CHIRAL SYMMETRY RESTORATION IN HADRON SPECTRA

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Abstract

The evidence and the theoretical justification of chiral symmetry restoration in high-lying hadrons is presented.

I. CHIRAL SYMMETRY RESTORATION OF THE SECOND KIND

It has recently been suggested that the parity doublet structure seen in the spectrum of highly excited baryons may be due to effective chiral symmetry restoration for these states [1]. This phenomenon can be understood in very general terms from the validity of the operator product expansion (OPE) in QCD at large space-like momenta and the validity of the dispersion relation for the two-point correlator, which connects the spacelike and timelike regions (i.e. the validity of Källen-Lehmann representation) [2,3].

Consider a two-point correlator $\Pi_{J_\alpha}$ of the current $J_\alpha(x)$ (that creates from the vacuum the hadrons with the quantum numbers $\alpha$) at large spacelike momenta $Q^2$, where the language of quarks and gluons is adequate and where the OPE is valid. The only effect that chiral symmetry breaking can have on the correlator is through the nonzero value of condensates associated with operators which are chirally active (i.e. which transform nontrivially under chiral transformations). To these belong $\langle \bar{q}q \rangle$ and higher dimensional condensates that are not invariant under axial transformation. At large $Q^2$ only a small number of condensates need be retained to get an accurate description of the correlator. Contributions of these condensates are suppressed by inverse powers of $Q^2$. At asymptotically high $Q^2$, the correlator is well described by a single term—the perturbative term. The essential thing to note from this OPE analysis is that the perturbative contribution knows nothing about chiral symmetry breaking as it contains no chirally nontrivial condensates. In other words, though the chiral symmetry is broken in the vacuum and all chiral noninvariant condensates are not zero, their influence on the correlator at asymptotically high $Q^2$ vanishes. This is in contrast to the situation of low values of $Q^2$, where the role of chiral condensates is crucial.

This shows that at large spacelike momenta the correlation function becomes chirally symmetric. In other words, two correlators $\Pi_{J_1}(Q^2)$ and $\Pi_{J_2}(Q^2)$, where $J_1 = UJ_2U^\dagger$, 

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$U \in SU(2)_L \times SU(2)_R$, become essentially the same at large $Q^2$. The dispersion relation provides a connection between the spacelike and timelike domains. In particular, the large $Q^2$ correlator is completely dominated by the large $s$ spectral density $\rho(s)$. (The spectral density has the physical interpretation of being proportional to the probability density that the current when acting on the vacuum creates a state of a mass of $\sqrt{s}$.) Hence the large $s$ spectral density must be insensitive to the chiral symmetry breaking in the vacuum. I.e. $\rho_1(s)$ and $\rho_2(s)$ must coincide at asymptotically large $s$. This is in contrast to the low $s$ spectral functions which are crucially dependent on the quark condensates in the vacuum. This manifests a smooth chiral symmetry restoration from the low-lying spectrum, where the chiral symmetry breaking in the vacuum is crucial for physics, to the high-lying spectrum, where chiral symmetry breaking becomes irrelevant and the spectrum is chirally symmetric.

Microscopically this is because the typical momenta of valence quarks should increase higher in the spectrum and once it is high enough the valence quarks decouple from the chiral condensates of the QCD vacuum and the dynamical (quasiparticle or constituent) mass of quarks drops off and the chiral symmetry gets restored. This phenomenon does not mean that the spontaneous breaking of chiral symmetry in the QCD vacuum disappears, but rather that the chiral asymmetry of the vacuum becomes irrelevant sufficiently high in the spectrum. The physics of the highly excited states is such as if there were no chiral symmetry breaking in the vacuum. One of the consequences is that the concept of constituent quarks (which are directly related to the quark condensates of the QCD vacuum), which is adequate low in the spectrum, becomes irrelevant high in the spectrum.

How should one refer to such a phenomenon? Typically under chiral symmetry restoration people understand that the chiral properties of the vacuum are changed with temperature or/and density. E.g. at critical temperature the phase transition occurs and the quark vacuum becomes trivial: all the quark condensates of the vacuum vanish. On the contrary, in our case one does not affect the QCD vacuum by exciting the hadrons. The symmetry restoration is achieved via different mechanism. Namely, by exciting the hadrons one decouples the valence degrees of freedom from the QCD vacuum. So even if the quark condensates of the vacuum (which break the chiral symmetry) are not zero, their role gets smaller and smaller once we go up in the spectrum and smoothly the chiral symmetry gets restored. We can refer to such a phenomenon as chiral symmetry restoration of the second kind.

