, while $\chi, \bar{\chi}$ are four flavors (or “tastes”) of Dirac fields. $D_0$ is now a covariant, interacting one-component staggered operator. For the blocking kernel we choose some covariant version of $Q_0$, denoted $Q^{(0)}$ below, defined in terms of the link variables on the fine lattice. Eqs. (3) and (5) provide explicit expressions for $D_1$, which is the resulting staggered operator in the flavor representation, as well as for $G_1$, whose (inverse) determinant is picked up when performing this non-trivial change of variables. (We will shortly return to the role of $\det(G_1^{-1})$.) Using the natural embedding $x \to \tilde{x}$ from the coarse lattice to the fine lattice (as described above Eq. (11)), both $D_1$ and $G_1$ are gauge-covariant functions of the link variables on the fine lattice.

Having thus constructed a “flavor representation” in the interacting theory, we may apply $n$ additional, ordinary RG blocking transformations to the fermions. These block transformation dilute the number of fermionic degrees of freedom while maintaining gauge invariance, provided that we keep choosing blocking kernels which are covariant with respect to gauge transformations on the original fine lattice. This is naturally realized if, for any point $x'$ and any point $y$ such that $Q^{(j)}(x', y) \neq 0$ in the free theory, we construct the covariant kernel $Q^{(j)}(\tilde{x}', \tilde{y})$ by summing over Wilson lines that go from $\tilde{x}'$ to $\tilde{y}$ on the original fine lattice.

An important question is what is the fate of the global symmetries of the original, interacting one-component staggered-fermion theory. As for the chiral $U(1)$ symmetry, we have seen that it becomes a GWL symmetry (Eq. (12)). Now, it may be impossible to preserve manifest hypercubic invariance in the construction of $Q^{(j)}$, and the same is true for the staggered shift symmetry [14]. Since the original one-component theory has exact hypercubic and shift symmetries, by Eq. (3d), any breaking of these symmetries induced in the above-defined flavor representation by $D_n$ must be exactly compensated by $\det(G_n^{-1})$. In other words, the effective action $S_n^{\text{eff}} = \log(\det(G_n))$ should automatically contain the local “counter-terms” needed to restore exact invariance.

It is known that the attempt to construct an interacting theory directly in the flavor representation gives rise to induced cutoff-scale masses which, on top of that, violate Lorentz invariance [17]. The interacting staggered theory is defined in the one-component formalism because its global symmetries, including in particular shift symmetry, forbid these disastrous mass terms [14]. It is therefore crucial that the cancellation mechanism proposed above will indeed be operative when the flavor representation is constructed though “RG blocking” from the one-component theory. Interestingly, there exists a general RG-blocking result which provides direct evidence that the above cancellation mechanism indeed works as it should. Introduce the
“telescopic sum” \[6\]

\[
D_0^{-1} = D_0^{-1} Q_0 D_n Q_n D_0^{-1} + \sum_{j=0}^{n-1} D_0^{-1} (Q_j^\dagger D_j Q_j - Q_{j+1}^\dagger D_{j+1} Q_{j+1}) D_0^{-1}
\]

\[
= S_n D_n^{-1} \mathfrak{S}_n + \sum_{j=0}^{n-1} S_j \bar{\Gamma}_{j+1} S_j,
\]

(20)

where we have set \(Q_0 = I\) and where

\[
S_j \equiv D_0^{-1} Q_j^\dagger D_j, \quad \mathfrak{S}_j \equiv D_j Q_j D_0^{-1},
\]

(21)

and

\[
\bar{\Gamma}_{j+1} \equiv D_j^{-1} - D_j^{-1} Q^{(j+1)\dagger} D_{j+1} Q^{(j+1)} D_j^{-1} = \Gamma_{j+1}.
\]

(22)

The equality \(\bar{\Gamma}_{j+1} = \Gamma_{j+1}\) follows by substituting Eq. (6c) for \(D_j^{-1}\) and re-expanding as a geometric series.

