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THE NON-PERTURBATIVE QUARK-GLUON PLASMA *

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ABSTRACT

After a brief and subjective evaluation of some talks of this School dealing with the high temperature phase of QCD a mechanism for quark confinement is presented.

I. BRIEF SUMMARY

Several talks of this Workshop have addressed the properties of the high temperature deconfined phase of QCD. I shall first briefly summarize the salient results announced in these talks.

Most of our knowledge of the deconfined phase comes from the numerical studies of lattice gauge theory. Though some reliable results have already been obtained for SU(3) lattice Yang-Mills system, the numerical study of QCD with dynamical quarks is in its infancy. This is due to the lack of sufficiently efficient simulation method to handle non-local integrands in the path integral formalism. Kreutz's lectures[1] describe the hybrid Monte Carlo algorithm which generates statistically independent gluon field configurations for massive quarks on conventional size lattice with the least amount of computational need. The problem I think comes up persistently in the attempts to fabricate a general purpose fermionic algorithm is the loss of ergodicity when the fermions are light and their zero modes become relevant. Lacking of the understanding of this large dimensional integral it is not clear how to handle the surfaces in the configuration space where the fermion determinant is zero. In fact, it may take unreasonably long time for the usual stochastic algorithms to penetrate such surfaces since the integrand vanishes there. The situation is made even worse when multi-fermion correlation functions are needed because these observables may be dominated by the configurations in the vicinity of the zero modes, the region which is avoided by the stochastic sampling based on the fermion determinant. The dynamical symmetry breaking is indicated in the numerical simulations by identifying a "non-symmetrical" region in the configuration space where the "simulation time" spent by the algorithm diverges with the physical volume of the system. We may lose this ability to recognize dynamical symmetry breaking with massless fermions if the "simulation time" necessary for the algorithm to cross the zero mode walls of the non-local integrand increases with the volume. The question of fermionic algorithm is one of the most challenging problem in this area since the intuitive understanding of the dynamics may lead to a breakthrough in the quantitative investigations.

Engels' talk[2] is a short summary of some of the bulk properties of the gluon plasma. There is clearly a feeling of accomplishment in announcing that the finite size effects are under control in the computations. Together with the rather satisfying status of the continuum limit we have a fairly complete numerical description of the SU(2) gluon plasma. This state of affairs should set the standards for the studies of the complete QCD. Another issue covered in this talk is the temperature dependence of the internal energy ϵ, pressure p and the trace of the energy momentum tensor $\epsilon - 3p$. It is worth to bear in mind that these intensive quantities have no immediate relation with the structure of the quasiparticles. These quantities may appear similar to their free counterpart even when the quasiparticles are rather complicated and non-perturbative in their origin. A simple example for this situation may be the ϕ^4 self-interacting scalar field theory in 3+1 dimensions. This model is supposed to have trivial renormalized scattering matrix and describe non-interacting particles. This expectation is supported by a vast amount of numerical works which show that the effective coupling strength responsible for the interactions of the dressed particles is several order of magnitude smaller than the bare coupling constant when the cutoff is high enough. But the simplicity of the model reveals itself in the language of the dressed particles rather than the bare ones. What the thermodynamical potentials of the quark-gluon plasma teach us is that there are light, weakly interacting quasiparticles with the same degeneracy factors than the "free" quarks and gluons[3].

An interesting and rather provocative suggestion was put forward in Zinovjev's talk[4] who conjectured that the formation of the deconfined phase is characterized by the dynamical breakdown of the charge conjugation invariance. The strong coupling effective model for the Polyakov line derived in his lecture suggests that $A_0(\vec{x})$ develops thermal expectation value in the deconfined phase. The dynamical symmetry breaking occurs in this case since A_0 changes sign under charge conjugation. Though the role of A_0 as a symmetry breaking field in the high temperature phase seems to be in conformity with the numerical results[5] one needs caution at this point. The reason is that this symmetry breaking should systematically be investigated in the framework of the effective three dimensional theory for the static modes which model is strongly interacting for any choice of the effective coupling constants and the symmetry breaking expectation values might be washed away by the quantum fluctuations[6]. Further works are clearly needed to clarify this interesting question.

