Whether composite fermion states with “wrong” chiralities dissolve into cuts

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Abstract

In the possible scaling region for lattice chiral fermions advocated in\textit{ Nucl. Phys. B}486 (1997) 282, no hard spontaneous symmetry breaking occurs and doublers are gauge-invariantly decoupled via mixing with composite three-fermion-states. However the strong coupling expansion breaks down due to no “static limit” for the low-energy limit ($p \sim 0$). We further analyze relevant Green functions of three-fermion-operators. It is shown that in the low-energy limit, the propagators of three-fermion-states with the “wrong” chiralities positively vanish due to the generalized form factors (the wave-function renormalizations) of these composite states vanishing as $O(p^4)$. This strongly implies that three-fermion-states with “wrong” chirality are “decoupled” in this limit.

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1 A possible scaling region of chiral fermions

The attempt to get around the “no-go” theorem [1] for the “vector-like” phenomenon of chiral fermions on a lattice is currently a very important issue of theoretical particle physics [2]–[13]. One attempt is the approach of multifermion couplings, which can be traced back from the recent lattice formulation of the Standard Model [14] to the pioneer model suggested in ref. [15] and successive work [16, 17, 18] in the past years. On the other hand, it was pointed out in a crucial paper [19] that the models proposed in ref. [15] fail to give chiral fermions in the continuum limit. The reasons are that an NJL [20] spontaneous symmetry breaking phase separates the strong-coupling symmetric phase from the weak-coupling symmetric phase, and the right-handed Weyl states do not completely disassociate from the left-handed chiral fermions. Further, another crucial paper [21] tried to extent the “no-go” theorem into interacting theories based on the plausible argument of such theories being local. The definite failure of the models so constructed has then been a general belief [2].

To more easily show the possibility and reason of dynamics for such constructed models to work, we further analyze the simple and anomaly-free model of multifermion couplings proposed in the ref. [17]. Note that \( \psi^i_L (i = 1, 2) \) is an \( SU_L(2) \) gauged doublet, \( \chi_R \) is an \( SU_L(2) \) singlet and both are two-component Weyl fermions. \( \chi_R \) is treated as a “spectator” fermion. \( \psi^i_L \) and \( \chi_R \) fields are dimensionful \( [a^{-2}] \). The following action for chiral fermions with the \( SU_L(2) \otimes U_R(1) \) chiral symmetries on the lattice is suggested:

\[
S = S_f + S_1 + S_2,
\]

\[
S_f = \frac{1}{2a} \sum_x \sum_\mu \left( \bar{\psi}^i_L(x) \gamma_\mu D_\mu^i \psi^i_L(x) + \bar{\chi}_R(x) \gamma_\mu \partial_\mu \chi_R(x) \right),
\]

\[
S_1 = g_1 \sum_x \bar{\psi}^i_L(x) \cdot \chi_R(x) \bar{\chi}_R(x) \cdot \psi^i_L(x),
\]

\[
S_2 = g_2 \sum_x \bar{\psi}^i_L(x) \cdot \left[ \Delta \chi_R(x) \right] \left[ \Delta \bar{\chi}_R(x) \right] \cdot \psi^i_L(x),
\]

where \( S_f \) is the naive lattice action for chiral fermions, \( a \) is the lattice spacing. \( S_1 \) and \( S_2 \) are two external multifermion couplings, where the \( g_1 \) and \( g_2 \) have dimension \( [a^{-2}] \), and the Wilson factor is given as,

\[
\Delta \chi_R(x) = \sum_\mu \left[ \chi_R(x + \mu) + \chi_R(x - \mu) - 2 \chi_R(x) \right],
\]

\[
2w(p) = \int_x e^{-ipx} \Delta(x) = \sum_\mu (1 - \cos(p_\mu)).
\]

Note that all momenta are scaled to be dimensionless, \( p = \tilde{p} + \pi_A \) where \( \pi_A \) runs over fifteen lattice momenta (\( \pi_A \neq 0 \)).

The action (1) has an exact local \( SU_L(2) \) chiral gauge symmetry,

\[
\sum_\mu \gamma_\mu D_\mu^i = \sum_\mu (U_\mu(x) \delta_{x,x+\mu} - U_\mu^\dagger(x) \delta_{x,x-\mu}), \quad U_\mu(x) \in SU_L(2),
\]
which is the gauge symmetry that the continuum theory (the target theory) possesses. The global flavour symmetry $U_L(1) \otimes U_R(1)$ is not explicitly broken in eq. (1), we will not discuss the property of violating fermion number in such a model.

It has been advocated\cite{17} there exists a plausible scaling region, which is a peculiar segment in the phase space of the multifermion couplings $g_1, g_2$,

$$\mathcal{A} = \left[ g_1 \to 0, g_2^{c,a} < g_2 < g_2^{c,\infty} \right], \quad a^2 g_2^{c,a} = 0.124, \quad 1 \ll g_2^{c,\infty} < \infty,$$

for chiral fermions in the low-energy limit. $g_2^{c,\infty}$ is a finite number and $g_2^{c,a}$ indicates the critical value above which the effective multifermion couplings associating to all doublers are strong enough, so that all doublers are gauge-invariantly decoupled. We qualitatively determined $a^2 g_2^{c,a} = 0.124$ in ref.\cite{17}. The crucial points for this scaling region to exist are briefly described in the following.

In segment $\mathcal{A}$ (4), the action (1) possesses a $\chi_R$-shift-symmetry \cite{22}, i.e., the action is invariant under the transformation:

$$\bar{\chi}_R(x) \to \bar{\chi}_R(x) + \bar{\epsilon}, \quad \chi_R(x) \to \chi_R(x) + \epsilon,$$  

where $\epsilon$ is independent of space-time. The Ward identity corresponding to this $\chi_R$-shift-symmetry is given as\cite{17} ($g_1 \to 0$),

$$\frac{1}{2a} \gamma_\mu \partial_\mu \chi'_R(x) + g_2 \Delta \left( \bar{\psi}_L(x) \cdot \Delta \chi_R(x) \psi_L(x) \right) - \frac{\delta \Gamma}{\delta \chi'_R(x)} = 0,$$  

where the “primed” fields are defined through the generating functional approach, and “$\Gamma$” is the effective potential. The important consequences of this Ward identity in segment $\mathcal{A}$ are:

1. the low-energy mode ($p \sim 0$) of $\chi_R$ is a free mode and decoupled:

$$\int_x e^{-ipx} \frac{\delta^{(2)} \Gamma}{\delta \chi'_R(x) \delta \chi'_R(0)} = i \frac{\gamma_\mu}{a \sin(p_\mu)};$$  

2. no hard spontaneous chiral symmetry breaking ($O(\frac{1}{a})$) occurs (see eqs. (30) and (31) in ref.\cite{17}),

$$\int_x e^{-ipx} \frac{\delta^{(2)} \Gamma}{\delta \psi_L^n(x) \delta \chi'_R(0)} = \frac{1}{2} \Sigma^i(p) = 0 \quad p = 0,$$  

in addition, we have:

$$\Sigma(p) = 0 \quad p \neq 0,$$

which is shown by the strong coupling expansion (see eq.(104) in \cite{17}).

