Nuclear axial currents from scale-chiral effective field theory *

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Abstract: By incorporating hidden scale symmetry and hidden local symmetry in the nuclear effective field theory, combined with the double soft-pion theorem, we predict that the Gamow-Teller operator coming from the space component of the axial current should remain unaffected by the QCD vacuum change caused by the baryonic density, whereas the first forbidden beta transition operator coming from the time component should be strongly enhanced. While the latter has been confirmed for some time, the former was given support by a powerful recent *ab initio* quantum Monte Carlo calculation for light nuclei, which also confirmed the old "chiral filter hypothesis." Formulated in terms of the Fermi-liquid fixed point structure of strong-coupled nuclear interactions, we offer an extremely simple resolution to the long-standing puzzle of the "quenched g_A," $g_A^{\rm eff} \approx 1$ [1], found in nuclear Gamow-Teller beta transitions, giant Gamow-Teller resonances, and double beta decays.

Keywords: effective field theory, hidden symmetry, fermi-liquid fixed point, quenched g_A

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1 Introduction

The behavior of the axial-vector coupling constant g_A in a nuclear medium has a long history of puzzles in the fields of nuclear physics, astrophysics, and particle physics. In nuclear physics, there has been the mysterious ~20% quenching of g_A in the shell-model calculations of nuclear beta decay, giant Gamow-Teller resonances, and double beta decay. In particle physics, there is the issue of the partial restoration of the chiral symmetry, which is an intrinsic property of the symmetry of quantum chromodynamics (QCD). Finally, the surprising role of the first-forbidden beta decay in nucleosynthesis can be found in astrophysics. Some of these issues were comprehensively reviewed in [2].

In this Letter, we propose a simple resolution of the g_A problem based on the scale-invariant chiral effective field theory combined with the "chiral filter" mechanism anchored on the current-algebra soft-pion theorems[1)]. We will argue that the recent *ab initio* quantum Monte Carlo calculation by Pastore et al. [3] goes a long way in giving support to our simple solution.

2 Scale-chiral symmetry

Among the symmetries that are most relevant to nuclear dynamics, QCD has chiral symmetry, which is explicitly broken by the quark mass, and scale symmetry, which is explicitly broken by the trace anomaly.

How chiral symmetry figures in the nuclear effective field theory is now well understood in the guise of the chiral perturbation theory and has been fairly well-established in the modern development of *ab initio* approaches since the paper by Weinberg [4]. Because the quark masses involved in nuclear dynamics are tiny compared with the chiral scale, ~1 GeV, one can talk about the "chiral limit," where one approximates by setting the quark mass equal to zero. In doing the calculations, it makes sense to theoretically "turn off" the chiral symmetry explicit breaking. When the chiral symmetry is spontaneously broken, Nambu-Goldstone bosons emerge, which are the well-known pions π.

In contrast, while QCD has classical scale invariance in the chiral limit, quantum mechanically, there is an anomaly, called the trace anomaly, which posits a scale, and hence the scale symmetry is explicitly broken. This anomaly is renormalization-group invariant, and hence cannot be "turned off" like the quark mass. Thus, there does not seem to be any point in discussing scale-invariance as is done with chiral symmetry. The explicitly broken scale symmetry can also be broken sponta-

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1) The soft-pion theorem that is exploited here is *trivially* encoded in the modern chiral perturbation theory, but seems to have a deep theoretical connection with the infrared structure of gravity, as well as the gauge theories. This point will be briefly mentioned at the end of this paper.

neously, generating a Goldstone boson, which is intrinsically massive as a result of the trace anomaly. Thus, it is a pseudo-Nambu-Goldstone boson, called a dilaton. The scale symmetry spontaneous breaking has subtlety, because unlike chiral symmetry, scale symmetry cannot be broken spontaneously but must be explicitly broken [5]. This raises a conundrum when introducing scale symmetry, i.e., the scalar meson dilaton, in nuclear physics [6].

In nuclear physics, there is a dire need for a local scalar field of mass for ∼600 MeV. Such a scalar plays an important role for the nuclear forces, as well as Walecka-type mean field approaches (i.e., an energy density functional) to nuclear matter, which is popular in nuclear physics. The particle data booklet includes a low-mass scalar $f_0(500)$, and an attractive possibility is identifying it as the dilaton. This is the proposal made by Crewther and Tunstall (CT) [7]. A similar idea was proposed by the authors of [8].