All this is in a very good analogy with the similar phenomenon in condensed matter physics. Consider a metal which is in a superconducting phase. There is a condensation of the Cooper pairs - which is analogous to the condensation of right-left quark pairs in the QCD vacuum - and as a consequence the low-lying excitations of the system are the excitations of quasiparticles - which are analogous to the constituent quarks in the low-lying hadrons. Now we want to study this superconductor by external probe, e.g. by photons. If we probe the superconductor by the low-energy photons $\hbar \omega \sim \Delta$, then the condensation of the Cooper pairs and the quasiparticle structure of the low-lying excitations are of crucial importance. One clearly sees a gap $\Delta$ and a quasiparticle structure in the response functions. However, if one probes the same superconductor by the high-energy photons $\hbar \omega \gg \Delta$, then the response of the superconductor is the same as of the normal metal - the high energy photons do not see quasiparticles and instead they are absorbed by the bare electrons. This is because the long-range phase coherent correlations in the superconductor become irrelevant in this case and the external probe sees a bare particle rather than a quasiparticle. Similar, in the QCD
case in order to create a hadron of a large mass one has to probe the QCD vacuum by the high energy (frequency) external probe (current) and hence the physics (masses) of the highly excited hadrons should be insensitive to the condensation of the chiral pairs in the vacuum.

II. A SIMPLE PEDAGOGICAL EXAMPLE

It is instructive to consider a very simple quantum mechanical example of symmetry restoration high in the spectrum. Though there are conceptual differences between the field theory with spontaneous symmetry breaking and the one-particle quantum mechanics (where only explicit symmetry breaking is possible), nevertheless this simple example illustrates how this general phenomenon comes about.

The example we consider is a two-dimensional harmonic oscillator. We choose the harmonic oscillator only for simplicity; the property that will be discussed below is quite general and can be seen in other systems. The Hamiltonian of the system is invariant under $U(2) = SU(2) \times U(1)$ transformations. This symmetry has profound consequences on the spectrum of the system. The energy levels of this system are trivially found and are given by

$$E_{N,m} = (N + 1); \ m = N, N - 2, N - 4, \ldots, -(N - 2), -N,$$

where $N$ is the principle quantum number and $m$ is the (two-dimensional) angular momentum. As a consequence of the symmetry, the levels are $(N + 1)$-fold degenerate.

Now suppose we add to the Hamiltonian a $SU(2)$ symmetry breaking interaction (but which is still $U(1)$ invariant) of the form

$$V_{SB} = A \theta(r - R),$$

where $A$ and $R$ are parameters and $\theta$ is the step function. Clearly, $V_{SB}$ is not invariant under the $SU(2)$ transformation. Thus the $SU(2)$ symmetry is explicitly broken by this additional interaction, that acts only within a circle of radius $R$. As a result one would expect that the eigenenergies will not reflect the degeneracy structure of seen in eq. (1) if the coefficients $R, A$ are sufficiently large. Indeed, we have solved numerically for the eigenstates for the case of $A = 4$ and $R = 1$ in dimensionless units and one does not see a multiplet structure in the low-lying spectrum as can be seen in Fig. 1.

What is interesting for the present context is the high-lying spectrum. In Fig. 1 we have also plotted the energies between 70 and 74 for a few of the lower $m$’s. A multiplet structure is quite evident—to very good approximation the states of different $m$’s form degenerate multiplets and, although we have not shown this in the figure these multiplets extend in $m$ up to $m = N$.