\footnote{The soft symmetry breaking for the low-energy modes ($p \sim 0$) is allowed.}
For the strong coupling \( g_2 \gg 1 \) in the segment \( \mathcal{A} \), the following three-fermion-states comprising of the elementary fields \( \psi^i_L \) and \( \chi_R \) are bound:

\[
\Psi^i_R = \frac{1}{2a}(\bar{\chi}_R \cdot \psi^i_L)\chi_R; \quad \Psi^n_L = \frac{1}{2a}(\bar{\psi}^n_L \cdot \chi_R)\psi^i_L. \tag{10}
\]

These fermion bound states possess the "wrong" chiralities in contrast with the "right" chiralities possessed by the elementary fields \( \psi^i_L \) and \( \chi_R \). The two-point Green functions of these charged \( \Psi^i_R \) and neutral \( \Psi^n_L \) have poles at the total momentum \( p = \pi_A \) \({}^7\),

\[
S^{ij}_{MM}(p) = \int_x e^{-ipx} \langle \Psi^i_R(0)\bar{\Psi}^j_R(x) \rangle \simeq \delta_{ij} \frac{i}{a^2} \frac{\sum \mu \sin p^\mu \gamma_\mu}{\sum \sin^2 p_\mu + M^2(p)} P_L; \tag{11}
\]

\[
S^n_{MM}(p) = \int_x e^{-ipx} \langle \Psi^n_L(0)\bar{\Psi}^n_L(x) \rangle \simeq \frac{i}{a^2} \frac{\sum \mu \sin p^\mu \gamma_\mu}{\sum \sin^2 p_\mu + M^2(p)} P_R; \tag{12}
\]

\[
M(p) = 8ag_2w^2(p), \tag{13}
\]

where \( p \sim \pi_A \) and \( w^2(p) \neq 0 \). And the two-point Green functions for doublers \((p \sim \pi_A)\) of the elementary fields \( \chi_R \) and \( \psi^i_L \) are given by,

\[
S^{ij}_{LL}(p) = \int_x e^{-ipx} \langle \psi^i_L(0)\bar{\psi}^j_L(x) \rangle \simeq \delta_{ij} \frac{i}{a^2} \frac{\gamma_\mu \sin(p)^\mu}{\sin^2 p + M^2(p)} P_R; \tag{14}
\]

\[
S_{RR}(p) = \int_x e^{-ipx} \langle \chi_R(0)\bar{\chi}_R(x) \rangle \simeq \frac{i}{a^2} \frac{\gamma_\mu \sin(p)^\mu}{\sin^2 p + M^2(p)} P_L. \tag{15}
\]

The three-fermion-states \((\mathcal{I}, \mathcal{II}, \mathcal{III})\) are Weyl fermions and respectively mix with the doublers of the elementary Weyl fields \( \bar{\chi}_R \) and \( \bar{\psi}^i_L \) \((\mathcal{I}, \mathcal{II}, \mathcal{III})\),

\[
S^{ij}_{ML}(p) = \int_x e^{-ipx} \langle \Psi^i_R(0)\bar{\psi}^j_L(x) \rangle \simeq \delta_{ij} \frac{M(p)}{\sin^2 p + M^2(p)} P_R; \tag{16}
\]

\[
S^n_{MR}(p) = \int_x e^{-ipx} \langle \Psi^n_L(0)\bar{\chi}_R(x) \rangle \simeq \frac{M(p)}{\sin^2 p + M^2(p)} P_L. \tag{17}
\]

As a result, the neutral \( \Psi^n \) and charged \( \Psi^i \) Dirac fermions are formed\(^2\),

\[
\Psi^i = (\psi^i_L, \Psi^i_R); \quad \Psi^n = (\psi^n_L, \chi_R), \tag{18}
\]

and the spectrum is vector-like and massive. These show that all doublers are decoupled as very massive Dirac fermions consistently with the \( SU_L(2) \otimes SU_R(1) \) chiral symmetries, since the three-fermion-states \((\mathcal{I})\) carry the appropriate quantum numbers of the chiral groups that accommodate \( \psi^i_L \) and \( \chi_R \).

\(^2\)The propagators of these Dirac fermions can be obtained by summing all relevant Green functions shown in above.

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Eqs. (11, 12) for doublers is obtained by the strong-coupling expansion in powers of $\frac{1}{g_2}$. For the strong coupling $g_2 \gg 1$, the kinetic terms can be dropped and the strong-coupling limit is given as,

$$Z = \Pi_{x\alpha} \int [d\bar{\chi}^\alpha_R(x) d\chi^\alpha_R(x)] [d\bar{\psi}^i_{L\alpha}(x) d\psi^i_{L\alpha}(x)] \exp (-S_2(x))$$

$$= (2g_2)^{4N} \left( \det \Delta^2(x) \right)^4, \quad g_2 \gg 1,$$

where the determinant is taken only over the lattice space-time and $N$ is the number of lattice sites. For the non-zero eigenvalues of the operator $\Delta^2(x)$ (2), which are associating to the doublers ($p \simeq \pi_A$) of $\psi^i_L(x)$ and $\chi_R(x)$, eq. (19) shows the existence of a sensible strong-coupling limit. Note that the operator $\Delta(x) \sim 2w(p)$ (2,19) has different eigenvalues 4 $\sim$ 6 with respect to different doublers ($p \sim \pi_A$), and the strong coupling expansion is actually in terms of powers of $\frac{1}{4g_2w^2(p)}$. As the consequence, the two-point Green functions (11-17) computed by the strong coupling expansion with respect to doublers should be a good approximation even for the intermediate coupling $a^2g_2 \sim O(1)$.

This discussion agrees with the qualitatively determined critical value $g_2^{c,a} = 0.124$ in (1), above which all doublers are decoupled via eqs. (18).

However, as for the zero eigenvalues of the operator $\Delta^2(x)$, which precisely correspond to the low-energy modes ($p \sim 0$) of $\psi^i_L(x)$ and $\chi_R(x)$, this strong-coupling limit is trivial and the strong-coupling expansion is actually in terms of powers of $\frac{1}{g_2}$. This physically means the weakness of the effective multifermion coupling for such low-energy modes of $\psi^i_L(x)$ and $\chi_R(x)$. Nevertheless, we cannot exclude the possibility of low-energy modes of three-fermion-states (10), which are represented by the poles $p \sim 0$ of the propagators (11,12).

On the other hand, as a consequence of the multifermion interacting action (1) being local, in the strong coupling limit the effective action (inverse propagator), which is bilinear in terms of interpolating fields, should be local and analytical in the whole Brillouin zone. Thus, the “no-go” theorem of Nielsen and Ninomiya is still applicable to this case [21]. Based on such an observation, one might argue the existence of the massless spectrum of the charged and neutral three-fermion-states (10) by the analytic continuation of their propagators (11,12) from $p \sim \pi_A$ to $p \sim 0$. As a result, the low-energy spectrum (18) is also vector-like.

