Confinement in three dimensional magnetic monopole–dipole gas

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Abstract

Confinement of electrically charged test particles in the dilute plasma of monopoles and pointlike magnetic dipoles is studied. We calculate the tension of the string emerging between the infinitely separated test particles. The string tension is an increasing function of the dipole density provided other parameters of the plasma are fixed. The relevance of our results to confining gauge theories is discussed.

1 Introduction

We study the confining properties the plasma of the Abelian magnetic monopoles with a fraction of the magnetic dipoles in three dimensional Euclidean space–time. This kind of the plasma may appear in gauge theories such as the Georgi–Glashow model which possesses the topologically stable classical solution called the ’t Hooft–Polyakov monopole \[.\] The dipoles are realized as the monopole–anti-monopole bound states. Since the long range gauge fields of the monopole are associated with an unbroken Abelian subgroup the long range properties of the monopoles and as a consequence, of the dipoles, are Abelian.

The confining properties of the dilute Abelian monopole–antimonopole plasma are well known \[2\]. The test particles with opposite electric charges are confined due to formation of a stringlike object between the charge and the anticharge. The string has a finite thickness of the order of the inverse Debye mass and a finite energy per the string length ("string tension"). Thus the potential between the test particles is linear at large distances \[2\].

At a sufficiently high temperature the plasma of the Abelian monopoles undergoes the Berezinsky–Kosterlitz–Thouless phase transition \[3, 4, 5\]. The vacuum in the high
temperature phase is filled with the neutral monopole–anti-monopole bound states\footnote{\cite{4, 6, 7}} obeying non–zero magnetic dipole moments\footnote{Note that the monopole binding is qualitatively similar to the formation of the instanton molecules in the high temperature phase of QCD suggested to be responsible for the chiral phase transition\cite{8}.}. The long range fields of the magnetic dipoles are weak and they can not induce a non–zero string tension between the electric charges separated by large distances. Thus the finite temperature phase transition separates the confining and deconfining phases of the model.

Since the formation of the magnetic dipole states is supported by the increase of the temperature we expect to have a mixed plasma of the Abelian monopoles and dipoles at the confining side of the phase transition. In this paper we investigate the behavior of the electric charges in the mixed plasma. Below we show that the magnetic dipoles affect the string tension even despite of the inability of the pure magnetic dipole gas to confine electric charges. For simplicity the dipoles are considered to be pointlike.

Note that in the finite temperature Georgi–Glashow model the charged bosons are supposed\cite{12} to play an important role. In this paper we are interested in the pure monopole–dipole gas without any charged matter fields. The structure of the paper is as follows. The path integral formulation as well as simplest properties of the monopole and dipole gas are discussed in Section 2. The interaction of the electric charges are studied both numerically and analytically in Section 3. Our conclusions are summarized in the last Section.

2 Gas of monopoles and magnetic dipoles

The partition function of the monopole–dipole gas has the following form:

\[
Z = \sum_{M=0}^{\infty} \frac{\zeta_m^M}{M!} \left( \prod_{a=1}^{M} \left( \sum_{q_a} \int d^3x_a \right) \right) \sum_{N=0}^{\infty} \frac{\zeta_d^N}{N!} \left( \prod_{\alpha=1}^{N} \int d\mathbf{n}_\alpha \int_0^\infty dr_\alpha \int d^3x_\alpha \right) \exp \left\{ -\frac{1}{2} \int d^3x \int d^3y \left[ \rho(x) + (\tilde{\rho}^{(\mu)}(x), \tilde{\partial}_x) \right] \left[ \rho(y) + (\tilde{\rho}^{(\mu)}(y), \tilde{\partial}_y) \right] D_{reg}(x - y) \right\},
\]

where \(\zeta_m (\zeta_d)\) is the fugacity of the monopole (dipole) component of the gas, and

