Confinement as crossover

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The order parameter of confinement together with the haaron model of the QCD vacuum is reviewed and it is pointed out that the confining forces are generated by the non-renormalizable, invariant Haar-measure vertices of the path integral. A hybrid model is proposed for the description of the crossover leading to the confining vacuum. This scenario suggests that the differences between the low and the high temperature phases of QCD should be looked for in the quark channels instead of the hadronic sector.

I. INTRODUCTION

Having no systematic derivation of the confining forces in pure Yang-Mills theories the studies of the long range properties of the hadronic vacuum are usually based on model computations [1]-[4]. In order to get closer to the understanding of the problem it is obviously better to use models which contain at least partially the original gluonic degrees of freedom.

The bag model [1] is based on weakly interacting quarks and gluons, the confinement is realized by the difference of the energy density of two different vacua: inside and outside of the bag. No colored degrees of freedom are supposed to exist outside. The boundary of the bag, considered as a dynamical degree of freedom is obviously non-renormalizable. In the Abelian dual superconductor model [2] the Schwinger relation which asserts that the electric and the magnetic charges are inversely proportional to each other indicates that the magnetic condensate is due to a non-renormalizable coupling constant. The stochastic confining model [3] does not shed light on this question since it is based on the cluster expansion leaving the origin of the correlations an open question. The invariant Haar-measure terms of the functional integral which yield the string tension within the haaron model [4] are non-renormalizable, as well.

It is worthwhile noting that haaron model is the only one where the non-renormalizable term which is responsible for the string tension is already present in the original asymptotically free Yang-Mills Lagrangian. In fact, the gauge invariant, non-perturbative lattice regularization is based on the invariant Haar measure for the gauge group valued link variables. The logarithm of the Haar measure, treated as a local interaction potential in the action is non-polynomial and thereby non-renormalizable. Nevertheless the cut-off can be removed by suppressing these vertices sufficiently fast in the continuum limit [5].

We are confronted with an interesting possibility: How can it happen that the leading infrared force of the Yang-Mills theory comes from non-renormalizable vertices? The string tension, being a dimensionful parameter, can be generated by a relevant operator only. Since the non-renormalizable terms are irrelevant these vertices influence the interaction at the cut-off scale only and could have been left out from the theory according to the universality. The solution of this apparent paradox is rather simple [6]: Any theory with internal scale has at least two scaling regimes, an UV and an IR one, separated by a crossover at the internal scale. The non-renormalizable vertices are indeed irrelevant in the UV scaling regime but they might become relevant at the IR side of the crossover, in the IR scaling regime, where the IR forces are generated.

Section II overviews briefly the symmetry and the order parameter related to the confinement. An effective theory, the haaron model, to describe confinement as a destructive interference is mentioned in Section III. The lesson of the manner the confining force is generated in this model is discussed in Section IV. The finite temperature aspects of confinement as a crossover phenomenon are touched upon in Section V. Finally, Section VI is for the conclusion.

II. ORDER PARAMETER AND DESTRUCTIVE INTERFENCE

The order parameter for confinement is given in terms of the analytically continued massive quark propagator,

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\[ \Omega(x,t) = \text{tr}(\psi(x,t)\bar{\psi}(x,t+i\beta)), \]  

(1)

where the trace is over the color and the spin indices and \( \beta = 1/T \). Note that for infinitely heavy quarks in a time independent environment this expression reduces to the Polyakov line

\[ \omega(x) = \text{tr}Pe^{i\int_0^T d\tau A_0(x,\tau)} \]  

(2)

up to a constant multiplicative factor, where \( A_0(x,\tau) \) is the temporal component of the Euclidean gauge field.

\( \Omega \) displays the status of the symmetry with respect to the center of the gauge group which is the group \( \mathbb{Z}_n \) for the gauge group \( SU(n) \). The order parameter can be easily introduced even for Yang-Mills models in continuous Minkowski space-time \([7]\) and it remains a manifestly gauge invariant, well defined observable.

