An Essay on Color Confinement

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Abstract

Color confinement is a consequence of an unbroken non-Abelian gauge symmetry and the resulting asymptotic freedom inherent in quantum chromodynamics. A qualitative sketch of its proof is presented.

1 Introduction

There has been an accumulation of evidence in favor of the quark model of hadrons [1] and we can no longer think of any other substitute for it. Yet, no isolated quarks have been observed to date, and we are inclined to think that observation of isolated quarks is, in principle, impossible. This is the hypothesis of quark confinement, and it has been further extended to that of color confinement that implies not only the unobservability of quarks but also of all the isolated colored particles such as quarks and gluons. Then a natural question is raised of whether or not we can account for this hypothesis within the framework of the conventional quantum chromodynamics (QCD) dealing with the gauge interactions of quarks and gluons. The answer to this question is affirmative and the detailed mathematical proof of color confinement has been published elsewhere [2-6]. In this article, therefore, we shall follow the flow of ideas underlying the proof in a qualitative manner.
The problem of color confinement may be decomposed into two steps. The first step consists of finding a consensus of interpretations of color confinement. Unless it is properly settled we do not know what we have to prove in the second step. Because of the importance of this subject many authors have proposed various interpretations. A typical example is Wilson’s area law for the loop correlation function in the lattice gauge theory [7]. When it is obeyed the interaction between a quark and an antiquark is given by a confining linear potential. Another example is given by coherent superposition of magnetic monopoles in the vacuum state [8-13]. This is dual to the superconducting vacuum based on coherent superposition of charged objects such as the Cooper pairs. Corresponding to the superconductor of the second kind a pair of magnetic monopoles can be connected by a quantized magnetic flux forming a hadronic string whose energy is proportional to the distance between them. Then the situation is similar to the preceding example.

In these examples one introduces a topological structure through monopoles, strings and instantons into the configuration space. In the present paper, however, we shall consider a different topological structure in the state vector space. For this purpose we look for a known example of confinement within the framework of known field theories, and we find a prototype example in quantum electrodynamics (QED) [2]. When the electromagnetic field is quantized in a covariant gauge, say, in the Fermi gauge, three kinds of photons emerge, namely, transverse, longitudinal and scalar photons, but only the transverse photons are subject to observation leaving the other two unobservable. We recognize that this is indeed a typical example of confinement, and we may be able to find some clues to color confinement by studying closely the mechanism of confinement of longitudinal and scalar photons in QED. For this reason we analyze its mechanism in Sec. 2 so that we can generalize it and apply it to QCD.

One of the profound features of gauge theories is the Becchi-Rouet-Stora (BRS) invariance [14] and its introduction is vital to the interpretation of confinement. Therefore, we shall describe some of the basic properties of this invariance in Sec. 3.

The strong interactions described by QCD possess a novel feature called asymptotic freedom [15,16], and in Sec. 4 we shall discuss how this aspect of strong interactions drew our attention and how the non-Abelian gauge theory entered the game. Finally in Sec. 5 we shall combine BRS invariance with asymptotic freedom to prove color confinement.

2 Quantum Electrodynamics and Indefinite Metric

When the electromagnetic field is quantized in a covariant gauge, say, in the Fermi gauge, we find transverse, longitudinal and scalar photons, but the latter two are never observed. We may interpret it as an example of confinement, and we have at least three alternative ways of explaining it. First, we can refer to the representations of the Poincaré group for massless particles [17,18]. Then, massless particles are known to have only two directions of polarization no matter what their spin is. Thus photons are always transversely polarized and the same would be true with gluons if they could be observed. The second method is to employ the Coulomb gauge by keeping only the transverse photons from the start. The remnants of unobservable photons manifest themselves in the form of the Coulomb potential. This method is applicable, however, only to the linear Abelian gauge theories such as QED. The third and the most useful method is the introduction of a subsidiary condition such as the Lorentz condition.
Quantization of the electromagnetic field in a covariant gauge forces us to introduce indefinite metric [19] which is inherited from the Minkowski metric. Thus the whole state vector space in QED can no longer possess the positive-definite metric, and for the physical interpretation of the theory we have to eliminate indefinite metric by imposing the Lorentz condition on the state vectors to select observable or physical states. In order to execute this program let us quantize the free electromagnetic field in the Fermi gauge and for a given momentum we have four directions of polarization, namely, two transverse, one longitudinal and one scalar. Thus we have four kinds of photons specified by the directions of polarization. The canonical quantization then implies that the scalar photons are represented by negative norm states. This is a consequence of the manifest covariance of the quantization of the vector field in the Minkowski space.