Alvarez-Estrada pointed out an important simplification in the infrared properties of the quark-gluon plasma[7]: The dominant contribution to the $A_0(\vec{x})$ correlation function in the infrared are coming from the Debye screening rather than the higher order effective vertices of the effective three dimensional model. The results like this are very important to select the subset of Feynman graphs which are relevant in a given physical problem.

All of these results refer to the static, bulk properties of the plasma at zero baryon density. The implementation of chemical potential for the baryon density or real time processes are pure numerical problems whose solution would immediately widen our knowledge about strong interaction enormously. These issues will not be pursued here and we have to be contended by some remark about the local quasiparticle structure of the high temperature phase presented in Section V below.

Another common feature of most of these results is that they address some non-perturbative aspects of the high temperature phase. One wonders why non-perturbative phenomena can be so important at high temperature where the effective coupling constant $g^2(T)$ is small. It is the subject of Section II to show by inspecting the finite

temperature Feynman graphs that perturbation expansion does not apply for the infrared modes at arbitrary high temperature. Section III gives a more detailed account of these non-perturbative modes in terms of an effective three dimensional model. We shall see that the gluon correlation functions are not "less non-perturbative" in the deconfined phase as in the vacuum. A simple mechanism for quark confinement based on the numerical experiences is described in Section IV. The application of this mechanism predicts some unusual features for the deconfined quarks in the high temperature phase which is the subject of the Section V. The last two sections may be considered as another demonstration that the high temperature phase of QCD is different and in fact much richer than a collection of weakly interacting "deconfined" quarks and gluons.

II. THE NON-PERTURBATIVE INFRARED MODES

The question we investigate in this section is the true expansion parameter of the perturbation expansion in QCD at finite temperature[8] $T = \frac{1}{\beta}$. Consider the one-loop contribution V to the gluon four point function which is depicted in Fig. 1. It is proportional to

$$V \sim g^4 \frac{1}{\beta} \sum_n \int \frac{d^3p}{(2\pi)^3} \frac{1}{(\vec{p}^2 + \omega_n^2 + M^2)^2} \tag{1}$$

where M is some mass gap to be determined later and the Matsubara frequency is given by $\omega_n = 2\pi n T$. One would naively think by comparing (1) with the tree level result which is $O(g^2)$ that the true expansion coefficient is $g^2 T$ at finite temperature. Though this is incorrect it is alarming to see loop integrals with higher and higher dimension as multiplicative factors in the perturbation expansion which must arise to preserve the dimension of the whole expression. This phenomenon characteristic to super-renormalizable models makes the higher order of the perturbation expansion infrared divergent and forces us to resum infinite number of diagrams. The mass gap M is generated by such a resummation method, it turns out to be $M_{\rm Debye} = O(gT)$ in the electric propagator and the magnetic sector remains massless when the conventional resummation methods are applied. We shall have some hand-waving argument in Section III that the mass gap is $M_{\rm mag} = O(g^2T)$ for the magnetic propagator. The infrared contribution in (1) comes form the region $p < M$ with $n = 0$ and is proportional to $g^2 \frac{g^2 T}{M} = g^2 \lambda$ where $\lambda = \frac{g^2 T}{M_{\rm mag}}$. It is easy to see that other higher-loop diagrams follow this pattern and the true expansion parameter is in fact λ which is not calculable in the perturbation expansion. The contributions of higher order Feynman graphs collapse into the same order in g and are separated according to λ. The increase of the temperature helps to send $g^2(T)$ to zero and to decrease the effects of the interactions since there is usually an overall factor $(g^2(T))^k$ in front of these contributions with k being independent of the order of the graph. But the actual shape of these contributions and the (asymptotic) convergence of the expansion remain unclear whatever high temperature is chosen. The high temperature plasma appears as the collection of weakly interacting plane wave quarks and gluons at short distances only where asymptotic freedom protects the perturbation expansion.