Indeed, due to the locality of action (1) presented in this paper, all Green functions must be analytically continuous functions in energy-momentum space, provided the dynamics is fixed by given $g_1$ and $g_2$. In the strong coupling symmetric phase (PSM) where $g_1 \gg 1$ and a sensible strong-coupling limit exists, two-point Green functions for three-fermion-states (11) are essentially indistinguishable from that of the elementary fermion fields $\psi^i_L(x)$ and $\chi_R(x)$ appearing in the Lagrangian. In such a phase, the analytical continuity of Green functions.
for both elementary and composite fields in the whole momentum space does result in the “vector-like” phenomenon, as asserted by the “no-go” theorem for a free-fermion theory. The only loophole would appear if the propagators of interpolating fields (three-fermion-states) properly vanished and no longer had poles at \( p \sim 0 \). This indicates that at \( p \sim 0 \), these three-fermion-states dissolve into three-fermion-cuts, where the “no-go” theorem is entirely inapplicable.

2 Three-fermion-cuts

We turn to discuss how these three-fermion-states (12,11) with the “wrong” chiralities dissolve into three-fermion-cuts in segment \( \mathcal{A} \) for the low-energy limit \( (p \to 0) \). These three-fermion-cuts\[18, 23\]:

\[
C[\Psi^L_n(x)], \quad C[\Psi^R_i(x)],
\]

are the virtual states of three individual chiral fermions with a continuous energy spectrum, provided the total momentum \( p \) is fixed. However, these virtual states carry exactly the same quantum numbers and total momentum \( p \) as that of three-fermion-states. Thus, gauge symmetries are preserved in such a phenomenon of dissolving. The dynamics of the three-fermion-states dissolving into their virtual state is that the negative binding energy of three-fermion-states goes to zero. In the energy plane, it was shown that due to the variety of effective interactions (potential), the poles for bound states can be analytically continued to the cuts for virtual states on the physical sheet\[24\]. Presumably, in this analytical continuation of effective interactions, no other dynamics, e.g. spontaneous symmetry breaking, takes place. In segment \( \mathcal{A} \), the weakness of effective multifermion couplings for the low-energy modes of \( \psi^L_n \) and \( \chi^R \) could lead to the vanishing of the binding energy of the three-fermion-states (10). We can conceive a “dissolving” scale (threshold) \( \epsilon \)

\[
\tilde{v} \ll \epsilon < \frac{1}{a},
\]

where \( \tilde{v} \) is the possible soft spontaneous symmetry breaking scale \( (a\tilde{v} \simeq 0) \). From inequality (22), we understand that no hard spontaneous symmetry breaking in segment \( \mathcal{A} \) is extremely crucial for the possibility of analytical continuation of propagators from poles for three-fermion-states to cuts for virtual states of three individual fermions. At the dissolving scale \( \epsilon \gg \tilde{v} \), we can approximately treat elementary massive fermions as massless.

In the relativistic Lagrangian approach, to discuss the property of three-fermion-states dissolving into three-fermion-cuts, we are bound to dynamically calculate two-point functions of three-fermion-states (Fig.1) to identify not only their poles, but also the corresponding residues. Using the strong coupling expansion, we approximately determined the simple poles for doublers \( (p \sim \pi A) \) in
The residues $Z_R(p)$ and $Z_L(p)$ of these simple poles are defined as,

$$ S_{MM}^{ij}(p) = \int_x e^{-ipx} \langle \Psi_R^i(0) \bar{\Psi}_R^j(x) \rangle \simeq \delta_{ij} \frac{Z_R(p)}{\alpha^2} \sum_{\mu} \sin^2 p_{\mu} Z_R(p) \frac{1}{\alpha^2} \sum_{\mu} \sin^2 p_{\mu} + M^2(p) P_L; \quad (23) $$

$$ S_{MM}^{n}(p) = \int_x e^{-ipx} \langle \Psi_L^n(0) \bar{\Psi}_L^n(x) \rangle \simeq \frac{Z_L(p)}{\alpha^2} \sum_{\mu} \sin^2 p_{\mu} Z_L(p) \frac{1}{\alpha^2} \sum_{\mu} \sin^2 p_{\mu} + M^2(p) P_R. \quad (24) $$

In fact, these residues represent the generalized form factors of three-fermion-states. The $Z_{L,R}(p)$ momentum dependence indicates that different doublers have different form factors, which implies the “size” of bound states is different from one doubler to another. This is clearly attributed to the momentum dependence of effective multifermion couplings in action (1). If these residues $Z_{L,R}(p = \pi A)$ are positive and finite constants with respect to each doubler, we can just make a wave-function renormalization of three-fermion-states with respect to each doubler,

$$ \Psi_R^i|_{ren} = Z_R^{-1}\Psi_R^i; \quad \Psi_L^n|_{ren} = Z_L^{-1}\Psi_L^n, \quad (25) $$

and the two-point Green functions (23,24) turn into eqs.(11,12) in terms of the renormalized fields (23).

The residues (generalized form factors) $Z_{RL}(p)$ (23,24) of the three-fermion-states (11) are given by one-particle irreducible (1PI) truncated Green functions (see Fig.2),

$$ Z_L(p) = \int_x e^{-ipx} \frac{\delta^{(2)}\Gamma}{\delta\bar{\Psi}_L^n(x)\delta\bar{\Psi}_R^j(0)}, \quad Z_R(p) = \int_x e^{-ipx} \frac{\delta^{(2)}\Gamma}{\delta\bar{\Psi}_R^n(x)\delta\bar{\Psi}_L^j(0)}. \quad (26) $$

The “primed fields” $\bar{\Psi}_R^n(x)$ and $\bar{\Psi}_L^n(x)$ of three-fermion-states are defined by eqs.(41) and (42) in ref.[17] through the generating functional approach. These generalized form factors $Z_L(p)$ and $Z_R(p)$ give the overlap between three-fermion-operators ($\bar{\Psi}_R^i(x)$, $\bar{\Psi}_L^n(x)$) and the interpolating three-fermion-states, that appear in the space of asymptotic states of the theory in the scaling region.

This description coincides with the renormalization of n-point 1PI functions with insertions of composite operators. In general, the renormalized n-point 1PI functions $\Gamma_{ren}^{(n)}$ with single and two insertions of composite operators are given by[25],

$$ \Gamma_{ren}^{(n)}(p_1, q_1, q_2, \cdots, q_n) = \frac{\Gamma_{reg}^{(n)}(p_1, q_1, q_2, \cdots, q_n)}{\alpha^2} P_L; \quad (27) $$

$$ \Gamma_{ren}^{(n)}(p_1, p_2, q_1, q_2, \cdots, q_n) = \frac{\Gamma_{reg}^{(n)}(p_1, p_2, q_1, q_2, \cdots, q_n)}{\alpha^2} P_R; \quad (27) $$

where $\Gamma_{reg}^{(n)}$ are the regularized n-point 1PI functions and $p_1$ and $p_2$ stand for the momenta entering the composite operators. Similarly given by eq.(26) (Fig.2)[25], the $Z$’s are the generalized “wave-function renormalizations” of composite operators. It is worthwhile to stress that for residues $Z_{L,R}$ being positively finite,

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3We do not want to have ghost states with negative norm.
the wave-function renormalization of composite fields is the exactly same as the wave-function renormalization of elementary fields. In fact, composite particles are indistinguishable from elementary particles in this case. However, the normal wave-function renormalizations of elementary fields appearing in the lagrangian is attributed to the fact that these elementary fields are defined at different scales rather than their "form factor". Note that the elementary fields $\psi^i_L(x)$ and $\chi^R(x)$ in the action (1) are bare fields and have not yet been renormalized.