We will follow here the approach by CT. In CT, the presence of an infrared fixed point in the QCD $\beta(\alpha_s)$ function is postulated, i.e., $\beta(\alpha_{IR})=0$. The difficulty is that it is not known whether such an IR fixed point is present in the QCD in the vacuum. There is no indication either for or against such an IR fixed point when the number of flavors N_f is less than ~ 8[1)]. This issue cannot be settled at the moment, as explained in detail in [6]. Here, we will bypass that conundrum by assuming that although it may not make sense in the vacuum, one can consider approaching the vicinity of an IR fixed point in a medium and study the fluctuations around, but not on top of, a potential IR fixed point. This is the standpoint we take here. Based on this assumption, we will make certain predictions to be confronted with nature.

There are two predictions in particular that are relevant to nuclear physics. One is the prediction of the properties of compact stars, which involves highly dense matter. Reference [9] addresses this issue. We will not go into this matter. The other is the g_A problem that we are interested in here.

As Yamawaki has argued [10], the scale symmetry that we are considering is present – though "hidden" – in a linear sigma model. Starting with a linear sigma model – to which the standard Higgs model belongs, changing a parameter in the model has shown that the linear sigma model can be driven to the familiar nonlinear sigma model, on which the chiral perturbation theory is built,

$$\mathcal{L}_{NL\sigma}=\frac{f_\pi^2}{4}\mathrm{Tr}\left(\partial_\mu U\partial^\mu U^\dagger\right)+\cdots \qquad (1)$$

where f_π is the pion decay constant, or to a scale invariant model with a dilaton coupled to a nonlinear sigma field with a dilaton potential that breaks the scale symmetry. It is the latter form that is relevant to us, and we suggest that it is the baryonic density that drives the coupling. In the chiral limit, it has the form

$$\mathcal{L}_{\mathrm{Scale}\sigma} = \mathcal{L}_{\mathrm{sinv}}-V(\chi) \qquad (2)$$

with

$$\mathcal{L}_{\mathrm{sinv}} = \frac{1}{2}(\partial_\mu\chi)^2+\frac{f_\pi^2}{4}\left(\frac{\chi}{f_\sigma}\right)^2\cdot\mathrm{Tr}\left(\partial_\mu U\partial^\mu U^\dagger\right)+\cdots. \qquad (3)$$

Here, f_σ is the σ decay constant; U is the usual chiral field, which is a scale singlet; and χ is the mass dimension 1 "conformal compensator field" $\chi=f_\sigma e^{\sigma/f_\sigma}$, where σ is the dilaton field, which is transformed into a singlet under chiral transformation and scale dimension 1 under scale transformation. The ellipsis stands for higher order terms. As given, the Lagrangian (3) is scale invariant and chiral invariant. Both the explicit and spontaneous scale symmetry breaking are put in the dilaton potential $V(\chi)$, which does not explicitly relate to the process with which we are concerned.

There is another hidden symmetry in the chiral Lagrangian that plays a role that is just as important as the scalar dilaton χ, and it involves the vector mesons ρ and ω. Because the mass involved is comparable to that of the scalar of ∼ 700 MeV, they need to be incorporated together. In fact the ω is essential in the Walecka-type relativistic mean field approach, providing the necessary repulsion, while the ρ comes into the nuclear tensor forces with a sign that is the opposite to that of the pion tensor. The symmetry associated with the vector mesons is the local gauge symmetry, and hence what is involved is the hidden local symmetry (HLS) [11]. This symmetry can be easily implemented to the scale-symmetric Lagrangian by exploiting the redundancy present in chiral field U to make the Lagrangian (3) hidden local symmetric. The resulting Lagrangian is scale-symmetric HLS [12], or sHLS for short.