Fig. 1: One-loop contribution to the gluon four point function.

III. THE EFFECTIVE MODEL

A more systematic method to isolate the non-perturbative effects in the high temperature phase is to eliminate all perturbative modes and try somehow to solve the effective model for the remaining degrees of freedom. Though this procedure is not completed yet one hopes to be able to solve the non-perturbative effective model because it contains less degrees of freedom and is much simpler than the complete QCD. We shall consider Yang-Mills theories without matter fields for simplicity.

The 3+1 dimensional Yang-Mills theory can be thought at finite temperature as a three dimensional Yang-Mills-Higgs system containing infinitely many fields. To see this use static temporal gauge $\partial_0 A_0(x) = 0$ and decompose the space components of the gluon field into Fourier-modes according to the Euclidean time dependence

$$\vec{A}^n(\vec{x}) = \frac{1}{\beta}\int_0^\beta d\tau e^{-i\omega_n\tau}\vec{A}(\vec{x},\tau) \tag{2}$$

The time integral can easily be carried out in the 3+1 dimensional Euclidean action written in terms of the fields $\vec{A}^n(\vec{x}), A_0(\vec{x})$ and generates a three dimensional model *with many* fields. The use of this construction stems from the fact that the free three dimensional propagator for $\vec{A}^n$ has the mass ω_n. It seems then reasonable to eliminate the heavy modes and look into the effective model for the light, static fields. The loop-integrals arising from the elimination of the non-static modes are infrared finite and the effective model for the static modes can in principle be computed in perturbation expansion. The tree level effective model is obtained by setting the non-static modes to zero leading to the three dimensional Yang-Mills-Higgs model

$$S_{\text{eff}}^{\text{tree}} = \frac{1}{g^2T}\int d^3x\{-\frac{1}{4}(F_{ij}^a)^2 + \frac{1}{2}(\vec{D}^a A_0)^2\} \tag{3}$$

with $\vec{A}^0$ and A_0 as the Yang-Mills and the Higgs field, respectively. The corrections to $S_{\text{eff}}^{\text{tree}}$ due to the massive non-static modes can be handled in the framework of the decoupling theorem[9] which suggests to separate the renormalizable and the non-renormalizable corrections $\Delta S_{\text{eff}}^{\text{ren}}, \Delta S_{\text{eff}}^{\text{non-ren}}$, respectively. The renormalizable contributions generate only a local potential for the Higgs field $\Delta S_{\text{eff}}^{\text{ren}} = \int d^3x V(A_0)$. It turns out[10] that the one-loop approximation to the renormalizable part of the effective model $S_{\text{eff}}^{\text{ren}} = S_{\text{eff}}^{\text{tree}} + \Delta S_{\text{eff}}^{\text{ren}}$ reproduces the static sector of the complete 3+1 dimensional Yang-Mills theory with accuracy $O(g^2(T))$.

An important consequence of the formalism outlined above results from the fact that the three dimensional Yang-Mills-Higgs system $S_{\text{eff}}^{\text{ren}}$ is highly non-perturbative and confining either in the symmetric[11] or the Higgs[6] phase and the Wilson loops of the

effective model follow the area law at arbitrary high temperature. The space-like Wilson loops of the 3+1 dimensional Yang-Mills theory differ only perturbatively from the corresponding Wilson loops computed in the effective model. Thus the space-like Wilson loops follow area behavior even in the high temperature phase. The permanent confinement in the effective model should not be mixed with the deconfining features of the high temperature 3+1 dimensional Yang-Mills system. The former is described by the space-like Wilson loops, the latter refer to the behavior of the space and time-like Wilson loops only. What is relevant in this discussion is that the equal-time thermal correlation functions for the gluon field experience strong "confinement-like" non-perturbative effects. Since asymptotic freedom protects the perturbation expansion at short distances these non-perturbative effects show up only in the infrared. The length scale ρ where perturbation expansion becomes unreliable was found to be the same $\rho \sim 0.23fm$ at $T \sim 2T_{\text{dec}}$ as in the vacuum[12]. The deconfined phase appears as "non-perturbative" as the vacuum beyond the confinement radius. Observe that the gauge coupling constant of the three dimensional model g^2T is dimensional suggesting that the mass scale generated for $\vec{A}^0$ in this model is $M_{\text{mag}} = O(g^2T)$.