However, in the analytical continuation of the propagators (23, 24) in momentum space, let us assume an interesting case that the residues $Z_{L,R}(p)$ positively vanish in the limit $p \to 0$ for the pole $p \sim 0$ in eqs.(23, 24),

$$Z_{L,R}(p) \to \mathcal{O}(p^n) \quad p \to 0,$$

(28)

with $n = 2, 4, 6, \cdots$. This property of $Z_{L,R}(p)$ could be realized by an appropriate momentum-dependence of the effective multifermion couplings $g_1, g_2$. Eq.(28) implies that $p \sim 0$ is no longer a pole for a relativistic particle in the propagators (23, 24). We are not allowed to make wave-function renormalization (25) with respect to $p \sim 0$. Eq.(28) indicates that the "size" of bound states (10) increases as $p \to 0$. These three-fermion-states may eventually dissolve into the virtual states of three individual fermions, whose possible configuration in momentum space is $(p, p, -p)$ \cite{18, 26}, where $p$ is the total momentum and the relative momentum ($q$) is zero. This dissolving phenomenon is entirely determined by both the dynamical and kinetic properties of the interacting theory.

This is reminiscent of the papers\cite{27} discussing whether helium is an elementary or composite particle based on the vanishing of wave-function renormalizations of composite states. It is normally referred to as the composite condition that the wave-function renormalizations of bound states go to zero ($Z \to 0$) \cite{28}. So far, we only give an intuitive and qualitative discussion of the dissolving phenomenon on the basis of the relations between the residues (generalized form factors) $Z_{L,R}(p)$, renormalized three-fermion-states and virtual states of three individual fermions (three-fermion-cut). Evidently, we are bound to do some dynamical calculations to show this phenomenon could happen.

### 3 The three-fermion-cut for neutral channel

The form factor $Z_L(p)$ in eq.(26) for the left-handed three-fermion-state $\Psi^L_n$ can be completely determined by the Ward identity (1). We appropriately take functional derivative of the Ward identity (1) with respect to the “primed” field $\Psi^p_L(x)$, and we obtain,

$$Z_L(p) = aM(p),$$

(29)

which are positive and finite constants for $p = \pi_A$ (see eq.(13)). Together with the propagators (14) obtained by the strong coupling expansion for $g_2 \gg 1$ and
\( p \sim \pi_A \), we conclude that the doublers of the neutral channel (24) are indeed relativistic massive particles, whose wave functions can be renormalized according to (25).

Given the strong coupling \( g_2 \gg 1 \), in spite of the propagator (24) for the neutral three-fermion-state \( \Psi^n_L \) resulting from the strong coupling expansion for \( p \sim \pi_A \), we can make an analytical continuation from \( p \sim \pi_A \) to \( p \sim 0 \). However, for the reason that \( Z^2_L(p) \to O(p^8) \) as \( p \to 0 \), we cannot conclude that \( p \sim 0 \) is a simple pole for a relativistic massless particle by an analytical continuation of the propagator (24) from \( p \sim \pi_A \) to \( p \sim 0 \), whereas we are not allowed to perform a wave-function renormalization (25) for \( p \sim 0 \). Actually, the propagator (24) is vanishing at \( p \sim 0 \). It is important to point out that at the limit of \( p \to 0 \), the vanishing of the propagator (24) is definitely positive, i.e., it is a double zero. This implies that ghost states with negative norm would not appear in low-energy spectrum.

However, on the other hand, we cannot conclude that the vanishing of the neutral propagator (24) for \( p \sim 0 \) indicates the virtual state of three fermions in the neutral channel. We must compute exactly the same two-point Green function for the neutral three-fermion-operator,

\[
\int_x e^{-ipx} \langle \Psi^n_L(0) \bar{\Psi}^n_L(x) \rangle, \tag{30}
\]

for \( g_2 \gg 1 \) and \( p \sim 0 \), as that (24) computed by the strong coupling expansion for \( p \sim \pi_A \) and \( g_2 \gg 1 \). Since the coupling \( g_2 \) in segment \( \mathcal{A} \) can be arranged such that

\[
a^2 g_2 w(p) < 1, \quad \infty > g_2 \gg 1, \quad p \sim 0, \tag{31}
\]

we can adopt the effective weak coupling expansion in powers of \( a^2 g_2 w(p) \) to calculate the Green function (30). This is due to the fact that \( \chi_R \) is always an external field in Feynman diagrams (see Fig.3) of computing this Green function (30) and each \( \chi_R \) is associated with \( g_2 w(p) \). This is a very reliable approximation as far as condition (31) holds. Here, we want to stress that the meaning of \( p \sim 0 \) is,

\[
\epsilon > p \gg a\tilde{v} \simeq 0. \tag{32}
\]

We can determine the value of \( g_2 \) in such a way that the threshold scale \( \epsilon \) is much larger than the soft symmetry breaking scale \( \tilde{v} \).

Clearly, no bound states in eq. (30) can be formed for such a weak effective coupling. Only the virtual state for the neutral three-fermion-cut (21) can be found in this Green function (30). We show this by calculating (30) in the effective weak coupling expansion as indicated in the Feynman diagrams (Fig.3). The leading order \( (O(1)) \) of this expansion is the first diagram in Fig.3, which corresponds to two parts,

\[
W_o(x) = -\left( \frac{1}{2a} \right)^2 \langle \psi^\gamma_L(0) \bar{\psi}^\delta_L(x) \rangle \langle \psi^\delta_R(0) \bar{\psi}^\gamma_R(x) \rangle \langle \psi^\alpha_L(0) \bar{\psi}^\beta_L(x) \rangle, \tag{33}
\]
\[ W'(x) = \left( \frac{1}{2a} \right)^2 \langle \bar{\psi}^\delta_L(0) \psi^\gamma_R(x) \rangle \langle \psi^\delta_R(0) \bar{\psi}^\gamma_L(x) \rangle \langle \bar{\psi}^\delta_L(0) \bar{\psi}^\gamma_R(x) \rangle, \]  

where \( \gamma, \delta, \beta, \alpha \) are spinner indices. In momentum space, eqs.(33) and (34) are given as,

\[ W_\circ(p) = -\int qS_{ji}LL(p + q) \text{tr} \left[ S_{RR}(k - \frac{q}{2})S_{ij}LL(k + \frac{q}{2}) \right] \left( \frac{1}{2a} \right)^2, \]  

\[ W'_\circ(p) = \int qS_{ji}LL(p + q) \text{tr} \left[ \Sigma^i(k - \frac{q}{2}) \right] \text{tr} \left[ \Sigma^j(k + \frac{q}{2}) \right] \left( \frac{1}{2a} \right)^2, \]  

where \( p \) is the total external momentum, \( k \) and \( q \) are the relative internal momenta. Since segment \( A \) is an entirely symmetric phase, i.e. \( \Sigma^i(k) = 0 \) (see eqs.(8,9)), eq.(36) identically vanishes. In general, we can write the internal propagators \( S_{ij}LL(k) \) and \( S_{RR}(k) \) in eq.(35) as follow,

\[ S_{ij}LL(k) = \delta_{ij}f_L(k^2)\gamma_\mu \sin k_\mu P_R, \]  

\[ S_{RR}(k) = f_R(k^2)\gamma_\mu \sin k_\mu P_L. \]  

These internal propagators \( S_{ij}LL(k) \) and \( S_{RR}(k) \) with internal momentum \( k \in (0, \pi_A) \) can be given by eqs.(14,15) calculated by the strong coupling expansion in segment \( A \). We define,

\[ R(k, q) = \text{tr} \left[ S_{RR}(k - \frac{q}{2})S_{ij}LL(k + \frac{q}{2}) \right], \]  

where we ignore indices \( i, j \) in the LHS. Since \( R(k, q) \) is an even function with respect to \( k \) and \( q \),

\[ R(k, q) = R(-k, q), \quad R(k, q) = R(k, -q), \]  

one can easily show for the external momentum \( p \sim 0 \),

\[ W_\circ(p) \simeq ac_\circ(i\gamma_\mu p^\mu) + O(p^2), \]  

where \( c_\circ \) is a constant.