\[
\rho(x) = g_m \sum_a q_a \delta^{(3)}(x - x_a), \quad \tilde{\rho}^{(\mu)}(x) = \sum_\alpha \tilde{\mu}_\alpha \delta^{(3)}(x - x_\alpha),
\]

are the densities of the magnetic charges and magnetic dipole moments, respectively. We use the latin, \(a\) (greek, \(\alpha\)) subscripts to denote the monopole (dipole) parameters. The magnetic charge of the \(a\)\textsuperscript{th} individual monopole in units of the fundamental monopole charge, \(g_m = 4\pi/g\), is referred to as \(q_a\). In the dilute monopole gas the (anti)monopoles have unit magnetic (anti)charges, \(|q_a| = 1\).

The magnetic moment of the \(\alpha\)\textsuperscript{th} dipole is \(\tilde{\mu}_\alpha = g_m \tilde{n}_\alpha r_\alpha\) (no sum is implemented), where \(r_\alpha\) is the "dipole size" and \(\tilde{n}_\alpha\) is the direction of the dipole magnetic moment.
For simplicity we consider a dilute gas of the point-like dipoles characterized only by a (fluctuating in general case) dipole moment and the space–time position. At the end of the paper we discuss the applicability of the results obtained in the pointlike dipole approximation to the real physical systems.

In eq.(1) the integration over the dipole moments \( \int d^3 \mu_\alpha \) is represented as the integration over direction \( \vec{n}_\alpha \) of the dipole moment and over the dipole size \( r_\alpha \), weighted with a normalized distribution function \( F(r) \), \( \int_0^\infty F(r) \, dr = 1 \). The function \( D(x) = 4\pi|x|^{-1} \) in eq.(1) is the inverse Laplacian, and the subscript “reg” indicates that the monopole and dipole self–interaction terms are subtracted.

The second line of eq.(1) can be rewritten as follows

\[
\int D\varphi \exp\left\{-\frac{g^2}{32\pi^2} \int d^3x \, (\vec{\partial}\varphi)^2 + i \sum_{a} q_a \varphi(x_a) + i \sum_{\alpha} (\vec{r}_\alpha, \vec{\partial}) \varphi(x_\alpha) \right\}. \tag{3}
\]

Substituting this equation into eq.(1), performing summation over the monopole charges \( q_a \) and integrating over the dipole moment directions \( \vec{n}_\alpha \), and the monopole (dipole), \( x_{a(\alpha)} \) positions, we get the following partition function for the gas of the monopoles and dipoles:

\[
\mathcal{Z} = \int D\varphi \exp\left\{-\int d^3x \, \mathcal{L}(x) \right\}, \tag{4}
\]

\[
\mathcal{L} = \frac{g^2}{32\pi^2} (\vec{\partial}\varphi)^2 - 2\zeta_m \cos \varphi - 4\pi\zeta_d \int_0^\infty dr \, F(r) \frac{\sin(r |\vec{\partial}\varphi|)}{r |\vec{\partial}\varphi|}, \tag{5}
\]

where \( |\vec{a}| = \sqrt{a^2} \).

We study the magnetic monopole–dipole plasma in the weak coupling regime. The density of the monopoles and antimonopoles, \( \rho_m \), and the density of the dipoles, \( \rho_d \), can be calculated as the following expectation values taken in the statistical sum (1):

\[
\rho_m = < M > \quad \text{and} \quad \rho_d = < N > .
\]

In the weak coupling limit the leading order contributions to the (anti)monopole and dipole densities are proportional to the corresponding fugacities [2, 9]:

\[
\rho_m = 2\zeta_m, \quad \rho_d = 4\pi\zeta_d. \tag{6}
\]

