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Extended model of elementary particles based on an analogy with superconductivity

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An extended model of elementary particles based on an analogy with superconductivity is presented. Starting from the modified Nambu–Jona-Lasinio model of a massless fermion field with the quartic self-interactions, we derive coupled differential equations for the self-consistent fermion mass based on the relativistic Hartree-Fock-Bogoliubov approximation. The physical implication of these equations is explored and is shown to suggest a stringlike picture for the extended hadrons.

I. INTRODUCTION

We develop an extended model of elementary particles in this article by an analogy with superconductivity.[1] In our view, the hadronic vacuum is made up of an infinite number of fermion-antifermion pairs, and a hadron is a complicated object with a finite spatial extension where the pairing of fermions of the vacuum is locally destroyed. The shift in the pair potential makes it possible to trap quasiparticles within a limited region of three-dimensional space. In these respects our picture resembles Abrikosov's vortex line in a type-II superconductor,[2] i.e., a limited region of the normal state surrounded by the superconducting vacuum.

Here we present a general field-theoretical framework to describe such an extended object. For this purpose we first extend the superconductor model[1] of elementary particles to the spatially inhomogeneous systems. This may be viewed as the construction of a field-theoretical model based on an analogy with the Ginzburg-Landau-Abrikosov-Gorkov (GLAG) theory of superconductivity[3] which generalizes the BCS theory to describe spatially inhomogeneous systems such as a type-II superconductor or a superconductor with magnetic impurities. In the next section we start from the Lagrangian model of a massless spinor field with quartic self-interaction. We derive coupled differential equations for the self-consistent fermion mass by making use of the relativistic Hartree-Fock-Bogoliubov approximation. These equations turn out to have a striking similarity to the so-called Ginzburg-Landau (GL) equations,[3,4] though their physical interpretations are significantly different. These mass equations are solved to determine the self-consistent pair potential for the quasiparticles which are trapped in the potential to constitute an extended hadron. Our task is to extract the physical implications of the pair potential for the internal structure of the hadron. Though we have not completely analyzed the equations yet, the immediate conclusion is that they suggest the stringlike structure of the hadrons, i.e., the potential deviates from its asymptotic value only in the one-dimensional region. This seems to be the most significant achievement of our investigations. In the next section we derive our mass equations by making use of the relativistic Hartree-Fock-Bogoliubov approximation, and in Sec. III we discuss their implications on the hadron structure. Section IV is devoted to the summary and the conclusion.

II. DERIVATION OF MASS EQUATIONS

Let us take as our model the following Lagrangian of the massless fermion with the quartic self-interaction:

$$L(x)=\bar\psi(x)i\gamma\cdot\partial\psi(x)+g[(\bar\psi\psi)^2-(\bar\psi\gamma_5\psi)^2]$$
$$-g'[(\bar\psi\gamma_\mu\psi)^2+(\bar\psi\gamma_5\gamma_\mu\psi)^2]\,. \qquad (1)$$

This is the most general form of four-fermion interactions invariant under the chirality transformation. Dimensional coupling constants g and g' are assumed to be positive and need not necessarily be the same. The relation between their magnitudes will be derived afterwards. This theory is, of course, not renormalizable. We are forced to introduce a cutoff as usual [see Eq. (10)]. We are not going to discuss in this paper whether this cutoff corresponds to a certain fundamental constant, to the mass of a heavy vector meson, or to something else. In the original Nambu-Jona-Lasinio model there is no term proportional to g', which we have inserted to achieve a sufficient attractive force in the axial-vector channel. An alternative form of Eq. (1) is obtained on making a Fierz transformation,

$$L(x)=\bar{\psi}(x)i\gamma\cdot\partial\psi(x)-(\tfrac{1}{2}g+g')(\psi\gamma_\mu\psi)^2 -(g'-\tfrac{1}{2}g)(\bar{\psi}\gamma_5\gamma_\mu\psi)^2. \tag{2}$$