How does this happen? The symmetry breaking interaction plays a dominant role in the spectroscopy for small energies. Indeed, at small excitation energies the system is mostly located at distances where the symmetry breaking interaction acts and where it is dominant. Hence the low-lying spectrum to a very large extent is motivated by the symmetry breaking interaction. However, at high excitation energies the system mostly lives at large distances, where physics is dictated by the unperturbed harmonic oscillator. Hence at higher energies the spectroscopy reveals the $SU(2)$ symmetry of the two-dimensional harmonic oscillator.
III. IMPLICATIONS FOR BARYON SPECTRA

If chiral symmetry restoration of the second kind happens in the regime where the spectrum is still quasidescrete (i.e. the successive excitations with the given spin are well separated), then the symmetry restoration imposes very strong constraints on the spectrum, which will be discussed below. A question then arises as to which extent and where the hadron spectrum is still quasidescrete? Clearly it is the case for the low-lying resonances. However, the resonances with the given spin are still well separated in the mass region $M \sim 2$ GeV and higher, which can be seen from the nucleon spectrum, see Fig. 2, as well as from the meson spectrum [5]. In addition, if the linear-like behaviour of both angular ($M^2 \sim J$) and radial ($M^2 \sim n$, where $n$ numerates radial excitations of the states with the given spin) Regge trajectories is valid up to large masses, then the spectrum should be still quasidescrete in the region of validity of Regge phenomenology.

What are the implications of the chiral symmetry restoration for a quasidiscrete spectrum? The equality of the spectral functions $\rho_1(s)$ and $\rho_2(s)$ means that both masses $m_1$ and $m_2$ as well as the amplitudes $\langle 0| J_1 | n_1 \rangle$ and $\langle 0| J_2 | n_2 \rangle$ coincide for all the successive radial excitations. In other words, the excited states must fill in the irreducible representations of the parity-chiral group [2,3]. Thus the task is to find all possible representations of the $SU(2)_L \times SU(2)_R$ group that are compatible with the definite parity of the physical state. We emphasize parity because the irreducible representations of the chiral group $(I_L, I_R)$ (with $I_L$ and $I_R$ being isospins of the left and right subgroups) are not generally eigenstates of the parity operator, because under parity operation the left quarks transform into the right ones and vice versa. However, a direct sum of two irreducible representations $(I_L, I_R) \oplus (I_R, I_L)$ does contain eigenstates of the parity operator and hence the corresponding multiplet should be considered as a set of basis states for physical hadrons. What crucially important is that such a multiplet necessarily includes baryons with opposite parity.

Since the total isospin $I$ of a baryon can be obtained from the left and right isospins $I_L$ and $I_R$ according to a standard angular momentum addition rules and since there are no baryons with isospin greater than $3/2$, one immediately obtains the following allowed...
parity-chiral multiplets \((1/2, 0) \oplus (0, 1/2), (3/2, 0) \oplus (0, 3/2), (1/2, 1) \oplus (1, 1/2)\). The first multiplet corresponds to parity doublets of any spin in the nucleon spectrum. The second one describes the parity doublets of any spin in the delta spectrum. However, the latter multiplet combines one parity doublet in the nucleon spectrum with the parity doublet in the delta spectrum with the same spin.

Summarizing, the phenomenological consequence of the chiral symmetry restoration of the second kind high in \(N\) and \(\Delta\) spectra is that the baryon states will fill out the irreducible representations of the parity-chiral group. If \((1/2, 0) \oplus (0, 1/2)\) and \((3/2, 0) \oplus (0, 3/2)\) multiplets were realized in nature, then the spectra of highly excited nucleons and deltas would consist of parity doublets. However, the parity doublet with given spin in the nucleon spectrum \(a-priori\) would not be degenerate with the doublet with the same spin in the delta spectrum; these doublets would belong to different representations, \(i.e.\) to distinct multiplets and their energies are not related. On the other hand, if \((1/2, 1) \oplus (1, 1/2)\) were realized, then the high-lying states in \(N\) and \(\Delta\) spectrum would have a \(N\) parity doublet and a \(\Delta\) parity doublet with the same spin and which are degenerate in mass. In either of these cases the high-lying spectrum must systematically consist of parity doublets. We stress that this classification is the most general one and does not rely on any model assumption about the structure of baryons. What is interesting is that the same classification can be trivially obtained if one assumes that the chiral properties of baryons in the chirally restored regime are determined by three \textit{massless} valence quarks.