The correlations of the fields \( \varphi \) in model (4) are short ranged indicating that the photon gets a non-zero mass in the presence of the monopole plasma [2]. Indeed the expansion of lagrangian (4) in small fluctuations of the field \( \varphi \) gives:

\[
\mathcal{L} = \frac{g^2}{32\pi^2} \left( 1 + \frac{4\pi}{3} \zeta_d \mu^2 \right) (\vec{\partial}\varphi)^2 + \zeta_m \varphi^2 + O(\varphi^4), \tag{7}
\]

where

\[
\overline{\mu^2} = g_m^2 r^2 = \frac{16\pi^2}{g^2} \int_0^{\infty} dr \, F(r) \, r^2, \tag{8}
\]

\[\text{Here and below we omit inessential constant factors in front of the path integrals.}\]
is the mean quadratic magnetic moment of a dipole. Thus the two–point correlation function \( \langle \varphi(0) \varphi(x) \rangle \) is exponentially suppressed at large distances as \( e^{-M_D |x|} \), where \( M_D \) is the Debye mass in the monopole–dipole plasma:

\[
m_D^2 = \epsilon^{-1} M_D^2, \quad M_D^2 = \frac{32\pi^2 \zeta_m}{g^2}.
\]

(9)

Here \( M_D \) is the Debye mass in the pure magnetic monopole plasma and \( \epsilon \) is the dielectric constant (permittivity) of the magnetic plasma,

\[
\epsilon = 1 + \frac{4\pi}{3} \zeta_d \mu^2.
\]

(10)

The presence of the dipole states in the plasma leads to charge renormalisation, \( g^2 \to \epsilon g^2 \) according to eqs. (7,9). In the next Section we show that the dipoles affect also the confining properties of the monopole plasma.

3 Electric charges in magnetic monopole–dipole gas

Let us consider two static infinitely heavy electrically charged test particles in the pure magnetic monopole plasma. The monopoles and anti-monopoles form a double layer at the sides of the minimal surface, spanned on the particles’ trajectories \([2]\): at one side of the surface the magnetic density is positive while at another side the density is negative. The stringlike structure between the charges has a finite thickness of the order of the inverse Debye mass, \( M_D^{-1} \), and a non-zero energy density per unit area of the surface. Thus, the electric charges are confined in the magnetic monopole plasma. At large separations \( R \) between the test particles the confining potential is linear \([2]\) while at \( RM_D \ll 1 \) the potential becomes of the form \( \text{const.} \cdot R^\alpha \), where \( \alpha \approx 0.6, \text{Ref. [10]} \).

The behavior of the interparticle potential in the pure plasma of pointlike dipoles is different. The potential is linear at small interparticle separations, \( V(R) = \text{const.} \cdot \rho_{\text{d}} \bar{r} \cdot R \), due to interaction of the dipole clouds gathered near the electric sources \([3]\). At large distances \( R \) the potential is of the Coulomb type with the renormalized coupling \( g, g^2 \to \epsilon g^2 \), where \( \epsilon \) is the dielectric constant given in eq.(10).

Thus, the pure gas of pointlike dipoles gives rise to the nontrivial modification of the interparticle potential only at small separations \( R \) since the fields of an individual dipole is decreasing faster than the monopole field and we could expect only a short range modification of the interparticle potential. However, in the mixed magnetic dipole–monopole plasma the role of the dipoles at large distances becomes nontrivial. Physically, the monopole fraction of the plasma leads to the formation of the stringlike structure between electrically charged particles while the dipole fraction of the plasma interacts with the electric field inside the string. This is the way how the dipoles modify the linear term of the interparticle potential at large distances. Below we study the influence of the dipole fraction of the gas on the long–range potential between electrically charged particles.

4
Let us consider the contribution of the monopole–dipole gas into the quantum average of the Wilson loop operator, \( W_C \), where contour \( C \) is the trajectory of the infinitely heavy test particle carrying the fundamental electric charge, \( g/2 \):

\[
W_C = \exp \left\{ i \frac{g}{2} \int d^3x \left( \rho(x) + (\rho^{(d)}(x), \vec{\partial}_x) \right) \eta^C(x)/(2\pi) \right\}.
\]  
(11)

The function \( \eta^C \) is defined as follows:

\[
\eta^C(x) = \pi \int_{\Sigma_C} d^2\sigma_{ij}(y) \varepsilon_{ijk} \partial_k D(x - y),
\]  
(12)

where the integration is taken over an arbitrary surface \( \Sigma_C \) spanned on the contour \( C \).