The emergence of indefinite metric indicates that observable states occupy only a portion of the whole state vector space called the physical subspace. In order to define such a subspace we introduce a subsidiary condition known as the Lorentz condition. Let us consider the four-divergence of the vector field, then it represents a free massless field even in the presence of the interactions. We decompose it into a sum of positive- and negative-frequency parts corresponding to destruction and creation operators, respectively. We find that the photons involved in this operator are special combinations of the longitudinal and scalar photons in the amplitude. We shall call them a-photons, then an a-photon state has zero norm. We can introduce an alternative combination of longitudinal and scalar photons called b-photons in such a way that a b-photon state also has zero norm. Thus for a given momentum we have two transverse (t-) photons, an a-photon and a b-photon. Although both an a-photon state and a b-photon state have zero norm, their inner product is non-vanishing so that they are metric partners.

A physical state is defined as such a state that is annihilated by applying the positive frequency part of the four-divergence of the vector field. This is the Lorentz condition. We can easily verify that the S matrix in QED transforms a physical state into another physical state since it commutes with the four-divergence. This is one of the general features of the subsidiary condition. Also we can easily verify that the b-photons are excluded from the physical subspace. Therefore, we have only t-photons and a-photons in the physical states. Then we can show that the inner product of a physical state involving at least one a-photon with another physical state vanishes identically. In other words, a-photons give no contributions to observable quantities, and both a- and b-photons escape detection. This is the confinement mechanism of the longitudinal and scalar photons. In QED only the transverse photons remain observable. In QCD, however, not only longitudinal and scalar gluons but also transverse gluons are unobservable. Thus, there are some essential differences in the nature of confinement between QED and QCD. In the former case confinement is kinematical in the sense that it could be understood without recourse to dynamics of the system, whereas in the latter case it is dynamical in nature as the proof depends sensitively on the dynamical properties of the system.

3 Quantum Chromodynamics and BRS Invariance

As we shall see in the next section strong interactions of quarks are mediated by a non-Abelian gauge field corresponding to the SU(3) color symmetry. Thus we shall discuss one of the most characteristic features of gauge theories known as the BRS invariance in this section [14].

Classical electrodynamics is gauge-invariant. Field strengths expressed in terms of the vector
field are invariant under the local or space-time-dependent gauge transformations of the latter. Given a source term, therefore, the solution of the equation for the vector field is not uniquely given, and this non-uniqueness is an obstacle to quantization. In order to overcome this difficulty we add to the gauge-invariant Lagrangian a term violating the local gauge invariance. This extra term is called the gauge-fixing term and was first introduced by Fermi. Later it has been generalized so as to include an arbitrary parameter called the gauge parameter. In the original form introduced by Fermi this parameter is equal to unity.

After quantization we find that we have to introduce indefinite metric into the state vector space and that the divergence of the vector field commutes with the S matrix. Because of the inclusion of the gauge-fixing term the field equation deviates from the classical Maxwell equation by a term proportional to the four-divergence of the vector field. It so happens that a matrix element of this four-divergence between two physical states vanishes identically because of the Lorentz condition, and the classical Maxwell equation is recovered in the physical subspace. In this way we find, despite the introduction of the gauge-fixing term, that expectation values of gauge-invariant quantities and the S matrix elements in the physical subspace are independent of the choice of the gauge parameter because of the congeniality between the gauge-fixing term and the subsidiary condition. In what follows we shall extend this approach to QCD.

There are many essential differences between QED and QCD, however. The former is an Abelian gauge theory described by a linear field equation, whereas the latter is a non-Abelian gauge theory described by a non-linear field equation. In both cases the gauge-invariant part of the Lagrangian is given by the square of the field strength. So, let us introduce the gauge-fixing term in QCD assuming the same structure as in QED. Then we recognize that it does not work because observable quantities depend explicitly on the gauge parameter. Another difficulty arises from the fact that the four-divergence of the gauge field is no longer a free field, and this prevents us from defining its positive frequency part. In other words, the Lorentz condition cannot be employed to define physical states in QCD. Thus we are obliged to find a device to overcome these difficulties and to this end we shall introduce the Faddeev-Popov ghost fields.

In order to eliminate the gauge-dependence of physically relevant quantities Faddeev and Popov have proposed a procedure of averaging the path integral over the manifold of gauge transformations. We skip the mathematical detail here and refer to the original paper [20], but we should mention that this procedure resulted in a new additional term in the Lagrangian called the Faddeev-Popov (FP) ghost term. This term involves a pair of Hermitian scalar fields, but they are anticommuting and consequently violate Pauli’s theorem on the connection between spin and statistics. For this reason they are called ghost fields. Pauli’s theorem is based on three postulates, (1) Lorentz invariance, (2) local commutativity or microscopic causality and (3) positive-definite metric for state vectors, and the FP ghost fields violate the last one obliging us to introduce indefinite metric into the theory.