Finally, it is necessary to consider how to set up the power series of the scale symmetry in conjunction with the chiral symmetry, the power counting of which is well established. Following the idea of CT, we consider expanding around the assumed IR fixed point at which the beta function is zero, $\beta(\alpha_{IR})=0$,

$$\beta(\alpha_s)=\delta\cdot\beta'+\cdots \qquad (4)$$

where

$$\delta = (\alpha_s-\alpha_{IR}) \qquad (5)$$

$$\beta' = \frac{\partial}{\partial\alpha_s}\beta(\alpha_s)\Big|_{\alpha_s=\alpha_{IR}} \qquad (6)$$

is the anomalous dimension of $G_{\mu\nu}^2$, with $G_{\mu\nu}$ being the gluon energy momentum tensor. The power counting

1) There is lattice indication for the presence at A large N_f near what is called the "conformal window," which has a connection with Higgs physics.

of the power expansion in the scale-chiral perturbation theory is then

$$O(\partial^2)\sim O(p^2)\sim O(m_\pi^2)\sim O(\delta). \tag{7}$$

Here, π stands for the octet pseudo-scalar NG bosons with momentum p. Note that β' signaling explicit scale symmetry breaking is $O(1)$ in the scale-chiral counting, in contrast to the chiral symmetry, where the chiral symmetry explicit breaking quark mass is counted as $O(p^2)$.

The above counting rule was recently implemented in deriving a scale-chiral expansion that incorporated both hidden symmetries, and the detailed discussions are given in [12]. The formalism is applied in [13] to derive Brown-Rho scaling in a medium, which is valid in the density regime up to $\sim 2n_0$, where n_0 is the normal nuclear matter density[1)].

The expressions of the Lagrangian beyond the leading order, namely, $O(p)$ in the baryon sector and $O(p^2)$ in the meson sector, which were given in [12, 13], are very complicated and involve a large number of unknown parameters. However, for the leading order (LO) in scale symmetry, it is simple to reproduce what was given in [12, 13]. Suppose one has a Lagrangian $\mathcal{L}^{(m)}(\Phi)$ involving Φ fields (baryons B, mesons ρ, ω, π) and the total scale dimension in the Lagrangian density is $m\leqslant 4$. Then, one makes the Lagrangian density have scale dimension 4, so that the action is scale-invariant, by multiplying it by the conformal compensator field χ as follows:

$$\bar{\mathcal{L}}=\left(\frac{\chi}{f_\sigma}\right)^{4-m}\mathcal{L}^{(m)}. \tag{8}$$

Then, the CT procedure is to write

$$\bar{\mathcal{L}}\to\left(\kappa+(1-\kappa)\left(\frac{\chi}{f_\sigma}\right)^{\beta'}\right)\bar{\mathcal{L}}, \tag{9}$$

where κ is an unknown constant. Now, when the dilaton field is turned off by setting $\sigma=0$, the β' dependence disappears. This will give the usual chiral perturbation theory, HLS, if vector fields are included. Then, there will be no footprint of scale symmetry breaking in it.

There are two ways that the CT Lagrangian reduces to the form of the hidden scale symmetric sHLS. One is that $\beta'\ll 1$, that is, weak explicit scale symmetry breaking. This is somewhat similar to the situation of a kaon mass with a dilaton mass of the same size. As there, perturbation expansion in β' could make sense. Another possibility is that $\kappa\approx 1$. This seems to be more consistent with the notion that the scale symmetry is hidden and in fact is favored in the treatment of compact stars [9], where scale invariance is considered to be "emergent" or un-hidden at high density. When expanded to higher orders in β', the physics may be quite different, but at this order, the resulting Lagrangian, with the explicit scale symmetry breaking entirely in the dilaton potential, can have an analogy to the usual chiral Lagrangian, where the explicit symmetry breaking is put entirely in the quark mass term. We shall call this the leading order scale symmetry (or LOSS) Lagrangian.

3 Nuclear axial currents

For our consideration of nuclear processes, we can restrict ourselves to chiral $SU(2)\times SU(2)$. Reducing from the three flavors for which the scale-chiral EFT is formulated [12, 13] to two flavors, one can extract the relevant part of the Lagrangian, which is found to be extremely simple:

$$\mathcal{L}= i\overline{N}\gamma^\mu\partial_\mu N-\frac{\chi}{f_\sigma}m_N\overline{N}N+g_A\overline{N}\gamma^\mu\gamma_5\tau_a N\mathcal{A}_\mu^a+\cdots, \tag{10}$$

where $\mathcal{A}_\mu$ is the external axial field. Note that while the kinetic energy term and particularly the nucleon coupling to the axial field are scale-invariant by themselves, and hence do not couple to the conformal compensator field, the nucleon mass term is multiplied by it. Put in the nuclear matter background, the bare parameters of the Lagrangian will pick up the medium VeV. Thus, in (10), the nucleon mass parameter will scale in density, while, significantly, g_A will remain unscaled:

$$m_N^*/m_N = \langle\chi\rangle^*/f_\sigma\equiv\Phi \tag{11}$$

$$g_A^*/g_A = 1 \tag{12}$$

where f_σ is the medium-free VeV $\langle\chi\rangle_0$ – sigma decay constant– and $*$ represents the medium quantities. The first is one of the scaling relations given in [14]. The dilaton condensate carries density dependence when the vacuum is warped by density. It is a part of the "induced density dependence (IDD)" inherited from the QCD valid at $n\lesssim 2n_0$ [9]. The second is new and says that the Lorentz-invariant axial coupling constant does not scale in density. This result was already indicated in the Skyrme term of the Skyrme model [14], but what is given here is more directly linked to the QCD symmetries. Combined with the "chiral filtering mechanism" to be specified below, this is the most important point in the present report.

The Lorentz invariance is spontaneously broken in a medium, which means space component $g_A^{\rm s}$ and time component $g_A^{\rm t}$ could be different. Indeed, writing out the space and time components of the nuclear axial current

1) Beyond $n_{1/2}\sim 2n_0$, where the skyrmion-half-skyrmion topology change takes place, the intrinsic density dependence coming from the matching of the correlators of the EFT and QCD at a matching scale, rather than the operative at density $n\lesssim 2n_0$, must be taken into account.

operators, one obtains the following:

$$\vec{J}_A^{\pm}(\vec{x}) = g_A^{\rm s}\sum_i \tau_i^{\pm}\vec{\sigma}_i\delta(\vec{x}-\vec{x}_i), \quad (13)$$

$$J_5^{0\pm}(\vec{x}) = -g_A^{\rm t}\sum_i \tau_i^{\pm}\vec{\sigma}_i\cdot(\vec{p}_i-\vec{k}/2)/m_N\delta(\vec{x}-\vec{x}_i), \quad (14)$$

where $\vec{p}$ is the initial momentum of the nucleon making the transition, and $\vec{k}$ is the momentum carried by the axial current. In writing (13) and (14), a nonrelativistic approximation is made for the nucleon. This approximation is valid not only near n_0 but also in the density regime $n \gtrsim n_{1/2} \sim 2n_0$. This is because the nucleon mass never decreases much after the parity-doubling sets in at $n \sim n_{1/2}$, at which $m_N^* \to m_0 \approx (0.6-0.9)m_N$ [9]. It will be related to the pion decay constants, as shown below.

A simple calculation that takes into account (11) and (12) gives

$$g_A^{\rm s} = g_A, \quad g_A^{\rm t} = g_A/\Phi \quad (15)$$

with Φ given by (11).

4 Chiral filtering effect

To confront with nature, we need two ingredients: (1) accurate nuclear wave functions and (2) reliable nuclear weak currents. In systematic EFT calculations, the two are treated on the same footing. To proceed, let us suppose that the wave functions are accurately calculable with an accurate potential. Regarding point (2), the soft-pion theorems are crucial. In the scale-chiral counting with the scheme espoused in this paper (which is essentially equivalent to chiral counting [15]), taking the axial current to be a "soft pion," there is a soft-pion exchange current that involves double soft-pions coupling to the nucleon dictated by the current algebras. This term was shown to be the most important exchange current contribution to axial current transitions in nuclei. This was first shown in 1978 using soft-pion theorems [16] and in 1991 using the chiral perturbation theory [17]. The phenomenon was dubbed the "chiral filter hypothesis" because at the time the high-powered computational techniques that have since been developed were not available to quantitatively check the arguments believed to be reasonable. It was predicted that (a) soft pions would cause a huge meson exchange correction to the one-body charge operator (14) governing the first-forbidden beta transitions, with the higher chiral corrections strongly suppressed (by three chiral orders) and (b) that Gamow-Teller transitions (in stark contrast, these would be unprotected by chiral filtering with the soft pions suppressed) would predominantly be caused by the leading order one-body operator (13), and the exchange current corrections, coming at chiral orders of three and higher, could not be reliably computed with only a few terms aided by the chiral symmetry. It can involve mixing to the states of high excitation energies, ~300 MeV, such as highly correlated nuclear states and Δ-hole states strongly coupled by the nuclear tensor force.