Substituting eq.(11) into eq.(1) and performing the transformations presented in Section 2 we derive the following representation for the contribution of the monopoles and dipoles to the Wilson loop:

\[
< W_C >_{\text{gas}} = \frac{1}{Z} \int D\varphi \exp \left\{ - \int d^3x \left[ \frac{g^2}{32\pi^2} (\bar{\partial}\varphi - \bar{\partial} \eta^C(x))^2 + 2\zeta_m (1 - \cos \varphi) + 4\pi \zeta_d \int_0^\infty dr F(r) \left( 1 - \frac{\sin(r |\bar{\partial}\varphi|)}{r |\bar{\partial}\varphi|} \right) \right]\right\},
\]  
(13)

where the expression in the square brackets has been normalized to zero at \( \varphi = 0 \) and we have shifted the field \( \varphi \rightarrow \varphi - \eta^C \).

In the weak coupling regime the densities of the monopoles are small and the Wilson loop in the leading order is given by the classical contribution to eq.(13). The corresponding classical equation of motion is:

\[
\bar{\partial}^2 \varphi - M_B^2 \sin \varphi - \frac{128\pi^3 \zeta_d}{g^2} \int_0^\infty dr F(r) \bar{\partial} \left[ \frac{\bar{\partial} \varphi}{(\bar{\partial} \varphi)^2} \left( \cos(r |\bar{\partial}\varphi|) - \frac{\sin(r |\bar{\partial}\varphi|)}{r |\bar{\partial}\varphi|} \right) \right]
= \pi \int_{\Sigma_C} d^2\sigma_{ij}(y) \varepsilon_{ijk} \partial_k \delta^{(3)}(x - y).
\]  
(14)

Below we consider the infinitely separated static charge and anticharge located at the points \((y, z) = (\pm\infty, 0)\). The \(x\)-coordinate is considered as a time coordinate. The string \( \Sigma_C \) in eq.(12) is chosen to be flat and is defined by the equation \( z = 0 \). Thus, the classical solution of eq.(14) does not depend on \( x \) and \( y \) coordinates, \( \varphi_{\text{cl}} = \varphi_{\text{cl}}(z) \). For the sake of simplicity we consider the case of the non–fluctuating dipole moments, \( G(r) = \delta(r - r_0) \) (below we use \( r \) instead of \( r_0 \)).

The classical contribution to the Wilson loop (13) gives the area low, \( < W_C >_{\text{gas}} = \text{const.} \exp\{-\text{Area } \sigma_{\text{cl}}\} \), where the classical string tension is

\[
\sigma_{\text{cl}} = g \frac{\sqrt{\rho_m}}{4\pi} \int_{-\infty}^{+\infty} d\xi \left[ \frac{1}{2} (\partial_\xi \varphi_{\text{cl}} - 2\pi \delta(\xi))^2 + 1 - \cos \varphi_{\text{cl}} + s \left( 1 - \frac{\sin(h |\partial_\xi \varphi_{\text{cl}}|)}{h |\partial_\xi \varphi_{\text{cl}}|} \right) \right],
\]  
(15)

\(^3\text{Note that the Wilson loop does not depend on the shape of this surface.}\)
and the $\varphi_{cl}$ is the solution of the rescaled classical equation of motion (14):

$$
\partial^2_\xi \varphi_{cl} - \sin \varphi_{cl} - s \partial^2_\xi \left[ \frac{1}{|\partial_\xi \varphi_{cl}|} \left( \cos(h|\partial_\xi \varphi_{cl}|) - \frac{\sin(h|\partial_\xi \varphi_{cl}|)}{h|\partial_\xi \varphi_{cl}|} \right) \right] = 2\pi \partial_\xi \delta(\xi). \tag{16}
$$

The dipole size in units of the Debye mass is denoted as

$$
h = r M_D, \tag{17}
$$

$\xi = z M_D$ is the rescaled $z$–coordinate and

$$
s = \frac{\rho_d}{\rho_m} \equiv \frac{2\pi \zeta_d}{\zeta_m} , \tag{18}
$$

is the ratio of the dipole density to the monopole density (”dipole fraction”), eq.(6).