Consider now an interacting staggered-fermion theory with light or massless quarks. Let us examine the long-distance behavior of the various kernels in Eq. (20). In ref. \[6\] it is proved that, in the free theory, \(S_j, \mathfrak{S}_j\) and \(\Gamma_j\) all have ranges which are \(O(1)\) in units of the coarse-lattice spacing. We expect the same to hold in the interacting case. Because the original kernel \(D_0^{-1}(\bar{x}, \bar{y})\) is long-ranged in physical units, the only way for Eq. (20) to hold is if \(D_n^{-1}(\bar{x}, \bar{y})\) is long-ranged too.

The physical interpretation of Eq. (20) is not completely straightforward because \(D_0^{-1}\) and \(D_n^{-1}\) are not gauge invariant all by themselves. Let us examine a physical observable. We will consider the gauge-invariant two-point function of the true Goldstone-boson field \(\pi(x)\). Its existence is implied by the symmetries of the one-component formalism \[14\], and so \(\langle \pi(x) \pi(y) \rangle\) decays like a power of \(|x - y|\) in the massless limit. We may construct this two-point function either from \(D_n^{-1}\) or, using Eq. (20), from \(D_n^{-1}\). Had any cutoff-scale masses been induced as in the case of ref. \[17\], \(D_n^{-1}\) would decay exponentially with a cutoff-scale rate, making it impossible to reproduce the correct long-distance behavior. We conclude that, whether or not the RG-blocked theory enjoys manifest shift and (full) hypercubic symmetries, the physical consequences of all the original symmetries remain intact!

Suppose that, instead, one had started with an interacting theory obtained by simply gauging the flavor representation using link variables that reside on the coarse lattice, as in ref. \[17\]. Assume also that the bare quark mass is set to zero. What would go wrong? Denoting the new interacting Dirac operator by \(D_0\) and the lattice spacing by \(a_0\). The two-point function of the would-be Goldstone boson is given explicitly by

\[
\langle \pi(x) \pi(y) \rangle = \langle \mathcal{G}(x, y) \rangle,
\]

(23a)

\[
\mathcal{G}(x, y) = \text{tr} \left( (\gamma_5 \otimes \tau_5) D_0^{-1}(x,y) (\gamma_5 \otimes \tau_5) D_0^{-1}(y,x) \right).
\]

(23b)

Since \((\gamma_5 \otimes \tau_5) D_0(\gamma_5 \otimes \tau_5) = D_0^\dagger\) one has \(\mathcal{G}(x, y) > 0\), i.e. this correlator is strictly positive. The decay rate of \(\langle \pi(x) \pi(y) \rangle\) is now \(O(1/a_0)\), because the explicit one-loop
calculation of ref. 17 shows that the quarks acquire (Lorentz-breaking) cutoff masses in this theory. Since \( G(x, y) > 0 \), it is ruled out that the short range could be generated by destructive interference between different gauge-field configurations; rather, given an ensemble of configurations, we must have \( \| D_0^{-1}(x, y) \|^2 \sim \exp(|x - y|/a_0) \) on each and every one of them.

One could have applied RG transformations, and Eq. 20 would again say that any long-distance physics of \( D_0 \) must be reproduced by the resulting \( D_n \). However, we have just seen that \( D_0 \) has no long-distance physics whatsoever. By Eq. 20, any number of RG transformations would not have fixed it.

Coming back to the interacting one-component staggered theory, the algebraic structure of \( G(x, y) \) is similar to Eq. (23b), and again \( G(x, y) > 0 \). In order to obtain the correct power-law behavior of the Goldstone-boson correlator \( \langle \pi(x) \pi(y) \rangle \), there must exist a finite probability to encounter configurations where \( G(x, y) \) has the same long-distance behavior. Finally, by Eq. 20, the long-distance behavior must be sustained if we use \( D_n^{-1} \) instead of \( D_0^{-1} \) to construct \( G(x, y) \). The long-distance physics contained in \( D_0^{-1} \) is faithfully reproduced by the RG-blocked \( D_n^{-1} \).