The coupling constants are temperature dependent in $V(A_0)$ so the effective model may have a phase transition at some temperature. The numerical results[5] are compatible with the conjecture that the effective model is in the symmetric and the Higgs phase for $T < T_{\text{dec}}$ and $T > T_{\text{dec}}$, respectively. The formation of the Higgs phase at high temperature might be understood by the close examination of the differences of the renormalization in three and four dimensions[13]. The semiclassical analysis of the Higgs phase of the three dimensional Yang-Mills-Higgs system[6] uses the magnetic monopole configurations as saddle points. This suggests that chromomagnetic monopoles may be long-living resonances of the deconfined phase. The chromomagnetic monopole density which can be computed numerically without using the semiclassical approximation was found to be finite in the continuum limit of the lattice QCD[3] indicating that these monopoles are indeed present in QCD. This circumstance has some relevance in the description of confinement as dual Meißner effect[14] and the prediction of some unusual quantum numbers in the deconfined phase[15].

IV. PERMANENT CONFINEMENT OF TRIALITY

The notion of triality has been introduced in the quark model to give a simple characterization of the observed multi-quark states. It is defined as $k = N - \bar{N}(\text{mod}3)$ for the state $|N,\bar{N}>$ consisting of N quarks and $\bar{N}$ antiquarks and only states with zero triality can be observed. Note that gluons give no contribution to the triality. A kinematical construction supported by the manner the confinement-deconfinement phase transition shows up in the numerical studies will be presented below to eliminate states with nonzero triality[15]. I believe that the understanding of the dynamical origin of this mechanism will ultimately exclude gluons from the asymptotic states as well. But there are some lessons to learn from this mechanism even without the proper understanding its dynamical source.

The first step is to give a more formal definition for triality. *The center of the group* G is the subgroup C formed by the elements commuting with all elements of G, $[C,G] = 0$. The center of the group $SU(3)$ is Z_3 which consists of 3×3 matrices which are the product of the 3×3 identity matrix and a cubic root of one $z_{ij}^{k} = \delta_{ij} e^{ik\frac{2\pi}{3}}, k = 0, 1, 2$. The gluon and the quark fields transform as

$$\vec{A}^{a}(x)\lambda^{a} \rightarrow \vec{A}^{ga}(\vec{x})\lambda^{a} = g(x)(\vec{\partial} + \vec{A}^{a}(x)\lambda^{a})g^{\dagger}(x) \tag{4}$$

and

$$q(x) \rightarrow q^{g}(x) = g(x)q(x) \tag{5}$$

respectively under gauge transformations. It is easy to see that the triality is the charge with respect of the center of the global gauge group,

$$z^{k}|N, \bar{N}> = e^{ik\frac{2\pi}{3}(N-\bar{N})}|N, \bar{N}> \tag{6}$$

The special feature of the center of the gauge group is that while it leaves the gluon field invariant $\vec{A}^{z}(\vec{x}) = z\vec{A}(\vec{x})z^{\dagger} = \vec{A}(\vec{x})$ and appears as an "invisible" discrete symmetry group of the gluon dynamics the multi-quark states with non-zero triality do feel it.

The conventional way to describe confinement is to try to generate linear potential between the confined constituents. I shall use a less detailed kinematical characterization of the confined states: they should have vanishing propagator since it seems impossible to prepare an isolated states made of a particle with vanishing propagator. This condition can naturally be used only after eliminating the local gauge invariance. The linearly rising quark potential is expected when the quark-antiquark vacuum polarizations are neglected. The quark propagator should vanish as long as states with zero triality appear as intermediate states. This latter condition seems more appropriate to impose on our approximation since it can be expressed in terms of a conserved quantum number.