The second diagram of Fig.3 denotes the contribution of order \( O[(g_2w(p))^2] \) to the Green function (30) given by \( \tilde{W}_1(p) \) given by,

\[ W_1(p) = [g_2w(p)]^2 V(p)\tilde{S}_{RR}(p)V(p), \]  

where \( \tilde{S}_{RR}(p) \) is the full propagator of \( \chi_R \), as indicated by a full circle in the middle of the Feynman diagram (Fig.3), which is summed over all contributions of this effective weak coupling expansion (Fig.4). We can adopt eq.(15) for \( \tilde{S}_{RR}(p) \)

\footnote{The contribution (34) is zero, and other possible contributions are identically zero, owing to \( \Sigma(k) = 0 \).}
by an analytical continuation from \( p \sim \pi_A \) to \( p \sim 0 \). On the other hand, as a consequence of the \( \chi_R \)-shift-symmetry (7), \( \tilde{S}_{RR}(p) \) for \( p \sim 0 \) is a free propagator,
\[
\tilde{S}_{RR}(p) \simeq \frac{a}{i\gamma_\mu p^\mu}.
\]
(43)

In eq.(42), \( V(p) \) is given by
\[
V(p) = \int_{qk} 4w(k - \frac{q}{2})S^{ji}_{LL}(k + q)tr \left[ S_{RR}(k - \frac{q}{2})S^{ij}_{LL}(k + \frac{q}{2}) \right],
\]
(44)
where the factor \( 4w(k - \frac{q}{2}) \) comes from the interaction vertex. Using a similar analysis to that giving eq.(41), we get for \( p \sim 0 \) (up to a finite constant),
\[
V(p) \sim -a(i\gamma_\mu p^\mu),
\]
(45)
which cancels the pole of eq.(43) at \( p \sim 0 \). As a result, we obtain \( w(p) \sim p^2 \), \( p \sim 0 \),
\[
W_1(p) \simeq (a^2 g_2 p^2)^2ac_1(i\gamma_\mu p^\mu),
\]
(46)
where \( c_1 \) is a finite number. Thus, for \( g_2 \gg 1 \) and \( p \sim 0 \), the Green function (34) is,
\[
\int_x e^{-ipx} \langle \psi^n_L(0) \bar{\psi}^n_R(x) \rangle \simeq W_0(p) + W_1(p),
\]
(47)
which is regular at \( p \sim 0 \). This shows that the neutral channel is a virtual state for the three-fermion-cut in region (31). Obviously, it is absolutely incorrect for doublers \( (p \sim \pi_A) \) since the effective coupling (31) can be extremely large, bound states \( \Psi^n_L(x) \) (10) must be formed.

Note that in this region (31), the mixing between the elementary field \( \chi_R \) and neutral three-fermion-state \( \Psi^n_L \) (10) calculated by the strong coupling expansion for \( p \sim \pi_A \) is given by (17). Analytical continuation of (17) from \( p \sim \pi_A \) to \( p \sim 0 \) shows the vanishing of the mixing at \( p \sim 0 \), while this mixing gives rise to the gauge-invariant mass terms for doublers \( p \sim \pi_A \). On the other hand, the mixing (17) can be calculated by the effective weak coupling expansion (31). The leading order is,
\[
\int_x e^{-ipx} \langle \psi^n_L(0) \bar{\psi}^n_R(x) \rangle = \int_x e^{-ipx} \left( \frac{1}{a} \right) \langle \psi^n_L(0) \bar{\psi}^n_R(x) \rangle \langle \psi_R(0) \bar{\psi}^n_L(0) \rangle + O(g_2 w(p)),
\]
(48)
which vanishes for no hard spontaneous symmetry breaking.

Based on the computations of the Green function for the neutral channel (30) at both \( p \sim \pi_A \) (24) and \( p \sim 0 \) (17) for \( g_2 \gg 1 \), we conclude that in the intermediate region (31) of the gauge-symmetric segment \( A \), the low-energy spectrum \( (p \sim 0) \) is

1. undoubled for all doublers \( (p \sim \pi_A) \) decoupled;
2. chiral because only exists a free right-handed mode \( \chi_R \), while \( \Psi^0_L \) is no longer a bound state, it is instead a virtual state \( \mathcal{C}[\Psi^0_L] \).

By the analytical continuity of the Green function (propagator) \((30)\) in terms of the total momentum \( p \), from eq.(24) for \( p \sim \pi_A \) to eq.(47), we must find the threshold scale \( \epsilon(22) \), where the neutral three-fermion-state \( \Psi^0_L(x) \) \((10)\) with the “wrong” chirality dissolves into its virtual states and only \( \chi_R \) remains as a relativistic particle at \( p \sim 0 \). We emphasize that this already violates the “no-go” theorem even though the chiral fermion is neutral. However, it is worthwhile to point out that such a mechanism violating the “no-go” theorem is only expected to work in the cases of neutral and anomaly-free theories. On the contrary, we recall that in the PMS phase of the Smit-Swift model \((29)\), the wave-function renormalization of the composite neutral field is given by the vacuum expectation value of scalar fields, \( z^2 = \langle V^\dagger(x)V(x+a) \rangle \) \((30)\), which is not momentum dependent.

4 The three-fermion-cut for the charged channel

The form factor \( Z_R(p) \) \((24)\) for the right-handed three-fermion-state \( \Psi^i_R(x) \) \((14)\) can not be determined by the Ward identity \((11)\). Instead, it can be calculated by using the results obtained in the strong coupling expansion for \( g_2 \gg 1 \) and \( p \sim \pi_A \).

The 1PI-vertex function associated to \( Z_R(p) \) is given by the truncated Green function \((26)\) that is defined as (as indicated in Fig.2),

\[
Z_R(p) = \int e^{-ipx} \frac{\delta(2)\Gamma}{\delta \Psi_R^0(0)\delta \bar{\Psi}_L^0(x)} = \int e^{-ipx} \int_{z_1,z_2} \left( G^{ij}_{MM}(x,z_2) \right)^{-1} G^{lk}_{ML}(z_2,z_1) \left( G^{ki}_{free}(z_1,0) \right)^{-1} \\
= \left( S^{ij}_{MM}(p) \right)^{-1} S^{lk}_{ML}(p) \left( S^{ki}_{free}(p) \right)^{-1},
\]

(49)

where

\[
G^{ij}_{MM}(x,z_2) = \langle \Psi^i_R(x)\bar{\Psi}^j_R(z_1) \rangle \rightarrow S^{ij}_{MM}(p) \\
G^{lk}_{ML}(z_2,z_1) = \langle \Psi^l_R(z_2)\bar{\Psi}^k_L(z_1) \rangle \rightarrow S^{lk}_{ML}(p') \\
G^{ki}_{free}(z_1,0) = \langle \bar{\Psi}^k_L(x)\bar{\Psi}^i_L(z_1) \rangle \rightarrow S^{ki}_{free}(p''), \quad g_1, g_2 = 0,
\]

(50)

and their transformations into the momentum space in the last line of eq.(49).