Now we have to analyze an empirical spectrum of baryons in order to see whether there are signs of chiral symmetry restoration. What is immediately evident from the empirical low-lying spectrum is that positive and negative parity states with the same spin are not nearly degenerate. Even more, there is no one-to-one mapping of positive and negative parity states of the same spin with masses below 1.7 GeV. This means that one cannot describe the low-lying spectrum as consisting of sets of chiral partners. It is not so surprising since the low-lying spectrum is mostly driven by the chiral symmetry breaking effects. The absence of parity doublets low in the spectrum is one of the most direct pieces of evidence that chiral symmetry in QCD is spontaneously broken (and very strongly).
FIG. 2. Excitation spectrum of the nucleon. The real part of the pole position is shown. Boxes represent experimental uncertainties. Those resonances which are not yet established are marked by two or one stars according to the PDG classification. The one-star resonances with $J = 1/2$ around 2 GeV are given according to the recent Bonn (SAPHIR) results.

Starting at the mass $M \sim 1.7$ GeV one observes three almost perfect parity doublets with spins $J = 1/2, 3/2, 5/2$. It is important that all these states are well established and have **** or *** status according to PDG classification. The next excitations with the same quantum numbers are not yet established though existing crude data support parity doubling phenomenon. The lowest excitations with spin 9/2 are also well established states and they represent another good example of parity doubling. There are well established states with $J = 7/2$ and $J = 11/2$ where the chiral partners have not yet been identified. According to a chiral symmetry restoration scenario they must exist. So it is a very interesting and important experimental task to find such states as well as to clear up a situation for $J = 1/2, 3/2, 5/2$ resonances around 2 GeV.

In the delta spectrum there are also obvious parity doublets starting from the mass region 1.9 - 2 GeV. Again, there are well established states (like $J^P = 7/2^+$ at 1950 MeV) where the chiral partner has not yet been identified. So similar to the nucleon spectrum, the upper part of the delta spectrum must be experimentally explored.

IV. IMPLICATIONS FOR MESON SPECTRA

Recent data on highly excited mesons also suggest an evidence for chiral symmetry restoration in hadron spectra [4]. Consider, as an example, the pseudoscalar and scalar mesons $\pi, f_0, a_0, \eta$ within the two-flavor QCD. The corresponding currents (interpolating fields) $J_\pi(x), J_{f_0}(x), J_{a_0}(x)$ and $J_\eta(x)$ belong to the $(1/2, 1/2) \oplus (1/2, 1/2)$ irreducible representation of the $U(2)_L \times U(2)_R = SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A$ group. Specifically the pairs $(J_\pi(x), J_{f_0}(x))$ and $(J_{a_0}(x), J_\eta(x))$ form the basis of the $(1/2, 1/2)$ representation of the chiral group $SU(2)_L \times SU(2)_R$. If the vacuum were invariant with respect to $U(2)_L \times U(2)_R$ transformations, then all four mesons, $\pi, f_0, a_0$ and $\eta$ would be degenerate (as well as all their excited states). Once the $U(1)_A$ symmetry is broken explicitly through the axial anomaly, but the chiral $SU(2)_L \times SU(2)_R$ symmetry is still intact in the vacuum, then the spectrum would consist of degenerate $(\pi, f_0)$ and $(a_0, \eta)$ pairs. If in addition the chiral $SU(2)_L \times SU(2)_R$ symmetry is spontaneously broken in the vacuum, the degeneracy is also lifted in the pairs above and the pion becomes a (pseudo)Goldstone boson. Indeed, the masses of the lowest mesons are

$$m_\pi \simeq 140\,\text{MeV}, \quad m_{f_0} \simeq 400 - 1200\,\text{MeV}, \quad m_{a_0} \simeq 985\,\text{MeV}, \quad m_\eta \simeq 782\,\text{MeV}.$$ 

This immediately tells that both $SU(2)_L \times SU(2)_R$ and $U(1)_V \times U(1)_A$ are broken in the QCD vacuum to $SU(2)_{I}$ and $U(1)_V$, respectively.