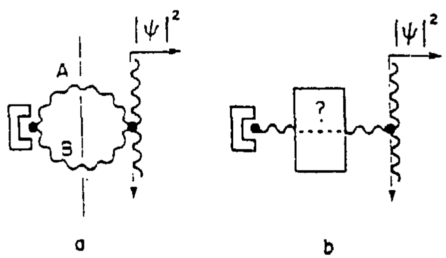

Fig. 2: (a) The two slit experiment illustrating the vanishing of the electron propagator at certain points. (b) The explanation of the absence of the electron becomes difficult if the slits are hidden.

The mechanism responsible for the vanishing of the propagator in the usual two slit thought experiment is illustrated in Fig. 2. Suppose that the detector is placed at the point where the length of the paths through the slits differ by half of the Compton wavelength. The amplitude finding the electron at that point is given by $\psi = c_1 e^{-iE_1 t} + c_2 e^{-iE_2 t}$ and is zero since $c_1 = -c_2, E_1 = E_2$. If the slits are hidden from us then the amplitude is expected of the form $\psi = ce^{-iEt}$ with no obvious reason to have $c = 0$.

The only way the absence of the electron can be understood is to have $E \rightarrow \infty$ since the Planck constant has an infinitesimal imaginary part for the causal propagators. This simple thought experiment shows that the absence of the particles can be understood in some circumstances as the completely destructive interference in their propagator rather than by assigning infinite energy to some states.

Another example of such a cancellation in the propagator which is closer to what happens in gauge theories is the vanishing overlap between states with integer and half integer spin in three dimensions. Consider a quantum rotator in a state with spin $\frac{1}{2}$

$$|\frac{1}{2}> = \int dg \mathcal{D}^{(\frac{1}{2})}(g)|\mathcal{R}(g)\vec{x}_0> \tag{7}$$

The invariant integration is over the group $SU(2)$, $\mathcal{D}^{(\frac{1}{2})}$ denotes a linear combination of the matrix elements of the spin $\frac{1}{2}$ $\mathcal{D}$-matrix and $\mathcal{R}(g)$ is the three dimensional rotation corresponding to the $SU(2)$ element g, $\mathcal{R}(g)\vec{x} = \vec{x}^{g}$

$$\vec{x}^{g}\vec{\sigma} = g\vec{x}\vec{\sigma}g^{\dagger} \tag{8}$$

where $\vec{\sigma}$ are the Pauli matrices. The overlap

$$< \frac{1}{2}|e^{-itH}|\vec{y}_0> = \int dg \mathcal{D}^{(\frac{1}{2})}(g) < \mathcal{R}(g)\vec{x}_0|e^{-itH}|\vec{y}_0> \tag{9}$$

vanishes for arbitrary t because the state $|\vec{y}_0>$ has integer angular momentum components only. The vanishing of (9) can be understood by pairing the contributions corresponding to g and $-g$. The sum of these two contributions is zero since $\mathcal{D}^{(\frac{1}{2})}(-g) = -\mathcal{D}^{(\frac{1}{2})}(g)$ and $\mathcal{R}(-g) = \mathcal{R}(g)$. This cancellation is demonstrated in Fig. 3 where the amplitude of propagating from the left arrow to the right one is the sum of two contributions which differ that the arrow has an additional rotation by 2π around some axis along the lower path. The two contributions differ in sign only when the state of the left and the right arrow has half-integer and integer spin, respectively. This cancellation is made possible because the two paths can not be transformed into each other smoothly[15], they are homotopically different.

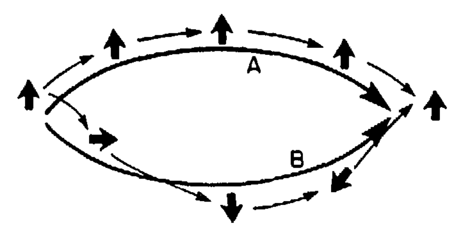

Fig. 3: The vanishing of the overlap between half-integer and integer spin states of the quantum rotator.