In the absence of the dipoles, $s = 0$, the solution of eq.(16) and the corresponding string tension are [2]:

$$
\varphi^{(0)}_{cl} (\xi) = 4 \text{sign}(\xi) \arctan e^{-|\xi|} , \quad \sigma^{(0)}_{cl} = \frac{2g \sqrt{\rho_m}}{\pi}, \tag{19}
$$

respectively. In the presence of the dipoles the string tension is modified:

$$
\sigma_{cl} = \sigma^{(0)}_{cl} H(s, h). \tag{20}
$$

We study the function $H$ below.

In the limit of the small dipole size, $h \ll 1$, eq.(16) can be expanded in powers of $h$ up to the second order and the classical string tension can be found analytically. The correction factor $H$ is the square root of the dielectric constant (10),

$$
H(s, h) = \left[ 1 + \frac{s}{3} \left( h^2 + O(h^4) \right) \right]^{\frac{1}{2}}, \quad h \ll 1. \tag{21}
$$

Obviously, this correction is due to the renormalisation of the coupling constant $g^2 \to \epsilon g^2$.

Unfortunately, eq.(16) can not be solved analytically for general values of the dipole fraction $s$ and the dipole size $h$. Therefore we solve this equation and calculate the correction factor $H$ numerically. The function $H$ at various (fixed) dipole fractions $s$ is shown in Figure 1(a) by solid lines. The dashed lines indicate the behavior of correction factor $H$, eq.(21) valid at the small dipole sizes $h$.

The behavior of function $H$ can be understood as follows. The size of the flux of the electric field coming from the test particles is of the order of the inverse Debye mass. The dipoles with small size $h$ do not feel the flux structure and the contribution of the dipole fraction of the gas to the string tension is solely due the polarization of the media of the magnetic dipoles by the electric field, eq.(21). The correction factor $H$ is a rising function of $h$ since the increase of the dipole magnetic moment (size) leads to the enhancement of the mean polarization energy of the dipole in the external field.
Figure 1: (a) The correction factor $H$, eq.(20), to the string tension vs. the dipole size $h$ for the different dipole fractions $s$: $s = 0.1, 0.2, 0.5$; (b) the correction factor $Q$, eq.(23), vs. the dipole fraction $s$ for the dipole sizes $h = 0.1, 1, 3$. The numerical results are shown by the solid lines and the corresponding leading contributions at small $h$ are shown by the dashed lines.

At $h \approx 1$ the function $H(h)$ starts to deviate from eq.(21) since the dipoles begin to feel the width of the electric flux. At $h \approx 2$ the rise of the function $H$ slows down, Figure 1(a), and at large values of $h$ the correction factor grows logarithmically, $H \sim s \log h$. To understand this effect we note that the pure dipole gas can not generate a finite correlation length [6, 11] contrary to the pure monopole gas. Therefore one can expect that the dipole fraction of the monopole–dipole gas has no essential influence on the Debye mass $4$. As a result, the derivative $|\partial_\xi \varphi_{cl}|$ is of the order of unity at the center of the electric flux and the derivative vanishes out of the flux core as $e^{-C|\xi|}$, $C \sim 1$. One can easily check that the dipole contribution to the string tension (the last term in eq.(15)) in the large $h$ limit is equal to $s$ at $|\xi| \lesssim \log h$ and vanishes quickly out of this region. Since in this limit the other terms in eq.(16) do not grow with the enlargement of $h$, the string tension is given by the following formula,

$$\sigma = \frac{sg \sqrt{\rho_m}}{4\pi} \left( \log h + O(1) \right), \quad h \gg 1.$$  \hspace{1cm} (22)