$^1\bar{\eta}$ represents the two-flavor singlet state; its mass for the lowest mesons can be extracted from the physical $\eta$ and $\eta'$ states.
Systematic data on highly excited mesons are still missing in the PDG tables. We will use recent results of the partial wave analysis of mesonic resonances from 1.8 GeV to 2.4 GeV obtained in \( p\overline{p} \) annihilation at LEAR, see Table below. We note that the \( f_0 \) state at 2102±13 MeV is not considered by the authors as a \( q\overline{q} \) state (but rather as a candidate for glueball) because of its very unusual decay properties and very large mixing angle. This is in contrast to all other \( f_0 \) mesons in this region, for which the mixing angles are small. Therefore these mesons are regarded as predominantly \( u, d = n \) states. Hence, in the following we will exclude the \( f_0 \) state at 2102±13 from our analysis which applies only to \( n\overline{n} \) states.

| Meson  | \( J^P \) | Mass (MeV) | Width (MeV) | Reference |
|--------|----------|------------|-------------|-----------|
| \( f_0 \) | 0 0\(^+\) | 1770 ± 12 | 220 ± 40 | [7] |
| \( f_0 \) | 0 0\(^+\) | 2040 ± 38 | 405 ± 40 | [5] |
| \( f_0 \) | 0 0\(^+\) | 2102 ± 13 | 211 ± 29 | | |
| \( f_0 \) | 0 0\(^+\) | 2337 ± 14 | 217 ± 33 | |
| \( \eta \) | 0 0\(^-\) | 2010\(^{+35}_{-60} \) | 270 ± 60 | | |
| \( \eta \) | 0 0\(^-\) | 2285 ± 20 | 325 ± 30 | | |
| \( \pi \) | 1 0\(^-\) | 1801 ± 13 | 210 ± 15 | | |
| \( \pi \) | 1 0\(^-\) | 2070 ± 35 | 310\(^{+100}_{-50} \) | | |
| \( \pi \) | 1 0\(^-\) | 2360 ± 25 | 300\(^{+100}_{-50} \) | | |
| \( a_0 \) | 1 0\(^+\) | 2025±? | 320±? | | |

The prominent feature of the data is an approximate degeneracy of the three highest states in the pion spectrum with the three highest states in the \( f_0 \) spectrum:

\[
\pi(1801 \pm 13) - f_0(1770 \pm 12),
\]

\[
\pi(2070 \pm 35) - f_0(2040 \pm 38),
\]

\[
\pi(2360 \pm 25) - f_0(2337 \pm 14).
\]

This can be considered as a manifestation of chiral symmetry restoration high in the spectra. The approximate degeneracy of these physical states indicates that the chiral \( SU(2)_L \times SU(2)_R \) transformation properties of the corresponding currents are not violated by the vacuum. This means that the chiral symmetry breaking of the vacuum becomes irrelevant for the high-lying states and the physical states above form approximately the chiral pairs in the \((1/2, 1/2)\) representation of the chiral group. The physics of the high-lying hadrons is such as if there were no spontaneous chiral symmetry breaking.

A similar behaviour is observed from a comparison of the \( a_0 \) and \( \eta \) masses high in the spectra:

\[
a_0(2025\pm?) - \eta(2010\^{+35}_{-60}).
\]

Upon examining the experimental data more carefully one notices not only a degeneracy in the chiral pairs, but also an approximate degeneracy in \( U(1)_A \) pairs \((\pi, a_0)\) and \((f_0, \eta)\) (in those cases where the states are established). If so, one can preliminary conclude that
not only the chiral $SU(2)_L \times SU(2)_R$ symmetry is restored, but the whole $U(2)_L \times U(2)_R$ symmetry of the QCD Lagrangian. Then the approximate $(1/2, 1/2) \oplus (1/2, 1/2)$ multiplets of this group are given by:

\[
\pi(1801 \pm 13) - f_0(1770 \pm 12) - a_0(?) - \eta(?); \quad (7)
\]

\[
\pi(2070 \pm 35) - f_0(2040 \pm 40) - a_0(2025\pm?) - \eta(2010^{\pm 35}_{-60}); \quad (8)
\]

\[
\pi(2360 \pm 25) - f_0(2337 \pm 14) - a_0(?) - \eta(2285 \pm 20). \quad (9)
\]