In the relativistic Hartree-Fock-Bogoliubov approximation, the self-consistent fermion mass is determined in such a way as to cancel the self-energy effects caused by the self-interaction in Eqs. (1) and (2). If we make use of the alternative expressions (1) and (2), the self-consistent mass is given by

$$\begin{aligned} m(x)=&\,2gi[\mathrm{Tr}(S_F(x,x))-\gamma_5\mathrm{Tr}(\gamma_5S_F(x,x))]\\ &-(4g'+g)i\gamma^\mu\mathrm{Tr}(\gamma_\mu S_F(x,x))\\ &-(4g'+g)i\gamma_5\gamma^\mu\mathrm{Tr}(\gamma_5\gamma_\mu S_F(x,x)). \end{aligned}\tag{3}$$

This may be written as

$$m(x)\equiv m_S(x)+i\gamma_5m_P(x)+\gamma^\mu m_\mu^V(x)+\gamma_5\gamma^\mu m_\mu^A(x), \tag{4}$$

where $m_S(x)$, $m_P(x)$, $m_V(x)$, and $m_A(x)$ are given by

$$m_S(x)=2gi\mathrm{Tr}(S_F(x,x)), \tag{5}$$

$$m_P(x)=2gi\mathrm{Tr}(i\gamma_5S_F(x,x)), \tag{6}$$

$$m_V^\mu(x)=-(4g'+g)i\mathrm{Tr}(\gamma^\mu S_F(x,x)), \tag{7}$$

$$m_A^\mu(x)=-(4g'-g)i\mathrm{Tr}(\gamma_5\gamma^\mu S_F(x,x)). \tag{8}$$

Here m_S, m_P, m_V, and m_A are assumed to be real. The propagator $S_F(x,y)$ satisfies

$$i\gamma\cdot\partial_xS_F(x,y)-m(x)S_F(x,y)=\delta(x-y). \tag{9}$$

Equations (4)–(9) form the basic ingredients of our theory. Note that in the Nambu theory the vacuum is taken to be Lorentz-invariant and hence $m_V(x)$ and $m_A(x)$ vanish identically. This is no longer true in our case.

We now attempt to solve Eqs. (4)–(9) approximately. Since we expect our hadrons to have only finite extensions, we assume that $m(x)$ goes to the Nambu value m_∞ when $x\to\infty$. m_∞ is obtained by solving the Eqs. (4)–(9) assuming a constant mass,

$$1=\frac{2gi}{(2\pi)^4}\int^\Lambda\frac{\mathrm{Tr}1}{p^2-m_\infty^2+i\epsilon}d^4p, \tag{10}$$

where the integral over the momentum is cut off at $\Lambda^2=p^2$.

Then we can formally solve Eqs. (4)–(9) by expanding the full propagator $S_F(x,y)$ around $S_F^{m\infty}(x,y)$,

$$\begin{aligned} S_F(x,x)=&\,S_F^{m\infty}(x,x)+\int S_F^{m\infty}(x,x')[m(x')-m_\infty]S_F^{m\infty}(x',x)d^4x'\\ &+\int S_F^{m\infty}(x,x')[m(x')-m_\infty]S_F^{m\infty}(x',x'')[m(x'')-m_\infty]S_F^{m\infty}(x'',x)d^4x'\,d^4x''\\ &+\int S_F^{m\infty}(x,x')[m(x')-m_\infty]S_F^{m\infty}(x',x'')[m(x'')-m_\infty]S_F^{m\infty}(x'',x''')[m(x''')-m_\infty]S_F^{m\infty}(x''',x)d^4x'd^4x''d^4x'''\\ &+\cdots. \end{aligned}\tag{11}$$

We find that each term in Eq. (11) is a loop integration with a larger number of mass insertions as we go up to higher-order terms. Hence, in the limit $\Lambda\to\infty$ only the first few terms in Eq. (11) diverge, while all the other terms remain convergent. A great simplicity arises, therefore, if we consider our theory in the limit of large Λ which we take to be the case hereafter. Then by retaining the terms only up to the order $\log\Lambda$, i.e., by neglecting all the other terms which stay finite in the limit of $\Lambda\to\infty$, we are able to perform the calculations in a closed form.