Thus we face again the problem of eliminating indefinite metric from the theory with the help of an appropriate subsidiary condition to select physical states out of the whole state vector space. When physical states are so defined as those that are annihilated by applying a certain operator, that operator should commute with the S matrix as does the four-divergence of the vector field in QED. In order to find such an operator a novel symmetry discovered by Becchi, Rouet and Stora is extremely useful. Although this symmetry was originally utilized in renormalizing QCD, it plays an essential role in the proof of color confinement in QCD.
In a classical gauge theory a local gauge transformation is specified by a function of the space-time coordinates called the gauge function and the classical theory is invariant under such a transformation. This local gauge invariance is lost when the gauge-fixing and FP ghost terms are introduced. Besides, local gauge transformations are defined only for the color gauge field and the quark fields, but they are not even defined for FP ghost fields. The BRS transformations for the color gauge field and the quark fields are given by replacing the gauge function by one of the FP ghost fields in infinitesimal gauge transformations. Since we have a pair of ghost fields we introduce, correspondingly, a pair of BRS transformations. Then a question is raised of how to define BRS transformations of the ghost fields since their gauge transformations are not defined. Fortunately, this problem has a simple but beautiful solution. Their BRS transformations are introduced by demanding the invariance of the total Lagrangian under them.

The total Lagrangian including the gauge-fixing and FP ghost terms is no longer invariant under local gauge transformations, but it is invariant under the global BRS transformations. Noether’s theorem then tells us that there must be a pair of conserved quantities corresponding to a pair of BRS symmetries. They are Hermitian and called the BRS charges. As mentioned before there are two kinds of Hermitian FP ghost fields and correspondingly a BRS charge must involve one of the ghost fields. In what follows we keep only one of these two charges for simplicity. The BRS charge that we keep is anticommuting just as the FP ghost field, and consequently the square of the BRS charge vanishes and it is called nilpotent. The Hermiticity and nilpotency of the BRS charge would imply indefinite metric since otherwise it would be a null operator [21,22]. The nilpotency is important and allows us to introduce the concept of cohomology in the theory. After a long detour we are going to introduce an appropriate subsidiary condition. Physical states are defined as those states that are annihilated by applying the BRS charge [23].

The FP ghost fields do not appear in the conventional QED but we can also introduce them although they are non-interacting fields. Then we can combine the Lorentz condition with the additional condition implying the absence of FP ghosts to define the physical states. When these conditions are satisfied, we can prove that physical states so defined are annihilated by the BRS charge in QED.

The BRS charge is the generator of the BRS transformation and the BRS transform of an operator is given by the commutator or anticommutator of that operator with the BRS charge, and this transformation is also nilpotent. An operator which is the BRS transform of another operator is called an exact operator, then it is clear that the matrix element of an exact operator between a pair of physical states vanishes.

The equation for the non-Abelian gauge field deviates from the classical Maxwell equation and in fact the divergence of the field strength plus the color current does not vanish but is equal to a certain exact operator, which will be referred to as an exact current hereafter. Therefore, the classical Maxwell equation is recovered when we take the matrix element of the field equation between a pair of physical states. Furthermore, the BRS charge commutes with the S matrix. Thus the scenario in QED is reproduced almost exactly.

When single quark states and single gluon states are unphysical these particles are unobservable and consequently confined. Thus the problem of color confinement reduces to that of proving that they are unphysical states. We shall evaluate the expectation value of the exact current in a single quark state or a single gluon state. If they should belong to physical states the expectation values in these states would vanish identically, so that non-vanishing of the
expectation values would be a direct indication that these particles are unphysical and confined.

The four-divergence of the exact current vanishes, and we can give a set of Ward-Takahashi identities for Green’s functions involving the exact current [2-4]. By making use of the above set of Ward-Takahashi identities we can prove that the expectation value of the exact current in a single colored particle state survives when the exact current as applied to the vacuum state does not generate a massless spin zero particle. Therefore, the absence of such a massless particle is a sufficient condition for color confinement [2-4]. In order to check its absence we introduce the vacuum expectation value of the time-ordered product of the gauge field and the exact current and evaluate the residue \( C \) of the massless spin zero pole of the Fourier transform of this two-point function. The four-divergence of this two-point function is proportional to this constant \( C \) except for a trivial kinematical factor, and the divergence can be cast in the form of an equal-time commutator.