Adopting the results \( S^{ij}_{MM}(p) \) \((14)\), \( S^{ij}_{ML}(p) \) \((16)\) and the free propagator \( S^{ki}_{free}(p) \) of the \( \psi^i_L \), we get,

\[
Z_R(p) = aM(p),
\]

(51)
which is the same as \( Z_L(p) \) (29) directly derived from the Ward identity (8). Eq. (51) is a positive and finite constant for doublers \( p \simeq \pi_A \) (see eq. (13)). Together with the propagator (23) obtained by the strong coupling expansion for \( g_2 \gg 1 \) and \( p \sim \pi_A \), we conclude that the doublers of charged channel (23) are indeed relativistic massive particles, whose wave functions can be renormalized according to (25).

In spite of eqs. (23) and (51) obtained by the strong coupling \((g_2 \gg 1)\) expansion for \( p \sim \pi_A \), we can analytically continue the momentum “\( p \)” in these equations to the limit of \( p \to 0 \). When \( p \to 0 \), \( Z_L^2(p) \to O(p^8) \), the propagator (23) of charged three-fermion-state vanishes. This implies that the low-energy state \((p \sim 0)\) of \( \Psi^i_R(x) \) with the “wrong” chirality is no longer a simple pole as a relativistic particle, instead it is a virtual state. We stress again that at the limit of \( p \to 0 \), the vanishing of the propagator of the charged three-fermion-state is definitely positive, i.e., it is a double zero, which implies that ghost states with negative norm do not appear and couple to the gauge field in the low-energy limit. Otherwise, the theory would be inconsistent [31].

On the other hand, in order to directly show that the low-energy state \((p \sim 0)\) of \( \Psi^i_R(x) \) is the virtual state for a three-fermion-cut. We have to compute exactly the same two-point Green function,

\[
\int_x e^{-ipx} \langle \Psi^i_R(0) \bar{\Psi}^j_R(x) \rangle,
\]

goals to the charged channel for \( p \sim 0 \) and \( g_2 \gg 1 \) as that (23) computed by the strong coupling expansion for \( p \sim \pi_A \) and \( g_2 \gg 1 \). However, unlike the case of neutral channel, we do not have the reliable method of the effective weak coupling expansion in powers of \( a^2 g_2 w(p) \) (31), since \( \psi^i_L(x) \) is always an external field in the Feynman diagrams computing the Green function (52) and each \( \psi_L \) does not associate with \( g_2 w(p) \). In this case, the effective weak coupling expansion used for computing the neutral channel (Fig.3) must breakdown for \( g_2 \gg 1 \).

Nevertheless, we observe that the last binding threshold is located at \( a^2 g_2^c = 0.124 \) (4). Above this threshold \((g_2 > g_2^c, a^2 g_2 \sim O(1)\) eq. (20)), all doublers are supposed to be decoupled via eqs. (11,12) and (13,18). The critical point \( a^2 g_2^c = 0.124 \) is a rather small number. Thus, at \( a^2 g_2 \sim O(1) \), we introduce \( N_f \), an addition number of fermion flavours of \( \psi^i_L \) and \( \chi_R \) so that,

\[
a^2 g_2 > 0.124, \quad \tilde{g}_2 = a^2 g_2 N_f < 1 \quad \text{fixed,} \quad N_f = 3 \sim 8,
\]

where the value of \( N_f \) depends on the value of \( a^2 g_2 \) considered. Therefore, in a certain intermediate region of \( a^2 g_2 \sim O(1) \), we adopt the large-\( N_f \) technique to control the convergence of the approximation (see Fig.3) and calculate the Green function (52) of the charged three-fermion-operators for the low-energy \( p \sim 0 \). Hence we can get a qualitative insight into the charged low-energy spectrum \((p \sim 0)\) within the intermediate region of \( g_2 \) (53). We expect that the dynamics.
of the interaction is not greatly changed by introducing more flavours. As for $a^2 g_2 > 1$, such large-$N_f$ technique is doomed to fail.

Analogous to the case of neutral channel \cite{11}, the non-trivial leading order $O(N_f)$ contribution, as indicated by the first diagram in Fig.3, is $(p \sim 0)$,

$$W^c_{\circ}(p) = -N_f \int_{q_k} S_{RR}(p+q)R(k,q)\left(\frac{1}{2a}\right)^2.$$  \hspace{1cm} (54)

The second order $(a^2 g_2 N_f)^2$ contribution, as indicated by the second diagram in Fig.3, is $(p \sim 0)$,

$$W^c_1(p) = (g_2 N_f)^2 V^c(p)\tilde{S}_{LL}(p) V^c(p)$$  \hspace{1cm} (55)

where we ignore indices $ij$ and $V^c(p)$ is given as $(p \sim 0)$

$$V^c(p) = -\int q_k 4w(p+q)w(k-\frac{q}{2})S_{RR}(p+q)R(k,q),$$  \hspace{1cm} (56)

where the factor $4w(p+q)w(k + \frac{q}{2})$ comes from interacting vertices. The full propagator $\tilde{S}_{LL}(p)$ of $\psi^i_L(x)$ in eq.(55), as indicated by a full circle in middle of the second diagram in Fig.3, is calculated by the train approximation (as indicated in Fig.4) for the external total momentum $p \sim 0$,

$$\tilde{S}^{ij}_{LL}(p) = \int_x e^{-ipx}\langle \psi^i_L(0)\bar{\psi}^j_L(x)\rangle \simeq Z_2^{-1}(p)S^{ij}_{LL}(p).$$  \hspace{1cm} (57)

The wave-function renormalization $Z_2(p)$ of the elementary field $\psi^i_L(x)$ is given by (see Appendix A),

$$Z_2(p) = 1 + (\bar{g}_2)^2 \frac{N_f}{\int_k q S_{RR}(p+q)S_{LL}(p)R(k,q)}.$$  \hspace{1cm} (58)

Note that in eqs.(54,56,58) and $R(k,q)$ \cite{39}, the internal propagators $S_{LL}(k)$ for $\psi^i_L(x)$, and $S_{RR}(k)$ for $\chi_R(x)$ are respectively given by eqs.(37,38) or approximately by eqs.(14,13). The reasons for such choices are that in the region \cite{53}, doublers $(k \sim \pi_A)$ are supposed to be decoupled via eq.(14,13) and the internal momentum $k$ runs from $k \sim 0$ to $k \sim \pi_A$. As for the propagator $S_{LL}(p)$ in eqs.(57,58) for the external total momentum $p \sim 0$, we adopt eq.(14) by analytical continuation from $p \sim \pi_A$ to $p \sim 0$.