After some lengthy but rather straightforward calculations we obtain (see Appendix)

$$\Box m_S(x)-2m_\infty^2m_S(x)+2m_S^3(x)+2m_P^2(x)m_S(x)-4m_A^2m_S(x)-2m_P(x)\partial_\mu m_A^\mu(x)-4\partial_\mu m_P(x)m_A^\mu(x)=0, \tag{12}$$

$$\Box m_P(x)-2m_\infty^2m_P(x)+2m_S^2(x)m_P(x)+2m_P^3(x)-4m_A^2m_P(x)+2m_S(x)\partial_\mu m_A^\mu(x)+4\partial_\mu m_S(x)m_A^\mu(x)=0, \tag{13}$$

$$(\Box g_{\mu\nu}-\partial_\mu\partial_\nu)m_V^\nu(x)+\frac{g}{(4g'+g)L}m_{V\mu}(x)=0, \tag{14}$$

$$(\Box g_{\mu\nu}-\partial_\mu\partial_\nu)m_A^\nu(x)+3[\partial_\mu m_P(x)m_S(x)-\partial_\mu m_S(x)m_P(x)]+6[m_S(x)^2+m_P(x^2)]m_{A\mu}(x)+\frac{g}{(4g'-g)L}m_{A\mu}(x)=0. \tag{15}$$

Here L is a positive infinitesimal quantity,

$$L=-\tfrac{4}{3}gi\frac{1}{(2\pi)^4}\int^\Lambda d^4p\frac{1}{(p^2-m_\infty^2+i\epsilon)^2}. \tag{16}$$

Note that Eqs. (12)–(15) imply

$$\partial_\mu m_V^\mu(x)=0, \tag{17}$$

and

$$\partial_\mu m_A^\mu(x)=0. \tag{18}$$

Equations (17) and (18) can also be deduced from Eq. (9) directly. The following remarks are in order on the above mass equations.

(i) While there is no cutoff dependence in the mass equations for $m_S(x)$ and $m_P(x)$, there appear cutoff-dependent mass terms in the equations for $m_V(x)$ and $m_A(x)$. This is because in Eqs. (5) and (6) the quadratically divergent terms cancel between the right- and left-hand sides of the equations. By setting the coefficients of the logarithmically divergent terms equal to zero we get Eqs. (12) and (13) (see Appendix). On the other hand, in the calculations of the vacuum-polarization tensors, the quadratically divergent terms vanish due to gauge invariance and hence there arises a mismatch in the order of Λ between the right- and left-hand sides of Eqs. (7) and (8). This is the cause of the large-mass terms in Eqs. (14) and (15). Note that we have to set $g' > \frac{1}{4}g$ in order for these masses to be real.

(ii) As one can see from Eqs. (12)–(15), the vector density decouples completely from the rest. This is reasonable because in our theory the vacuum and quasiparticles retain a well-defined charge-conjugation property and hence the nonzero $m_V(x)$ would lead to C violation. Since we do not want C violation in our theory we set $m_V = 0$. Next let us rewrite Eqs. (12), (13), and (15) in a more transparent fashion. By defining

$$m_S(x) + i m_P(x) \equiv \phi(x)\,, \tag{19}$$

we obtain

$$[\partial_\mu + 2i m_{A\mu}(x)]^2 \phi(x) - 2m_\infty{}^2 \phi(x) + 2\,|\,\phi(x)\,|^2\,\phi(x) = 0\,, \tag{20}$$

$$(\Box g_{\mu\nu} - \partial_\mu \partial_\nu) m_A^\nu(x) + \frac{3}{2i}\,[\phi^*(x)\partial_\mu \phi(x) - \phi(x)\partial_\mu \phi^*(x)] + 6\,|\,\phi(x)\,|^2\, m_{A\mu}(x) + \frac{g}{(4g'-g)L}\, m_{A\mu}(x) = 0\,. \tag{21}$$