By checking this equal-time commutator closely we find that \( C \) is the sum of a constant \( a \) and the Goto-Imamura-Schwinger (GIS) term. The constant \( a \) is equal to the inverse of the renormalization constant of the color gauge field. These constants \( C \) and \( a \) satisfy distinct renormalization group (RG) equations and boundary conditions. We shall not enter this subject here since the mathematical detail has been given elsewhere [2-5], but we infer the fact that vanishing of \( a \) automatically leads to vanishing of \( C \) and color confinement is realized. Indeed, it has been known for some time that gluons are confined when \( a \) vanishes [24,25], but now with the help of the BRS invariance we could conclude that not only gluons but also all the colored particles are simultaneously confined. We shall come back to this subject again in Sec. 5.

### 4 Asymptotic Freedom

In this section we shall review briefly how and why our attention was drawn to the non-Abelian gauge theory in describing strong interactions. In particle physics strongly interacting particles such as nucleons and pions are called hadrons. Hadrons are composite particles of quarks and antiquarks, however, and we have to study the origin of the strong interactions of quarks.

We already know that strong interactions are mediated by the color gauge field and the quanta of this field are called the gluons since they glue up quarks together to form hadrons. Dynamics of quarks and gluons is called QCD as mentioned before. In the sixties experiments on the deep inelastic scattering of electrons on protons had been carried out. The differential cross-section had been measured by specifying the energy and direction of electrons without observing the hadrons in the final states. Then, apart from kinematical factors this differential cross-section can be expressed as a linear combination of two structure functions. They are functions of the square of the momentum transfer and the energy loss of the electron in the laboratory system. When these two variables increase indefinitely the two structure functions tend to be functions of the ratio of these two variables except for trivial kinematical factors. This characteristic behavior of structure functions is called the Bjorken scaling [26], and it is considered to be an empirical manifestation of the properties of strong interactions. What do we learn from this? In 1969 Feynman proposed the parton model and assumed that a nucleon consists of point-like partons moving almost freely inside the nucleon [27]. In order to keep the partons inside the nucleon, however, the four-momentum of a parton must be equal to a fraction \( x \) of the total four-momentum of the nucleon. The partons may be identified with the quarks.
and since $x$ is identified with the ratio of the two kinematical variables referred to in the above
the distribution of the fraction $x$ has been shown to be related to the structure functions.

From the success of the parton model in reproducing the Bjorken scaling we may infer
that quarks inside the hadrons are almost free and that the interactions of quarks turn out
to be weaker at shorter distances. This is a distinctive feature of strong interactions and we
may express it in the momentum space as follows: The probability of a process involving large
momentum transfer in strong interactions is small.

We look for a model satisfying this condition and find that only non-Abelian gauge interactions meet this requirement with the help of RG [15,16].

The concept of RG was first introduced by Stueckelberg and Petermann in 1953 [28], and it
was further advanced by Gell-Mann and Low in QED in 1954 [29]. Let us consider a dielectric
medium and put a positive test charge inside, then the medium is polarized, namely, negative
charges are attracted and positive ones are repelled by this test charge. As a consequence it
induces a new charge distribution in the medium. The total charge inside a sphere of radius
$r$ around the test charge is a function of $r$ and we call it the running charge. The vacuum
is an example of the dielectric media because of its ability of being polarized – the vacuum
polarization. In this case the test charge is called the bare charge and the total charge inside
a sphere of a sufficiently large radius is called the renormalized charge. The running charge
is a function of the radius $r$, but it can also be regarded as a function of momentum transfer
through the Fourier transformation. The bare charge then corresponds to the limiting value of
the running charge for infinite momentum transfer.

Gell-Mann and Low have proved on the basis of the RG method that given a finite renormalized charge the bare charge is equal to a certain finite constant independent of the value of the
renormalized one or it is divergent [29]. The Bjorken scaling phrased in terms of RG implies that
the bare coupling constant must be equal to zero. We shall refer to this property as asymptotic
freedom (AF), and the non-Abelian gauge theory is the only known example in which AF is
realized as clarified by Gross and Wilczek and by Politzer [15,16]. The origin of AF may be
traced back to the fact that the vector field introduces indefinite metric needed to realize AF
and that the non-Abelian gauge theory is the only example involving non-linear interactions of
the vector field.

Thus starting from the empirical Bjorken scaling we have finally reached the non-Abelian
gauge theory of strong interactions, namely, QCD.

5 Color Confinement

Now we are ready to present the proof of color confinement, at least verbally, by combining
arguments given in preceding sections.