The dynamical quarks make the picture more complicated and we should distinguish two competing confinement mechanisms, a hard and a soft one. Both are driven by the increase of the effective coupling strength as the color charges are separated. The hard confining mechanism of the Yang-Mills models is responsible for the flux tube formation and the linearly rising potential between a static quark-anti quark pair. The soft mechanism is due to the Dirac-sea polarization and, similarly to the supercritical vacuum of QED \([8]\), shields the isolated quarks \([9]\). The soft mechanism cuts short the hard one and saturates the linearly rising potential when the flux tube between a static quark-anti quark pair is broken by the polarization of the Dirac-sea.

This ”deconfining ” vacuum-polarization effect appears in the dynamics of our order parameter, as well: The formal center symmetry is broken by the fermion determinant in the grand canonical ensemble. But it is easy to see that the Legendre transformation of the baryon number between the canonical and the grand canonical ensemble is ill defined in the thermodynamical limit due to the confinement mechanism. In fact, the free energy is infinitely large for states with non-vanishing triality \([10]\) which turns the free energy into a non-differentiable function of the baryon number density in the thermodynamical limit, and the control of the baryon number by a chemical potential into a highly non-trivial problem \([11]\). It is the canonical ensemble for the triality \([10], [11]\) which should rather be used in this case and this ensemble is formally center symmetrical. But this formal symmetry is broken spontaneously at low temperatures \([11]\). At high temperature the center symmetry is broken dynamically by the gluon kinetic energy. There is no reason to expect that the two unrelated symmetry breaking mechanisms would generate the same expectation value for \( \Omega \) thereby \([1]\) remains to be an order parameter which experiences a non-analytic dependence on the environmental variables at the deconfinement transition \([12]\).

The dynamical picture of confinement with \( SU(n) \) as color gauge group is the following \([7]\): The configuration space for global gauge rotations, \( SU(n)/\mathbb{Z}_n \), is multiply connected. Its fundamental group, the center \( \mathbb{Z}_n \), can be used to lump the time-dependent gluon field configurations into \( n \)-tuples in such a manner that the trajectories of an \( n \)-tuple correspond to the same initial or end points in the multiply connected space \( SU(n)/\mathbb{Z}_n \) but differ on the covering space, \( SU(n) \). The center symmetry of the pure gluonic system makes the action \( S \) of the trajectories of an \( n \)-tuple degenerate in the absence of quarks. Thus the contribution of \( n \) trajectories of an \( n \)-tuple to the transition amplitude is

\[ A_E = ne^{-S}, \quad A_M = ne^{iS} \]  

(3)

in Euclidean and Minkowski space-time. When a spectator quark is propagating along with the gluons then it picks up the \( \mathbb{Z}_n \) phase of the center transformation and the contribution of an \( n \)-tuple is vanishing due to the destructive interference between the homotopy classes,

\[ A_E = \sum_{\ell=1}^{n} e^{i\frac{2\pi}{n}\ell-S}, \quad A_M = \sum_{\ell=1}^{n} e^{i\frac{2\pi}{n}\ell+iS}. \]  

(4)

Notice that the non-positive definite phase factor comes from the projection operator which is supposed to install Gauss’ law and may lead to a destructive interference even for imaginary time. The semiclassical expansion, saturated by Wu-Yang monopoles in the Prasad-Sommerfeld limit \([12]\) supports the confinement as a destructive interference phenomenon, as well.

For an \( SU(2) \) gauge model the center symmetry expresses the invariance of a three vector under rotation by angle \( 2\pi \) and the destructive interference is due to the factor \(-1\) the spinors of the fundamental representation collect during a rotation by \( 2\pi \).