This preliminary conclusion would be strongly supported by a discovery of the missing $a_0$ meson in the mass region around 2.3 GeV as well as by the missing $a_0$ and $\eta$ mesons in the 1.8 GeV region. That these missing mesons should indeed exist is also supported by the hypothesis of the linear radial Regge trajectories for highly excited states \[1,6\]. We have to stress, that the $U(1)_A$ restoration high in the spectra does not mean that the axial anomaly of QCD vanishes, but rather that the specific gluodynamics (e.g. instantons) that are related to the anomaly become unimportant there. It should also be emphasized that the only restoration of $U(1)_V \times U(1)_A$ symmetry (without the $SU(2)_L \times SU(2)_R$) is impossible. This was discussed in ref. \[2\]. The reason is that even if the effects of the explicit $U(1)_A$ symmetry breaking via the axial anomaly vanish, the $U(1)_V \times U(1)_A$ would still be spontaneously broken once the $SU(2)_L \times SU(2)_R$ were spontaneously broken. This is because the same quark condensates in the QCD vacuum that break $SU(2)_L \times SU(2)_R$ do also break $U(1)_V \times U(1)_A$.

V. IMPLICATIONS FOR MODELING OF HADRONs

It is quite natural to assume that the physics of the highly excited hadrons is due to confinement in QCD. If so it follows that the confining gluodynamics is still important in the regime where chiral symmetry breaking in the vacuum has become irrelevant. Then it follows that the mechanisms of confinement and chiral symmetry breaking in QCD are not the same.

The phenomenon of chiral symmetry restoration high in the spectra rules out the potential description of high-lying hadrons in the spirit of the constituent quark model. Clearly, the chiral symmetry restoration by itself implies that constituent quarks as effective degree of freedom (whose mass is directly related to spontaneous chiral symmetry breaking in the vacuum) become inadequate high in the spectrum, though it is a fruitful concept for the low-lying hadrons. That the potential description is incompatible with the parity doubling is also seen from the following.

Consider, for instance, mesons. Within the potential description of mesons the parity of the state is unambiguously prescribed by the relative orbital angular momentum $L$ of quarks. For example, all the states on the radial pion Regge trajectory are $^1S_0$ $q\bar{q}$ states, while the members of the $f_0$ trajectory are the $^3P_0$ states. Clearly, such a picture cannot explain the systematical parity doubling as it would require that the stronger centrifugal
repulsion in the case of $^3P_0$ mesons (as compared to the $^1S_0$ ones) as well as the strong and attractive spin-spin force in the case of $^1S_0$ states (as compared to the weak spin-spin force in the $^3P_0$ channel) must systematically lead to an approximate degeneracy for all radial states. This is very improbable. Similar conclusions can be easily obtained for baryons [1]. More generally, the chiral symmetry restoration of the second kind is in contradiction with all models where chiral symmetry breaking is induced by confinement.

The potential picture also implies strong spin-orbit interactions between quarks while the spin-orbit splittings are absent or very small for excited mesons and baryons in the $u, d$ sector. The strong spin-orbit interactions inevitably follow from the Thomas precession (once the confinement is described through a scalar confining potential) [2], and this very strong spin-orbit force must be practically exactly compensated by other strong spin-orbit force from e.g. the one-gluon-exchange interaction in this picture. In principle such a cancellation could be provided by tuning the parameters for some specific (sub)families of mesons. However, in this case the spin-orbit forces become very strong for other (sub)families. This is a famous spin-orbit problem of constituent quark model.

This picture should be contrasted with the string description of highly excited hadrons [3]. Within the latter one these hadrons are the relativistic strings (with the color-electric field in the string) with practically massless bare quarks at the ends; these massless quarks are combined into parity-chiral multiplets. The string picture is compatible with the chiral symmetry restoration because there always exists a solution for the right-handed and left-handed quarks at the end of the string with exactly the same energy and total angular momentum. Since the nonperturbative field in the string is pure electric and the electric field is ”flavor-blind”, the string dynamics itself is not sensitive to the specific flavor of a light quark once the chiral limit is taken. This picture explains the empirical parity-doubling because for every intrinsic quantum state of the string there necessarily appears parity doubling of the states with the same total angular momentum of hadron. Hence the string picture is compatible not only with the $SU(2)_L \times SU(2)_R$ restoration, but more generally with the $U(2)_L \times U(2)_R$ one. In addition, there is no spin-orbit force at all once the chiral symmetry is restored. This is because the helicity operator does not commute with the spin-orbit operator and a motion of a quark with the fixed helicity is not affected by the spin-orbit force.

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Note also that a scalar potential explicitly breaks the chiral symmetry in contradiction to the requirement that the chiral symmetry must be restored high in the spectra.
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