The previous argument seems suspicious since the wave function of the state $|\frac{1}{2}>$ has two different values at the same point. To make things better defined it is advised to separate those two, physically equivalent positions where the wave function has different values. This amounts to distinguishing the rotation by 2π from the identity. Fig. 3 suggests that the physical amplitude is the sum of the amplitudes corresponding to mathematically different but physically equivalent points. The notions introduced here are well known in mathematics[16]. A space Q is called multiply connected if a point moving in it can have trajectories which can not be transformed into each other, *i.e.* are homotopically different. The space Q' obtained from Q by splitting the points of Q according to the homotopy classes of the paths reaching it is called the covering space. The points on the covering space which belong to the same physical location q are called the image point of q. Functions defined on multiply connected spaces may assign different values to the same points as long as only local operators are applied on them and the values satisfy some compatibility conditions. The well known multi-valued functions are $y = \sqrt{x}$ or $y = \log(x)$ with the additional rule that local operators *e.g.* derivation act on the same Riemann sheets only. The different values of the wave function may differ in a constant phase only in order to have well defined observables. Note that the multi-valued functions become single-valued on the covering space.

It is not difficult to imagine the cancellation in the quark propagator by comparing (8) with (4) when g is constant. The propagation of the quark is thought in the color space as in the previous case and three different contributions to the quark propagator can be grouped together in a manner that they cancel each other. The difference between these contributions is that at some time *between* the creation and the annihilation of the quark the system undergoes a global gauge transformation by a different center element. The gluon dynamics is insensitive to such transformations but the quark propagator picks up the appropriate phase which sums up to zero. It is worth while noting that glueballs may carry triality charge despite the invariance of the gluon field under global gauge transformations from the center. Such states can be constructed analogously to (7) where the wave function changes under the application of the center of the group $SU(2)$ though the coordinate remains invariant. The sum of *infinitely* many integer may be fractional.

The cancellation mechanism above is purely kinematical. The dynamical part of this picture is whether the summation over homotopically different paths should be included in the path integral or not. The remarkable fact is that gauge theories are consistent in either case! To understand how this happens we have to go back to the proof of the path integral expression for a simple quantum system

$$\psi(Q,T) = \int \mathcal{D}[q(t)] e^{-iS[q]} \tag{10}$$

where the integration is over paths with end points $q(0) = 0, q(T) = Q$. The variation of the paths at the end point and the Taylor expansion of the integrand gives quickly that $\psi(Q,T)$ satisfies the appropriate Schrödinger equation[17]. The point is that this proof can be repeated when the integral is restricted to homotopically equivalent paths. One may multiply the contribution of different homotopy classes by different coefficients since the Schrödinger equation is linear. This ambiguity is fixed by requiring that the

observables should remain the same when the mathematically different but physically equivalent points are exchanged on the covering space. This amounts to requiring that observables should be invariant under rotation by 2π in the case of the rotator. There is no compelling reason to require that global gauge transformations from the center should leave the observables invariant. The choice of keeping only one homotopy class in the path integral corresponds to the dynamical breakdown of the center symmetry. Thus quark confinement is equivalent with the presence of the center symmetry in the vacuum.

One expects that the center symmetry breaks dynamically in short time processes. The rotator travel through more distance along the lower path in Fig. 3 which corresponds to the extra rotation by 2π. Its kinetic energy will suppress this path in the path integral for short time. The "long" paths may be missing completely for time $t < t_{dec}$ when the system has infinite degrees of freedom and the "rotation" is a collective coordinate like global gauge transformations. The deconfined phase is where the quarks propagate in one homotopy sector only. Observe that the center symmetry is broken at high temperature or high energy rather than at long time or in the vacuum.