\footnote{The \( n \)-ality of a multi-quark state with \( N \) quarks and \( \bar{N} \) anti-quarks in an \( SU(n) \) gauge model is defined as \( t = N - \bar{N} (\text{mod} n) \).}
The high temperature deconfining transition is due to the too high kinetic energy barrier for gluons to follow the trajectories in the whole homotopy class. In the high temperature phase "there is no time" to realize all homotopy class, the destructive interference is prohibited and quarks can propagate. Note that the kinetic energy driven dynamical symmetry breaking occurs at high energy or in short time processes contrary to the potential energy governed spontaneous symmetry breaking which is observed at low energy or long time. When dynamical quarks are present then the spontaneous breakdown of the center symmetry selects a homotopy class which dominates the sum and leads to the screening of the isolated triality charge.

III. HAARON MODEL

We start this brief summary of the haaron model with a remark about the importance of keeping the exact gauge invariance in a computation to extract the string tension in Yang-Mills theory. Suppose that gauge invariance is implemented in an approximate manner and the state with a static quark-anti quark pair is $|q\bar{q}\rangle_0 + |q\bar{q}\rangle_1$ where $|q\bar{q}\rangle_0$ has the proper transformation rule under gauge transformations and $|q\bar{q}\rangle_1$ not. The non-covariant component $|q\bar{q}\rangle_1$ appears to contain uncontrollable charge distribution. The expression

$$
\langle q\bar{q}|_0 + \langle q\bar{q}|_1 \rangle \mathcal{H} (|q\bar{q}|_0 + |q\bar{q}|_1) = \langle q\bar{q}|_0 \mathcal{H} |q\bar{q}|_0 + \langle q\bar{q}|_1 \mathcal{H} |q\bar{q}|_1
$$

for the static potential shows that the charges in the component $|q\bar{q}|_1$ will break the flux tube for a sufficiently large separation of the test charges and saturate the potential. The gauge non-covariant components shield off the string tension when they are present with any small amplitude.

There is another indication of the strong relation between gauge invariance and the confining forces. The effective theory for the Polyakov line obtained in the strong coupling expansion shows that the minimum of the effective potential is at a vanishing value of the order parameter at low temperature due to the presence of the invariant Haar-measure for the gauge field. The replacement of the Haar-measure with a non-compact measure may keep the gauge invariance intact at any finite order of the loop expansion but certainly would break it at a non-perturbative level.

Guided by these remarks we wish to base our effective theory for the confining forces on the invariant Haar-measure in the path integral. Consider $SU(2)$ Yang-Mills theory for simplicity where the center transformation amounts to $\Omega \to -\Omega$. The lattice regularized path integral is given as

$$
\int D_H[aA_\mu(x)]e^{-S_{YM}[aA_\mu(x)]},
$$

where $a$ stands for the lattice spacing. The invariant integration measure for $A_\mu^j(x) = u(x)\omega^j(x), j = 1, 2, 3, (\omega^j(x))^2 = 1$ can be written as

$$
D_H[aA_\mu(x)] = D[\omega(x)]D[u(x)]e^{\sum_j \log \sin^2 au(x)}
$$

in terms of the flat integration measure $D[u(x)]$ for $u(x)$. The center transformation, $u(x) \to u(x) + \pi/a$, performed at a given equal-time hypersurface is a symmetry of the periodic potential $a^{-4} \log \sin^2 au$ appearing in (3). Let us integrate out the UV modes from the path integral and lower the cut-off. This step makes the dimensional transmutation explicit and induces an effective gauge theory model with dimensional parameters. First we choose a gauge in this theory where the temporal component $A_0(x)$ is diagonal, $\omega^a(x) = \delta^{a,3}$ and then we set the non-diagonal components of the gauge field to zero, leaving behind a compact $U(1)$ gauge model. The Feynman gauge is chosen for this model. As the final step, we set the spatial component of the Abelian gauge field to zero. What we find at the end is an effective theory for the diagonal, temporal component of the original gauge field, $u(x)$. The corresponding effective action will be approximated in the framework of the gradient expansion by the form

$$
S_{eff}[u] = \int d^4x \left[ \frac{1}{2} (\partial_\mu u(x))^2 - V(u(x)) \right].
$$

2This does not present serious problem in QED where the gauge non-covariant contributions are not gaining importance by non-perturbative effects in the absence of the confining forces.