Analogous to that \cite{11} in the neutral channel, for the total momentum $p \sim 0$, the leading order contribution \cite{54} becomes \cite{35},

$$W^c_{\circ}(p) \sim a N_f (i\gamma_\mu p^\mu).$$  \hspace{1cm} (59)

As for the second order contribution \cite{53}, we find for $p \sim 0$,

$$V^c(p) \sim a (i\gamma_\mu p^\mu),$$  \hspace{1cm} (60)

$^5 \mu \sim \pi$ indicates up to a finite constant.
which cancels the pole at $p = 0$ stemming from the propagator $\tilde{S}_{LL}$ in eq.(57). And the wave-function renormalization $Z(p)$ of the elementary fields $\psi_L^i(x)$ at $p = 0$ is given by,

$$Z_2(0) = 1 + \text{const.},$$

(61)
due to eq.(39). This indicates that the relativistic particle $(p = 0)$ of $\psi_L^i$ receives a wave-function renormalization $Z_2(0)$. As a result, the second-order contribution (55) reads,

$$W_1^{\psi}(p) \sim (\tilde{g}_2)^2 a(i\gamma_{\mu}p^\mu).$$

(62)

Thus, for $p \sim 0$ the Green function (52) is approximately computed as,

$$\int_x e^{-ipx} \langle \Psi^c_R(0) \bar{\psi}_L^j(x) \rangle \simeq W_c^c(p) + W_1^{\psi}(p),$$

(63)

which is regular at $p \sim 0$. This implies that in this intermediate region (53), the charged channel of three-fermion-operators at $p \sim 0$ is not a simple pole for a massless relativistic particle with the “wrong” chirality, rather it is regular for a virtual state of three-fermion-cut. This agrees with the result obtained by analytical continuation of the propagator (23) from $p \sim \pi_A$ to $p \sim 0$.

An important consistent check is to examine the equation (55) for the “doublers”, i.e., the external momentum $p \sim \pi_A$. In fact, for $p \sim \pi_A$,

$$V^c(p) \simeq -4w(\pi_A) \int q \frac{w(k - q)}{2} S_{RR}(p + q) R(k, q).$$

(64)

This results in the coupling $\tilde{g}_2$ in the second-order contribution (54) being enhanced up a factor of $w^2(\pi_A)$ (see eq.(2)),

$$W_1^{\psi}(p) \sim (\tilde{g}_2 w(\pi_A))^2 V^c(p) \tilde{S}_{LL}(p) V^c(p).$$

(65)

Analogously, for $p \sim \pi_A$,

$$Z_2(p) \simeq 1 + \frac{2(\tilde{g}_2 w(\pi_A))^2}{N_f} \int_{k,q} \left( 4w(k - q) \right)^2 S_{RR}(p + q) S_{LL}(p) R(k, q).$$

(66)

The consequence is the complete breakdown of the large-$N_f$ expansion (53), which is not convergent for $p \sim \pi_A$, to calculate the two-point Green function (52) of charged three-fermion-operators. This implies that bound states (three-fermion-states) with the total momentum $p \sim \pi_A$ should be formed, consistently with the bound states that we find by the strong coupling expansion for $p \sim \pi_A$.

We turn to the computation of Green function for the mixing between the elementary field $\psi_L^i$ and charged three-fermion-state $\Psi^j_R$ (11). This mixing can be calculated by the large-$N_f$ expansion (54). The non-trivial leading order $O(N_f)$ is explicitly written as,

$$\int_x e^{-ipx} \langle \Psi^c_R(0) \bar{\psi}_L^j(x) \rangle = \int_x e^{-ipx} \left( \frac{1}{a} \right) \langle \psi_R(0) \bar{\psi}_L^j(x) \rangle \langle \psi_L^i(0) \psi_R(0) \rangle + O \left( \frac{\tilde{g}_2}{N_f} \right),$$

(67)
where \( O \left( \frac{g_2}{N_f} \right) \) stands for higher order contributions. Eq.\((67)\) vanishes for non spontaneous symmetry breaking (see eq.\((8)\)). Consistently, an analytical continuation of the Green function \((14)\) from \( p \sim \pi_A \) to \( p \sim 0 \) shows vanishing of the mixing at \( p \sim 0 \) as well.

Finally we check that in the intermediate region of \( g_2 \) \((53)\), whether the Green function \((57)\) for the elementary field \( \psi^i_L \) has a simple pole at \( p \sim 0 \) representing a relativistic massless particle. Due to \( Z_2(0) = 1 + \text{const.} \) \((11)\), one can check the propagator \((57)\) for \( \psi^i_L \) has a simple pole at \( p = 0 \), which indicates a massless, charged left-handed \( \psi^i_L \) in the low-energy spectrum. This agrees with the analytical continuation of the propagator \((14)\) from \( p \sim \pi_A \) to \( p \sim 0 \).

In the continuation from eq.\((23)\) for \( p \sim \pi_A \) to eq.\((63)\) for \( p \sim 0 \), we must meet the “dissolving” threshold \( \epsilon \), which should be the same as that in the neutral channel. We need to point out that the computation of charged channel is different from the computation of neutral channel. In the neutral channel, the Green functions \((13,17,23)\) for \( p \sim \pi_A \) and \((43,47,48)\) for \( p \sim 0 \) are both consistently calculated in \( g_2 \gg 1 \). However, in the charged channel, the Green functions \((57,58,67)\) for \( p \sim 0 \) are computed in the intermediate region \( a^2 g_2 \sim O(1) \) \((53)\), while the same Green functions \((14,23,10)\) for \( p \sim \pi_A \) are computed in the region \( g_2 \gg 1 \). This may raise a question whether the dynamics we explore for \( p \sim 0, a^2 g_2 \sim O(1) \) and for \( p \sim \pi_A, a^2 g_2 \gg 1 \) are consistent. We argue that such studies are qualitatively justified, since the computations in the strong coupling expansion with respect to doublers \((p \sim \pi_A)\) are valid as well for \( a^2 g_2 \sim O(1) \) \((20)\) as discussed in the section 1.

In conclusion, on the basis of approximate calculations of relevant two-point Green functions of elementary and composite three-fermion-operators with respect to doublers \( p \sim \pi_A \) and low-energy mode \( p \sim 0 \), we qualitatively explore a possible scaling window in segment \( \mathcal{A} \), where in the low-energy spectrum, the three-fermion-states with the “wrong” chirality turn into their corresponding virtual states, and elementary fermion states with the “right” chirality remain as massless states. Evidently, full non-perturbative numerical simulation to explore this scaling window is very inviting and necessary in particular for any solid conclusions in the charged channel.

5 Some remarks

In this paper, we have discussed the features of the spectrum of neutral and charged sectors appearing in segment \( \mathcal{A} \) \((4)\). It is interesting to point out that in this scenario, doublers \((p \sim \pi_A)\) are decoupled by a gauge-invariant mass term, while the low-energy modes with the “wrong” chirality are “decoupled” by the vanishing of their generalized form factor. However, it is still far from a definitive demonstration that chiral gauge theories in the low-energy limit can be achieved in this way, we need to have numerical simulations to show that this scenario is
indeed realized in segment \( \mathcal{A} \).