The enhancement factor for the axial-charge operator is very simple to calculate. It involves the soft-pion exchange. The ratio of the two-body over one-body matrix elements R can be computed almost independently of the nuclear model [18, 19]. It comes out to be $R = 0.5 \pm 0.1$ ranging from $A = 12$ to $A = 208$. An extremely simple calculation shows that with the two-body effect taken into account, the effective axial-charge operator is obtained by making a substitution in (14) using [20] $g_A^{\rm t} \to g_A^{\rm t*} = \epsilon g_A$, with $\epsilon = \Phi^{-1}(1+R/\Phi)$. The scaling factor Φ is related to the pion decay constant in medium f_π^* as $\Phi \approx f_\pi^*/f_\pi$ [9]. At the nuclear matter density, one gets $\Phi(n_0) \approx 0.8$ from deeply bound pionic nuclei [21]. The enhancement factor at the nuclear matter density is then $\epsilon(n_0) \approx 2.0 \pm 0.2$, within the range of theoretical uncertainty in ratio R. This is confirmed by what was found in Pb nuclei [22], $\epsilon^{\rm exp}(n_0) = 2.01 \pm 0.05$. The results for $A = 12, 16$ [23] are compatible with the Pb result. It is an understatement to say that this is a gigantic correction as an exchange current effect[1)].

Now we turn to the other aspect of the chiral filter mechanism, i.e., the Gamow-Teller coupling constant. Here, the soft-pions are rendered powerless. Hence, whatever corrections occur for the leading one-body operator must be suppressed relative to the leading $O(1)$ operator, accounting for the "chiral filtering." The quantum Monte Carlo calculations for $A = 6-10$ nuclei made by Pastore et al. verified this feature at the $\rm N^4LO$. At that order, the filter seems to work remarkably well, say, at the level of $\lesssim 3\%$. It is even more striking that with high-order correlations accounted for in the Monte Carlo approach, there is no indication of g_A quenching, which means the high-order nuclear correlations in the wavefunctions, rather than a basic modification of the axial current, could have been responsible for the "g_A problem." If that is the case, then this calculation provides support for the prediction that g_A should not be affected by a decrease in the chiral condensate with density, as given by the following: (15)[2)]. This may appear somewhat surprising because, as is generally accepted,

1) Needless to say, *ab initio* high-powered calculations on this matter for light nuclei would be highly desirable to further confirm this prediction.

2) Without going into details that involve the behavior of the nuclear tensor force mediated by the ρ exchange in the sHLS Lagrangian with baryons implemented [9], which the lack of space does not allow us to go into here, we should point out that nontrivial support for the thesis developed here is provided by the highly successful explanation by Holt et al. [24] of the strongly suppressed Gamow-Teller matrix element in the C14 dating process with the "Brown-Rho scaling" but with the exchange-current contributions totally ignored.

the pion decay constant should follow in some way the chiral condensate, which is considered to decrease with an increase in the temperature or density, going to zero at the chiral restoration, whereas the axial current that crucially couples to the pions does not change with the density. In the same vein, it has been considered plausible that the axial coupling constant must approach 1 as chiral symmetry is realized in the Wigner-Weyl mode, as is indicated at the dilaton-limit fixed point [9]. The prediction made in this paper and the powerful *ab initio* quantum Monte Carlo calculation indicate that this is not the case.

5 Landau-fermi fixed point $g_A^{\rm L}$

Given the above explanation of where the quenched g_A is located, the question that remains is why g_A is quenched "universally" by $\sim 20\%$ in the nuclear shell model calculations [1].