3The factor $a^{-4}$ is to compensate the space-time integration volume $d^4x$ of the potential in the action. It makes the measure term vanishing in dimensional regularization.
The potential, the remnant of the invariant Haar-measure is periodic \( V(u+2\pi/\ell) = V(u) \). The periodicity reflects the discrete center symmetry of the original theory and allows the Fourier representation \( V(u) = \sum_n v_n \cos n \ell u \). The center symmetry requires that the symmetry \( u \rightarrow u + 2\pi/\ell \) is respected by the vacuum. The obvious consequence of this symmetry is the absence of any barrier in the effective potential between the periodic minima. This leads to the flattening of the effective potential for \( u \) and the masslessness of the field \( u(x) \).

It is well known that the sine-Gordon model is equivalent with a Coulomb gas. The simplest way to see this is to expand the generating functional in the coupling constant \( v \),

\[
Z[\rho] = \int D[u] e^{-\frac{i}{4} u \cdot G^{-1} u + i u \cdot \rho + \frac{1}{2} v_m} \int dx (e^{im\ell u(x)} + e^{-im\ell u(x)})
\]

\[
= \sum_n \frac{(v_m/2)^n}{n!} \prod_{j=1}^n \int dx_j \sum_{\sigma_j = \pm 1} \int D[u] e^{-\frac{i}{4} u \cdot G^{-1} u + i u \cdot (\rho + \sigma)}
\]

\[
= e^{-\frac{1}{2} \rho \cdot G \cdot \rho} \sum_n \frac{(v_m/2)^n}{n!} \prod_{j=1}^n \int dx_j \sum_{\sigma_j = \pm 1} e^{-\frac{1}{2} \sigma \cdot G \cdot \sigma} e^{-\rho \cdot G \cdot \sigma},
\]

where \( G \) is the massless propagator and \( \sigma(x) = m\ell \sum_{j=1}^n \sigma_j \delta(x - x_j) \). This is the grand canonical partition function of a four dimensional gas of particles interacting with the inverse of the massless propagator, the Coulomb potential. The first exponential in the last line represents the perturbative self-interaction of the external source \( \rho \), the second stands for the self-interaction of the particles and finally the third one describes the interaction between the source and the particles. Let us ignore the inter-particle forces and the partition function for non-interacting particles can be resummed,

\[
Z[\rho] \approx e^{\frac{1}{2} \rho \cdot G \cdot \rho + \frac{1}{2} v_m} \int dx \cos(i m \ell \int dy G(x-y)\rho(y)).
\]

These steps, repeated for each Fourier modes give

\[
Z[\rho] = \int D[u] e^{-\int d^4x \left[ (\partial_\mu u(x))^2 - V(u(x)) - i \rho(x) u(x) \right]} \approx e^{-\frac{1}{2} \rho \cdot G \cdot \rho} \int dx V(i \int dy G(x-y)\rho(y)).
\]

The particles representing the vertices of the perturbation expansion are called haarons since their contributions come from the invariant measure of the original path integral.

It is worth mentioning three applications of the partial resummation (111). The leading long range part of the static potential between a quark-anti quark pair separated by \( x \) is

\[
- \int d^3y V \left( \frac{i}{|y|} - \frac{i}{|y-x|} \right) \approx -2V''(0)|x|,
\]

giving the string tension

\[
\sigma = -2V''(0).
\]

Notice that in the center symmetry broken phase there is no protection against mass generation and the massive propagator does not give linearly rising potential. Thus the measure term which gives vanishing contribution in the UV regime and is included to assure the full gauge invariance only actually generates the leading long range force. Since it generates a new dimensional parameter, the string tension, the measure term must be relevant in the IR regime.