The covering space for gluons consists of the field configurations $\vec{A}^{\ell}(\vec{x}), \ell = 0, 1, 2,$ where the index ℓ distinguishes the image points and the "invisible" global gauge transformations from the center change its value $z\vec{A}^{\ell}z^{\dagger} = \vec{A}^{\ell+1(\mathrm{mod}3)}$ The gluon states with triality k have the wave functional $\Psi[\vec{A}^{\ell}]$ on the covering space satisfying $\Psi[\vec{A}^{\ell+1(\mathrm{mod}3)}] = e^{ik\frac{2\pi}{3}}\Psi[\vec{A}^{\ell}]$. The suppression of some homotopy classes observed in the deconfined phase can equivalently be described by allowing certain linear combination of the triality classes, *e.g.* one can verify easily that the sum over triality classes with equal weight suppresses the homotopically nontrivial "long" paths in the path integral. Consider now a deconfined quark with non-vanishing propagator when the initial and the final gluon fields are $\vec{A}^{(i)}$ and $\vec{A}^{(f)}$, respectively. What is the triality of the gluon state corresponding to this deconfined quark ? The amplitude of propagating from the image point $\vec{A}^{(i)\ell}$ to $\vec{A}^{(f)\ell'}$ is $\mathcal{A}e^{i(\ell-\ell')\frac{2\pi}{3}}$ where the phase factor is generated by the quark propagating with the gluons. The only way the sum over the image points can be non-zero is that the oscillating phase of the quark propagator $e^{i\Theta}$ becomes cancelled by the phase factor coming from the wave functional $e^{i\alpha}$ (*c.f.* Fig. 4). Thus the gluon state has triality which screens the triality charge of the deconfined quark. Triality charges are either confined (confined phase) or completely screened (deconfined phase), no state with triality can be observed in QCD[15]. The deconfined phase is where a new kind of screening is available for the quarks, the localized excitations of the high temperature phase continue to have zero triality charge.

V. THE DECONFINED-SCREENED QUARKS

We have seen in the previous section that the deconfined-screened quark is a composite particle. This can be illustrated by an imaginary world where only the charmed quarks participate in the electro-weak currents. The charmed mesons appear as deconfined-screened quarks for the experimentalists of this world The role of the non-charmed quarks is played in our world by the gluon states which carry similar color charge than a quark but these states are not available in our familiar environment at $T < T_{dec}$. The separation of the screening gluon cloud from the quark in the high

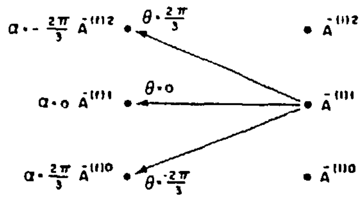

Fig. 4: Contributions to the quark propagator on the covering space.

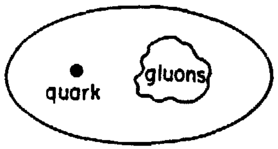

Fig. 5: The deconfined-screened quark.

temperature phase would generate a string of gluonic "mesons" in complete analogy with the situation in the vacuum in the presence of dynamical quarks.

The analogy with the imaginary world is even more complete, the screening gluon cloud may appear as a fermion. To see this we return to the chromomagnetic monopoles.Their characteristic property is that they are gauge hedgehogs, three-space rotations can be compensated by global gauge transformations. Consider $SU(2)$ gauge theory for simplicity where the wave functional of these states changes sign under global gauge rotation by 2π (center element!) in color space. The hedgehog property guarantees that the rotation by 2π in three-space gives the same change of sign. Thus these particles have half-integer spin and negative exchange parity according to the spin-statistics theorem for kinks[18]. This mechanism which is identical to the way how skyrmions can be quantized as fermions applies in the case of $SU(3)$ as well.

Perform a phase shift analysis with a deconfined-screened quark to determine its spin. We separate the gluon state corresponding to a deconfined-screened quark into two components, one which contains chromomagnetic monopoles and the other which does not. The former component of the quark-gluon composite particle behaves as a boson in this experiment. The states contributing to the partition function in the high temperature phase have no well defined triality. Since the chromomagnetic monopoles relate the triality charge to the spin-statistics the deconfined-screened quarks have bose *and* fermi components.

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