The whole spectrum in segment \( \mathcal{A} \) is gauge symmetric, and Ward identities of gauge symmetry are preserved. We can straightforwardly turn on the perturbative gauge interaction. By the strong coupling expansion in powers of \( \frac{1}{g^2} \), we compute three-point vertex function \( \langle \Psi_R^i(x) \bar{\Psi}_R^j(y) A_\mu(z) \rangle \) and obtain the vertex of the \( SU_L(2) \) gauge field coupling to the charged three-fermion-state (the leading order of gauge coupling \( O(g) \)),

\[
\Lambda_{\mu RR}^{(1)}(p, p') = ig \frac{\tau^a}{2} \gamma_\mu P_R \cos \left( \frac{p + p'}{2} \right),
\]

where the momenta of three-fermion-states \( p, p' \sim \pi_A \). According the renormalization (27) of truncated Green functions with two three-fermion-operator insertions, the vertex of gauge coupling to the three-fermion-states is given by (Fig.5),

\[
\Lambda_{\mu RR}^{(1)}(p, p') = ig \frac{\tau^a}{2} \gamma_\mu P_R \cos \left( \frac{p + p'}{2} \right) Z_R(p) Z_R(p').
\]

For \( p \sim \pi_A \) and \( p' \sim \pi_A \), \( Z_R \)'s are positive definite constants, we thus renormalize the wave functions with respect to each doubler (27), as discussed in section 2. As a result, gauge coupling vertex \( (69) \) turns into \( (68) \). Although eq. (69) is obtained for \( p, p' \sim \pi_A \) and \( a^2 g^2 \gg 1 \), it can be analytically continued to \( p, p' \sim 0 \). We find that in the limit of \( p \to 0 \) and \( p' \to 0 \), the coupling vertex \( (69) \) of three-fermion-operators and gauge boson vanishes as \( O(p^8) \). This consistently corresponds to the dissolving of three-fermion-states into three-fermion-cuts in the low-energy limit. Since the propagator of charged three-fermion-states positively vanishes, there are no ghost states with negative norm coupling to gauge field through the Ward identity stemming from the gauge symmetry in segment \( \mathcal{A} \).

The model presented in this paper cannot be considered as a realistic model reflecting all aspects of chiral gauge theories. We need to completely understand the relationship between the anomaly-free condition and the realization of such dynamics discussed in the paper. We also need to understand what kind role of 't Hooft condition for anomaly matching\([35]\), fermion-number violation and Witten’s \( SU(2) \) global anomaly\([36]\) would play in such dynamics.

In this approach, the left-handed field \( \psi_L^i(x) \) is the doublet of the \( SU_L(2) \) chiral gauge group. However, it can be generalized to be the left-handed field (complex representations) of any anomaly-free chiral gauge group (e.g. \( SU(5) \) and \( SO(10) \)). A right-handed field (spectator) \( \chi_R \), that is a singlet of the chiral gauge group, can be introduced and coupled to the left-handed field in the same way as the multifermion couplings given in (1). As for the right-handed field \( \psi_R^i \) of chiral gauge groups, we can analogously introduce a left-handed spectator field \( \chi_L \) that is a singlet of the chiral gauge groups, and couple it to the right-handed field \( \psi_R^i \) in the same way as the multifermion couplings in the action (1) with \( L \leftrightarrow R \). This indicates that such a formulation of chiral gauge theories is actually quite general.
To be more specific, we take the anomaly-free chiral gauge group of the Standard Model as an example. In this realistic case, $\psi^i_L(x)$ can be both left-handed lepton doublets and left-handed quark doublets. The candidate for a right-handed spectator field $\chi_R$ could be the right-handed neutrino $\nu_R$. As for the right-handed fields $\psi^i_R$ with respect to $U_Y(1)$, we can introduce an additional left-handed spectator field $\chi_L$ (a $SU_L(2) \otimes U_Y(1)$ singlet) coupling to the right-handed fields $\psi^i_R$ as that in eq.(1). These spectator fields $\nu_R$ and $\chi_L$ are free and decouple from other particles due to the $\nu_L$- and $\chi_L$-shift-symmetry acting on them. In this way, we can in principle have a gauge-invariant formulation of the Standard Model on the lattice. In practice, non-perturbative analysis, which can be done analytically to show whether such a formulation gives the low-energy Standard Model, should be more or less similar to (certainly more complicated than) that discussed in this paper and references[17, 18]. However, the spectator field $\chi_L$ might not be necessary. Alternatively, the ’t Hooft vertex in the lattice formulation of the Standard Model[14] provides a scenario in which fermion-number conservation is explicitly violated, and all chiral fermions find their patterns with opposite chirality within the Standard Model. The formulation of a realistic Standard Model with multifermion couplings and all features of dynamics discussed in this paper are an extremely interesting subject for future studies.

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Appendix A

The propagator of $\psi^i_L$ is given by eq.(14) ($p \sim 0$),

$$S^{ij}_{LL}(p) = \delta_{ij}P_LP_R,$$

$$\hat{p} = \frac{\gamma^\mu \sin p_\mu}{\frac{4}{a^2} \sum_\mu \sin^2 p_\mu + M^2(p)}.$$ (70)

The Feynman diagram (see Fig. 6) is given by,

$$\sigma(p) = -N_f \int_{qk} \lambda S_{RR}(p+q) \text{tr} \left[ S_{RR}(k - \frac{q}{2}) S_{LL}(k + \frac{q}{2}) \right]$$

$$= -N_f \int_{k,q} \lambda S_{RR}(p+q) R(k,q),$$ (71)

where $S_{LL}, S_{RR}$ are (14,15) and $R(k,q)$ is eq. (39) and

$$\lambda = \left(4g_2w(p-k)w(k + \frac{q}{2})\right)^2 = \frac{1}{N_f^2} a^{-4} \left(4g_2w(p-k)w(k + \frac{q}{2})\right)^2.$$ (72)
The wave-function renormalization $Z_2(p)$ of $\psi^i_L(x)$ in eq. (57) can be calculated by using the train approximation (see Fig. 4),

$$Z_2^{-1}S^{ij}_{LL}(p) = P_L(\hat{p} + \hat{p}\sigma\hat{p} + \hat{p}\sigma\hat{p}\sigma\hat{p} + \cdots)P_R\delta^{ij}$$

$$= S^{ij}_{LL}(p) \left( \frac{1}{1 - \sigma\hat{p}} \right),$$

(73)

and one gets

$$Z_2 = 1 - \sigma\hat{p}.$$  

(74)

With eq. (71) one can get

$$\sigma\hat{p} = -\frac{1}{N_f} \int_{k,q} \left( 4\tilde{g}_2 w(p-k)w(k+q/2) \right)^2 S_{RR}(p+q)S_{LL}(p)R(k,q).$$

(75)

By substituting eq. (73) into (74), one gets eq. (58).

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**Figure Captions**

**Figure 1:** The two-point Green functions of composite three-fermion-operators at strong coupling $g_2 \gg 1$ for $p \sim \pi_A$, indicating three-fermion-states.

**Figure 2:** 1PI truncated Green functions of an elementary field and one three-fermion-operator insertion, i.e., the generalized form factors of three-fermion-states.

**Figure 3:** The two-point Green functions of composite three-fermion-operators in the effective weak coupling for $p \sim 0$, indicating three-fermion-cuts.
Figure 4: The train approximation to the propagators of the elementary fields $\psi^i_L$ and $\chi^i_R$.

Figure 5: Gauge boson coupling to three-fermion-states.

Figure 6: A single bubble diagram in the weak-coupling expansion (see Figs. 3, 4).