The second application of the resummation gives a confining version of the NJL model. We start with the Lagrangian

\[
L = \frac{1}{2} (\partial_\mu u)^2 - V(u) + i \bar{\psi} \gamma^\mu \gamma^5 \psi - igj,
\]

where \( j = u \bar{\psi} \gamma^\mu \gamma^5 \psi \) and perform the free haaron gas resummation yielding

\[
L_{\text{eff}} = i \bar{\psi}(x) \partial_\mu \gamma^\mu \psi(x) - V \left( ig \int d^4y G(x-y)j(y) \right)
\]

\[
\approx i \bar{\psi}(x) \partial_\mu \gamma^\mu \psi(x) - \frac{1}{2} g^2 V''(0) \int d^4y j(x)G_2(x-y)j(y)
\]

\[
= i \bar{\psi}(x) \partial_\mu \gamma^\mu \psi(x) + g \sqrt{-V''(0)} j(x) \phi(x) + \frac{1}{2} \phi(x) \Box^2 \phi(x) ,
\]

where \( j = u \bar{\psi} \gamma^\mu \gamma^5 \psi \).
where

\[ G_2(x - y) = \int d^4 z G(x - z)G(z - y) = \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^4} e^{-i p(x - y)}, \]

(16)

and the auxiliary field \( \phi(x) \) was introduced in order to render the Lagrangian local. Notice that the \( 1/p^4 \) propagator of the auxiliary field, coupled to the quarks in the same manner as \( u(x) \), confines the color charges with a linearly rising potential and the string tension is \( [\bar{u}] \).

The third application of the resummation is the computation of the quenched quark propagator. The grand canonical partition function, the last line of Eq. \( [1] \) when \( \rho \) is replaced by the quark current \( j(x) \) shows that the quarks are propagating in the imaginary long range field

\[ u_{y,n}(x) = \frac{\text{inf}}{4\pi^2 (x - y)^2} \]

(17)
of the haarons. After performing the Wick rotation into Minkowski space-time this external field becomes real. The destructive interference between the homotopy classes appears in this effective model as the destructive interference between the scattering processes of a quark off the gas of haarons. The long range haaron field makes the phase shift diverging and the fast rotating phase of the scattered state cancels the quark propagator when the averaging over the haaron distributions is performed. In order to understand the propagation of a meson qualitatively let us assume that the haaron field at \( x \) and \( y \) is identical or completely uncorrelated when \( |x - y| < \xi \) or \( |x - y| > \xi \), respectively where \( \xi \approx \Lambda_{\text{QCD}}^{-1} \) is the correlation length of the haaron gas. As long as the quark and the anti-quark of the mesons propagate within the distance \( \xi \) the phase shift suffered by them is canceled and the haarons do not influence much the propagation. When the color charges are separate from each other more than \( \xi \) then the statistically independent phase shift suppresses the amplitude. The result is that the world lines can not separate more than the distance \( \xi \), the confinement radius.

IV. CROSSOVER IN THE VACUUM

Let us consider the thought-experiment when the hadronic matter is viewed by a microscope of adjustable space-resolution. When details below the distance scale \( \Lambda_{\text{QCD}}^{-1} \) are considered we find partons, i.e. quarks and gluons. As the resolution becomes worse and details on the scale well above \( \Lambda_{\text{QCD}}^{-1} \) are seen only then hadrons and glueballs are found. The interactions between quarks and gluons on the one hand, and between hadrons on the other hand, are very different. This difference can simply be recorded by following the scale dependence generated by them.