In QED the square of the ratio of the renormalized charge to the bare one is equal to the
renormalization constant of the electromagnetic field. It is equal to the inverse of the dielectric
constant of the vacuum relative to the empty geometrical space. Usually the dielectric constant
of a dielectric medium is defined relative to the vacuum, but here we define it relative to the
empty geometrical space or the void.
This dielectric constant of the vacuum is larger than unity as a consequence of the positive-definite metric of the physical subspace, or more intuitively, it is a consequence of the screening effect due to the vacuum polarization. Then, let us consider a fictitious case in which the dielectric constant of the vacuum is smaller than unity. In this case we have antiscreening instead of screening when a test charge is placed in this fictitious vacuum, and such a vacuum is realized when a pair of virtual charged particles of indefinite metric should contribute to the vacuum polarization. In this case the running charge would be an increasing function of the radius $r$ at least for small values of $r$. Next we shall consider an extreme case of the vanishing dielectric constant, then a small test charge would attract an unlimited amount of like charges around it thereby bringing the system into a catastrophic state of infinite charge. Nature would take safety measures to prevent such a state from emerging, and a possible resolution is to bring another test particle of the opposite charge. The total charge of the whole system is equal to zero and charge confinement would be realized. In QED, however the dielectric constant of the vacuum or the inverse of the renormalization constant is larger than unity, and the above scenario reduces to a mere fiction.

The situation in QCD is completely different since it allows introduction of indefinite metric in the vacuum polarization and AF is one of its manifestations. In QCD what corresponds to the dielectric constant of the vacuum in QED is the inverse of the renormalization constant of the color gauge field denoted by $a$ in Sec. 3. If $a$ should vanish we would encounter a scenario similar to the one mentioned above and a test color charge would induce an intolerable catastrophic state. In Sec. 3 we have shown that such a state is excluded by means of the subsidiary condition that selects physical states. Therefore, what can be realized are states of zero color charge and this is precisely color confinement. Unlike electric charge, color charge is not a simple additive quantum number but a member of a Lie algebra $su(3)$, so that physically realizable states should belong to the one-dimensional representation of this algebra. Thus the entire problem of color confinement reduces to the proof that the constant $a$ vanishes.

Before presenting its proof we have to introduce the concept of the equivalence class of gauges [2,4,5]. When the difference between two Lagrangian densities is an exact operator we say that these two Lagrangian densities belong to the same equivalence class of gauges. For instance, two Lagrangian densities corresponding to two distinct values of the gauge parameter belong to the same equivalence class. In QCD hadrons are represented by BRS invariant composite operators [30-32], and the S matrix elements for hadron reactions are obtained by applying the reduction formula of Lehmann, Symanzik and Zimmermann [33] to Green’s functions defined as the vacuum expectation values of the time-ordered products of the BRS invariant composite operators. Then we can readily prove that the S matrix elements for hadron reactions are the same within the same equivalence class of gauges [2,4,5]. Color confinement signifies that the unitarity condition for the S matrix in the hadronic sector is saturated by hadronic intermediate states. That means that quarks and gluons have no place to show up in the unitarity condition just as longitudinal and scalar photons never appeared in the S matrix elements in QED. Therefore, we may take it for granted that the concept of color confinement is gauge-independent within the same equivalence class.

Then we come back to the evaluation of the constant $a$. First, it should be stressed that $a$ can be evaluated exactly as a function of the gauge coupling constant and the gauge parameter thanks to AF [2,5]. These two parameters define a two-dimensional parameter space, which is then decomposed into three domains according to the value of $a$, namely, zero, infinity and
finite. It should be stressed here that the existence of these three domains can be proved without recourse to perturbation theory. Of these three domains color confinement is manifestly realized in the first one, and also in the other two confinement should prevail because of the gauge-independence of the concept of color confinement. Evaluation of $a$ by means of RG based on $AF$ is a very interesting mathematical problem, but we shall refer to the original paper for the technical detail [5].

Finally, it should be stressed that confinement as has been discussed in this paper is realized only when we have an unbroken non-Abelian gauge symmetry [2]. When a certain gauge symmetry is spontaneously broken the exact current generates a massless spin zero particle as the Nambu-Goldstone boson and our proof of confinement breaks down. For instance, the electroweak interactions are formulated on the gauge group $SU(2) \times U(1)$, but spontaneous symmetry breaking reduces the gauge symmetry to the Abelian $U(1)$ corresponding to the electromagnetic gauge symmetry. Thus the electroweak interactions do not possess any unbroken non-Abelian gauge symmetry and are not capable of confining any particle.

To conclude, we have presented the flow of ideas towards intuitive understanding of the mechanism of color confinement without recourse to mathematical detail, but interested readers are encouraged to refer to the original articles.

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