These two equations may be deduced from the Lagrangian

$$L(x) = |\,[\partial_\mu + 2i m_{A\mu}(x)]\,\phi(x)\,|^2 - \tfrac{1}{3} F_{\mu\nu}(x) F^{\mu\nu}(x) + \frac{2g}{3(4g'-g)L}\, m_A^\mu(x)^2 + 2m_\infty{}^2\,|\,\phi(x)\,|^2 - |\,\phi(x)\,|^4\,; \tag{22}$$

here

$$F_{\mu\nu} = \partial_\nu m_{A\mu} - \partial_\mu m_{A\nu}\,.$$

(iii) Our next remark is on the problem of the massless pseudoscalar mode. The asymptotic form of the solution to Eq. (20) shows that there exists no long-range pair potential in our theory. This means that the massless pseudoscalar mode does not imply a long-range force. In contrast with the case of Freundlich and Lurié,[5] the massless pseudoscalar mode is not absorbed into the external axial-vector gauge field. The original scalar field of mass $2m_\infty$ and the massless pseudoscalar field combine into a complex scalar field $\phi(x)$ whose radial $|\phi(x)|$ and angular $\arg\phi(x)$ components are now the eigenstates of the mass with the eigenvalues $2m_\infty$ and 0, respectively.

At this point readers may note the striking similarity of Eqs. (20)–(22) to the GL equations and suspect that our derivation is a field-theoretic analog of Gorkov's derivation of the GL equation from the microscopic theory of superconductivity near the critical temperature.[6] The crucial difference between the GLAG theory of superconductivity and our mass equations exists, however, in the following point: In contrast with the case of the GL equation the quantity $m_A(x)$ in our theory is not an external field. It describes the collective excitations of quasiparticle pairs just as the scalar and the pseudoscalar densities do. We should also remark that the similarity between our mass equations and the GL equations already suggests the string-like picture for the extended hadrons.

III. EXTENDED MODEL OF HADRONS

Our extended model of hadrons is formulated in two steps. The first is to solve the mass equations (12), (13), and (15) and determine the one-particle Hartree-Fock-Bogoliubov field. This provides us an inhomogeneous, locally excited vacuum to trap quasiparticles in the Hartree potential. The next step is to solve the potential problem

$$i\gamma\cdot\partial\psi(x) = m(x)\psi(x)\,, \tag{23}$$

to determine various bound and continuum states. $\psi(x)$ is decomposed into a sum of the orthonormal set of solutions to Eq. (23),

$$\psi(x) = \sum a_i u_i(x) + \sum b_j^\dagger v_j(x)\,. \tag{24}$$

Our excited vacuum is defined by

$$a_i\,|0\rangle = 0\,, \quad b_j\,|0\rangle = 0\,. \tag{25}$$

The scalar and other densities characterize its internal structures,

$$2gi\,\mathrm{Tr}\,\langle 0\,|\tfrac{1}{2}[\bar\psi(x), \psi(x)]\,|0\rangle = -m_S(x)\,, \tag{26}$$

$$2gi\,\mathrm{Tr}\langle 0\,|\tfrac{1}{2}[\bar\psi(x), i\gamma_5\psi(x)]\,|0\rangle = -m_P(x)\,, \tag{27}$$

$$-(4g'-g)i\,\mathrm{Tr}\langle 0\,|\tfrac{1}{2}[\bar\psi(x), \gamma_5\gamma_\mu\psi(x)]\,|0\rangle = -m_\mu^A(x)\,, \tag{28}$$

$$\langle 0\,|[\bar\psi(x), \gamma_\mu\psi(x)]\,|0\rangle = 0\,. \tag{29}$$

Our extended hadrons are constructed by applying to our vacuum the creation operators of bound states determined by Eq. (23),

$$|\text{hadron}\rangle = \prod_{i,j}^{\text{bound}} a_i^\dagger \cdots b_j^\dagger \cdots |0\rangle . \tag{30}$$

What picture should we have for our excited vacuum? Since $m_A^\mu(x)$ has an infinitesimal penetration depth due to its infinite mass, the region where it is different from zero is limited to the stringlike one-dimensional configuration (the origin of the stringlike structure will be discussed below). Along this stringlike extension $m_S(x)$ and $m_P(x)$ also deviate from Nambu-BCS values. Therefore we can picture our vacuum as the Nambu-BCS vacuum broken in a cylindrical region in space where the pairing of fermions and antifermions is destroyed and their spins are lined up to yield nonzero spin density $m_A(x)$. The radius of the cylindrical region is $1/2m_\infty$, as is known in the GL theory. Since the chirality current is conserved $[\partial_\mu m_A^\mu(x)=0]$, the cylinder cannot have open ends, i.e., it is a closed loop (see Fig. 1).