There are at least two different scaling regimes in any non-scale invariant theory, an UV and an IR one separated by a crossover at the internal scale of the theory, \( \Lambda_{\text{QCD}}^{-1} \) in our case, c.f. Fig. \( [1] \). The UV scaling reflects asymptotically free forces between quarks and gluons in the UV regime and short ranged Yukawa interactions among the asymptotic states, hadrons, on the IR side. In pure Yang-Mills theory glueballs are the asymptotic states in the IR and color charges remain strongly bound by the linearly rising potential. What was surprising in the haaron model picture is that the measure vertices of the action which are non-renormalizable, i.e. irrelevant in UV scaling regime can generate the leading long range force. There must be a change in the behavior of the measure vertices as we move towards the IR directions which explains their increased importance in the confining forces. The most natural scenario is that these operators, being irrelevant in the UV scaling regime become relevant in the IR side of the crossover.

This scenario raises a more general question, the possibility that non-renormalizable operators might play an important role in low energy physics. It is easy to see that this surprising phenomenon does not take place in models with mass gap \( m \neq 0 \). These models display a correlation length \( \xi \approx 1/m \) and the evolution of the running coupling constant slows down at distance \( x \gg \xi \). In fact, the evolution of the coupling constants is driven by the contribution of the modes around the running cut-off and the fluctuations at the scale \( x \gg \xi \) are suppressed by \( \exp(-x/\xi) \). The absence of runaway trajectories of the renormalization group flow indicates that all non-Gaussian operators are irrelevant in the IR scaling regime\(^4\).

Theories without mass gap may develop new relevant operators by the help of collinear or simple IR divergences which may drive the run-away trajectories. The \( \phi^4 \) model in the mixed phase possesses a non-renormalizable operator which is relevant at low energies \([2,3]\). The condensation mechanism in general can easily generate radically new

\(^4\) An irrelevant coupling constant may naturally be important if its fixed point value is not small.
their quark content. The deconfined quarks are rendered colorless and only their flavor quantum numbers reveal when dynamical quarks are present. A sort of soft confinement mechanism is operating in the high temperature phase potential between triality charges by vacuum-polarization, a mechanism similar to the polarization of the Dirac-sea.

The absence of screening mechanism leads to confining forces. The real RHIC experiment is different, we are interested in the long distance correlations and quasiparticle structure at high temperature assuming that thermal equilibrium is an acceptable approximation. How does the temperature modify the scaling laws and the renormalized trajectory? It is obvious that the renormalized trajectory is in good approximation temperature independent at the observational length scale $x \ll 1/T$ and the temperature induced effects show up around $x \approx 1/T$, as shown qualitatively in Fig. 1. For $T < T_{dec}$, the clusterization of the color charges can be best understood as the impossibility of screening the $1/3$ color charge of a quark by multi-gluon states whose color charge is sum of integers, $\sum \pm 1 \neq 1/3$. The absence of screening mechanism leads to confining forces. The deconfining phase transition can be characterized in the Hamiltonian description by the improper implementation of the Gauss’ law projection operator which does not exclude certain states with infinitely many gluons. These states carry the color charge of a quark or anti-quark \[1\]. The result is the possibility of screening a quark color charge by a gluon cloud whose wave functional is multi-valued, the rearrangement of the infinite sum $\sum \pm 1$ in such an order that it converges to $1/3$. The effect of the temperature at the deconfining transition is the removal of the linear potential between triality charges by vacuum-polarization, a mechanism similar to the polarization of the Dirac-sea when dynamical quarks are present. A sort of soft confinement mechanism is operating in the high temperature phase of the pure glue system. The deconfined quarks are rendered colorless and only their flavor quantum numbers reveal their quark content.

V. CROSSOVER AT HIGH TEMPERATURE

The fields $\chi$ and $\Psi$ correspond to baryon and meson states and $S_H[\bar{\chi}, \chi, \bar{\Psi}, \Psi]$ is the action for a hadronic field theory. Since we are interested in the low energy phenomena we can fix the original cut-off at a sufficiently high but finite energy scale $\Lambda_0$. The coupling constants $G_\chi$ and $G_\Psi$ govern the strength of interactions between the hadronic and the colored states and their initial value is $G_\chi(\Lambda_0) = G_\Psi(\Lambda_0) = 0$, together with the hadronic coupling constants in $S_H$. This scheme is not a double counting since it is casted in the path integral formalism, it is a possible parameterization of the effective action.