Another interesting question is the dependence of the string tension $\sigma$ on the dipole fraction $s$ at fixed total number of the monopoles. The total (anti)monopole number includes both the free (anti)monopoles and the constituent (anti)monopoles of the magnetic

\footnote{A numerical study of the profiles of the field $\varphi_{cl}(z)$ shows that this is indeed the case for all studied values of $h$ and $s$.}
dipole states. Thus, $\rho_{\text{tot}} = \rho_m + 2\rho_d = \rho_m (1 + 2s)$, where the dipole fraction $s$ is defined in eq.(18). If a part of the free monopoles transforms into the dipole states the string tension acquires positive a contribution due to the increase of the dipole states (the function $H$ grows) and a negative contribution due to the decrease of the monopole density ($\sigma_{\text{cl}}^{(0)}$ diminishes). The correction factor $Q$ due to the monopole binding is defined by the following equation:

$$\sigma(s, h) = \sigma_{\text{tot}} Q(s, h), \quad Q(s, h) = (1 + 2s)^{-1/2} H(s, h).$$

(23)

where $\sigma_{\text{tot}} = 2g \sqrt{\rho_{\text{tot}} / \pi}$ is the string tension in the absence of the dipole states. Note that $\sigma_{\text{tot}}$ remains constant during the binding process.

In Figure 1(b) we show the function $Q$ vs. the dipole fraction $s$ for a few fixed values of the dipole sizes $h$. In accordance with physical intuition the string tension (solid lines) decreases due to the monopole binding. Moreover, the larger the dipole size the smaller the decrease is. This is not an unexpected fact since the contribution of the dipole to the Wilson loop grows with the increase of the dipole magnetic moment (equivalently, ”size” $h$ of the dipole).

Let us discuss the applicability of the pointlike dipole approximation to a physical system, for example, to the Georgi–Glashow model. The dipoles may be treated as point–like particles provided the typical distance between the plasma constituents is much larger than the dipole size $r$. The corresponding condition reads as follows:

$$(\rho_m + \rho_d) r^3 \ll 1.$$

(24)

Using eqs.(17,18) this equation can be rewritten as follows:

$$(1 + s) h^3 n_m \ll 1, \quad n_m = \rho_m M_D^{-3},$$

(25)

where $n_m$ is the number of the monopoles in the unit Debye volume. The string tension can be calculated in the classical approximation (15,16) only in the weak coupling limit which is equivalent to the condition $n_m \gg 1$. One can easily check that the condition $h \ll 1$ can be fulfilled in the Georgi–Glashow model and the correction factor (21) to the string tension can be arbitrarily large. The condition $h \gg 1$ can not be fulfilled and the result (22) is not valid in the Georgi–Glashow model in the weak coupling regime.

**Conclusion and acknowledgments**

We have studied the confining properties of the gas of the Abelian magnetic monopoles and the pointlike dipoles. Both the increase of the dipole fraction at a fixed monopole density and the enlargement of the dipole magnetic moment amplify the coefficient in front of the linear potential (“string tension”). At the same time the monopole binding at a fixed total number of magnetic charges leads to the natural decrease of the string tension.
The described effects may have interesting physical applications since the monopole–dipole gases are realized in various gauge theories. For example, the monopole binding mechanism leads to the high temperature deconfining phase transition in the three-dimensional compact electrodynamics and in the Georgi–Glashow model\textsuperscript{7}. Another example is the electroweak model in which the formation of the Nambu monopole–antimonopole pairs has been observed \textsuperscript{13} in the low temperature (Higgs) phase. In this model the string tension for the spatial Wilson loops is non-zero in the high temperature (symmetric) phase in which the monopoles form a gas. At low temperatures the string tension vanishes \textsuperscript{14}.

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