It will be interesting to test the resulting operators numerically, and see whether the properties established in the free theory persist. For the RG-blocked operator \( D_n \), the crucial properties are locality, suppression of the flavor-mixing (or taste-mixing) part with \( n \), and convergence to a GW operator in the massless case. As for \( G_n \) (Eq. (23)), the question is whether this operator indeed describes only excitations with masses of the order of the cutoff, so that the effective action \( S_{\text{eff}}^n = \log(\det(G_n)) \) is local. Last, it should be verified that the mechanism that protects the (physical consequences of the) symmetries of the original one-component formulation is indeed operative.

On the theoretical side, an interesting idea is to make use of the notion of admissibility condition. The original concept introduced in ref. 18 asserts that the lattice gauge field is constrained such that, for every plaquette, \( \text{Re} \, \text{tr}(1 - U_{\mu\nu}(x)) < \epsilon_0 \), where \( U_{\mu\nu}(x) \) is the ordered product of link variables around the given plaquette, and \( \epsilon_0 > 0 \) is a fixed (small) number. It is believed that an admissibility condition does not change the universality class. The utility of an admissibility condition is the following. Given a free lattice operator whose spectrum satisfies a certain bound, one expects that this bound will be modified only by \( O(\epsilon_0) \) if we promote the operator to a covariant one, while allowing only for gauge fields that satisfy the admissibility condition.

In the case at hand, we may envisage imposing an admissibility condition on the link variables of the original fine lattice. This should imply that, apart from \( O(\epsilon_0) \) modifications, the operator bounds established in ref. 6 will continue to hold during the first few blocking steps, and the same should follow for the locality properties.

Unlike in ref. 18, however, we now face the following problem. When the number of blocking transformations becomes of the order of \( n \sim 1/\epsilon_0 \), we no longer have any useful bound on the gauge field. We will thus propose a stronger notion of admissibility. Considering the product of (covariant) blocking kernels \( Q_n = Q^{(n)}(n-1) \ldots Q^{(1)}Q^{(0)} \), we will constrain the gauge field on the original fine lattice
by demanding that every Wilson loop $W$ occurring in the product $Q_nQ_n^\dagger$ will satisfy the constraint $|W| < \epsilon_0$. The reasoning behind this new admissibility condition is the following. After $n$ blocking steps, the length of the loops contained in $Q_nQ_n^\dagger$ will be $O(2^n)$ in units of the original lattice spacing $a_0$. However, as explained earlier, we are really interested in the “two cutoff” situation where $a_0 = 2^{-n}a$. When measured in units of the coarse-lattice spacing $a$, all these Wilson loops have length smaller than some $l_0 = O(1)$. Our new admissibility condition is therefore a natural generalization of the same concept to a “two cutoff” situation. With this new admissibility condition, it is plausible that the free-theory bounds will continue to hold for arbitrarily large $n$, up to $O(\epsilon_0)$ modifications. In other words, the deviations from the free-theory bounds will be $O(\epsilon_0)$ independently of $n$. In addition, for very large $n$, the flavor-breaking part of the RG-blocked operator will become very small. It should be possible to find large enough, but finite, $n$, such that full flavor (“taste”) symmetry is recovered to any desired accuracy.\(^8\)

In conclusion, using the machinery of RG block transformations I have shown that the fourth-root trick is consistent for free staggered fermions. In the limit of infinitely many RG steps, the free staggered-fermion determinant is equal to the fourth power of the determinant of a one-flavor local operator which, in the massless case, satisfies the GW relation, times the determinant of a local operator whose excitations have cutoff masses. The fourth root of the latter operator is local too. While a similar result for interacting staggered fermions is unlikely to be established anytime soon, I have discussed how to construct the flavor representation in the interacting theory, while in effect maintaining all the symmetries of the one-component formalism. I have also suggested avenues for numerical tests, as well as a theoretical framework which appears to be best suited for generalizing some of the rigorous free-theory results to the interacting case.

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\(^8\)Since in this framework the independent link variables always reside on the original fine lattice, the limit $n \to \infty$ cannot be taken independently of the continuum limit.
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