We offer an extremely simple answer in terms of the Landau Fermi-liquid fixed point theory using the scale-chiral EFT Lagrangian, sHLS. The key point is that the mean-field approximation with the sHLS Lagrangian endowed with the IDDs inherited from QCD corresponds to the Landau-Fermi liquid fixed point theory, similar to the Wilsonian renormalization group to many-fermion systems with a Fermi surface [9, 25]. In the large N limit, where $N{=}k_{\rm F}/(\Lambda{-}k_{\rm F})$ and Λ is the cutoff on top of the Fermi surface, the Landau mass $m_{\rm L}$ and quasiparticle interactions $\mathcal{F}$ are at the fixed point, with $1/N$ corrections suppressed [26]. The relation between the Landau mass $m_{\rm L}$ and the effective $g_A^{\rm L}$, both taken at the fixed point, is given by [14, 25]

$$\frac{m_{\rm L}}{m_N}=1+\frac{1}{F_1}=\left(1-\frac{\tilde{F}_1}{3}\right)^{-1}\approx\Phi\sqrt{\frac{g_A^{\rm L}}{g_A}}, \tag{16}$$

where $\tilde{F}_1$ is related to Landau parameter F_1 by $\tilde{F}_1=(m_N/m_{\rm L})F_1$. Applying the mean field argument, this relation gives

$$\frac{g_A^{\rm L}}{g_A}\approx\left(1-\frac{1}{3}\Phi\tilde{F}_1^{\pi}\right)^{-2}, \tag{17}$$

where $\tilde{F}_1^{\pi}$ is the pion Fock term contribution to the Landau parameter $\tilde{F}_1$. The Fock term is a loop contribution, so naively $O(1/N)$. However, because the pion is "soft," it plays an indispensable role, just as it does for the anomalous orbital gyromagnetic ratio δg_l^p [25]. It is noteworthy that both δg_l^p and $g_A^{\rm L}$ depend only on Φ and $\tilde{F_1}^{\pi}$.

Let us consider the $g_A^{\rm L}$ at the nuclear matter density. With $\Phi(n_0)\approx 0.8$ inferred from deeply bound pionic systems [21] and $\frac{1}{3}\tilde{F}_1^{\pi}(n_0)=-0.153$, which is precisely given by the pion exchange, we get, with a current value of $g_A{=}1.27$,

$$g_A^{\rm L}(n_0)\approx 0.79 g_A\approx 1.0. \tag{18}$$

This is the precise $g_A^{\rm eff}$ needed in the shell-model calculations [1] and in the giant Gamow-Teller resonances [27]. Note here the crucial role of the pionic contribution interlocked with the dilaton condensate for the quenching. It turns out that the density dependence of Φ (which decreased with the density) nearly cancels the density dependence of $\tilde{F}_1^{\pi}$ (which increases with the density). Thus, the product $\Phi\tilde{F}_1^{\pi}$ becomes more or less independent of the density. The values of g_A^L differ by less than 2% between densities of $\frac{1}{2}n_0$ and n_0. Thus, the Landau g_A (18), although evaluated for nuclear matter, is robust and could be applied to nuclear matter as well as finite nuclei, both light and heavy.

How does this $g_A^{\rm L}$ correspond to $g_A^{\rm eff}$ in the shell-model calculations?

To answer this question, recall that at the Fermi-liquid fixed point in our formulation, the beta functions for the quasiparticle interactions $\mathcal{F}$, mass $m_{\rm L}$, Gamow-Teller coupling $g_A^{\rm L}$, etc. at a given density should be suppressed. This means in particular that the quasiparticle loop corrections of the effective g_A should be suppressed. It is therefore the effective coupling constant, duly implemented with density-dependent condensates inherited from QCD and with high-order quasiparticle correlations subsumed, that is applied to non-interacting quasiparticles, that is, simple particle-hole configurations in shell model calculations. This corresponds effectively to what is captured microscopically in the *ab initio* quantum Monte Carlo calculation of [3].

6 "Soft-pion theorem triangle"

The prominent role played by soft pions in the previously addressed processes raises a potentially deep issue in nuclear physics. According to the lore of effective quantum field theory, it makes good sense to integrate out the pion for processes involving energy scales much lower than the pion mass, leading to what is known as pionless ($\not\pi$) EFT. It turns out that such a $\not\pi$EFT works fairly well in numerous low-energy nuclear processes. Now the question is how the soft-pion effect, which is crucial in certain nuclear processes such as first-forbidden beta transitions, can manifest when pion fields do not figure explicitly? The interplay between the in-medium vacuum condensate Φ and the pionic Landau parameter F_1^{π} is mysterious. This intriguing question may have an answer in the recent development involving soft theorems in the web of triangles "echoed" in a variety of infrared structures in the gauge – and gravity – theories [28]. It may act as a sort of "memory effect" in the triangle with the soft-pion theorem sitting on one corner.

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