Such a hybrid model should hold the key to the understanding of the confinement phenomenon because it offers a singularity-free description of the crossover. As the cutoff is lowered the non-renormalizable coupling strengths remain small and the asymptotically free coupling $g$ grows. When the crossover is reached then $g$ explodes in perturbative QCD and an IR Landau-pole arises because the long range correlations of the ground state are supposed to be generated by the asymptotically free vertices. But such a hybrid model offers the following alternative. In the presence of operators which are important in the IR scaling regime there is a chance that $g$ stays finite because the desired long range correlations can be established first in the colored and after that in the neutral sector by the renormalization of the measure term as in the haaron model and the hadronic coupling constants, respectively.
The renormalization group flow of the pure glue system should have two different manifolds of IR fixed points, for \( T > T_{\text{dec}} \) and \( T < T_{\text{dec}} \), as shown in Fig. 1. The high temperature fixed points should be qualitatively similar to those of full QCD at low temperature, the role of quarks are being played by gluon states with multi-valued wave functionals.

For full QCD the difference between the low and the high temperature fixed point manifolds is more subtle since the soft confinement mechanism is operating in both phases, by means of quark or gluon states with multi-valued wave functionals. I believe that the quasiparticles of the high temperature fixed point are similar to those of the low temperature phase except that a new quark “flavor” appears in the form of gluonic states with multi-valued wave functionals. The difference between the real and this fake quark can be detected by electro-weak currents only, the color charges being equivalent. The main features of this scenario are the screening of the color charge of the deconfined quark and the presence of the usual hadronic bound states. As of the former, a careful numerical study should be carried out measuring the gluonic color charge polarization around a deconfined quark. There are two indirect numerical evidences supporting the latter conjecture, the presence of the usual hadronic bound states. The first is the observation that the spacelike string tension is to a large extent temperature independent and remains non-vanishing even at high temperature \([17]\), indicating that the equal time long-range correlations of the multi-quark states do not change at the deconfinement phase transition. Another result, indicating the unimportance of the string tension from the point of view of the structure of the hadronic states at \( T = 0 \) is that the hadronic structure functions are qualitatively reproduced after cooling, a modification of the gluonic configurations which removes the string tension.

**VI. CONCLUSION**

It was argued in the framework of the haaron model that the leading long range forces between a quark-anti quark pair are generated by non-renormalizable vertices. This phenomenon motivates a look into the confinement problem following the strategy of the renormalization group and suggests that the confinement characterizes the IR scaling regime.

A lesson learned in dealing with the renormalization group is that all important operators should be present in the initial Hamiltonian even with vanishing coupling strength in order to understand the appearance of the dynamically generated, new kind of forces. Such a point of view motivates a hybrid model which contains both the quark-gluon and the hadronic fields. Their difference is set by the initial condition for the renormalization group flow only: a finite value for the asymptotically free coupling constant and a vanishing strength for the hadrons. The non-renormalizable measure term is supposed to generate a crossover in this model where the hadronic coupling constants turn on and induce the interactions of nuclear physics.

Such a description of the long range structure, together with the screening mechanism of the quark color charges by gluons available at high temperature suggests that the main difference between the high and the low temperature phases is not in the hadronic but rather in the quark sector of QCD. Possible examples are the following: The chromomagnetic monopoles, being "hedgehog" configurations relate color and spin. The gluon polarization cloud around a deconfined quark with even or odd number of monopoles possesses integer or half-integer spin, respectively. In this manner the violation of the superselection rule for the charge induces a similar violation for the spin and the
The deconfined quark state is actually the sum of components with Bose and Fermi statistics \[7\]. Another triality-related effect is the deviation of the temperature of the quark and gluonic degrees of freedom at the deconfinement phase transition point, \(T_q = T_{gl}/3\).\[8\]