The origin of the stringlike structure is understood in the following way. As we have already noted in the preceding section, in the calculations of the right-hand sides of Eqs. (7) and (8) the quadratically divergent term vanishes due to the gauge invariance. This is physically reasonable since $m_V(x)$ and $m_A(x)$ are identically zero in the case of homogeneous vacuum and hence their contributions come only from the locally excited part of our inhomogeneous vacuum. Because these contributions are logarithmically divergent we can conclude that the excitation of the vacuum occurs in the one-dimensional region on the basis of the dimensional argument.

IV. SUMMARY AND CONCLUSION

Starting from a Lagrangian model of a fundamental massless spinor field with self-interaction, we have developed above an extended model of hadrons based on an analogy with superconductivity. In the foregoing sections we have derived self-consistent mass equations for the quasiparticles and have outlined the qualitative features of the solutions. Most strikingly these equations are shown to suggest a string picture of the hadrons.

Recently there has been much interest in constructing the field-theoretical models of extended hadrons.[7] In these theories authors generally assume the existence of the so-called Higgs scalar meson and attempt to build an extended object with three-dimensional extension. The stringlike picture of our theory seems to be in sharp contrast to these theories. It will be extremely interesting if our theory indeed has some fundamental connection with the conventional string model of hadrons.[8]

Finally we should comment on the recovery of the lost symmetries in our one-particle Hartree field which we have left unsolved in this paper. In contrast with the case of the Nambu-BCS theory our one-particle field violates Lorentz invariance as well as the invariance under chirality transformations. It is therefore necessary to restore the translational and the rotational symmetries of our original Lagrangian by diagonalizing the residual interactions. This will be achieved, for instance, by the use of the relativistic generalization of the generator coordinate methods of the theory of nuclear collective motions.[9] Then, as an analog of the nuclear rotational levels, Regge recurrence will appear. All these problems need further investigation.

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APPENDIX

Here we outline the derivation of the mass equations, taking Eq. (12) as an example. All the other equations can be derived in a more or less similar manner [however, note remark (i) of Sec. II].

We start from Eq. (5):

$$m_S(x) = 2gi\, \mathrm{Tr}\big(S_F(x,x)\big) , \tag{A1}$$

where

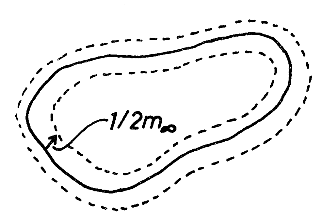

FIG. 1. Schematic picture of a typical solution of the mass equations. On the closed string shown by the solid line $m_\mu^A(x)$ has nonzero values. Inside the cylinder shown by the dashed lines $\phi(x)$ deviates from its asymptotic value m_∞. On the solid line $\phi(x)$ goes to zero.

$$\begin{aligned}S_F(x,x)=S_F^{m_\infty}(x,x)&+\int S_F^{m_\infty}(x,x')[m(x')-m_\infty]S_F^{m_\infty}(x',x)d^4x'\\&+\int S_F^{m_\infty}(x,x')[m(x')-m_\infty]S_F^{m_\infty}(x',x'')[m(x'')-m_\infty]S_F^{m_\infty}(x'',x)d^4x'd^4x''\\&+\int S_F^{m_\infty}(x,x')[m(x')-m_\infty]S_F^{m_\infty}(x',x'')[m(x'')-m_\infty]S_F^{m_\infty}(x'',x''')[m(x''')-m_\infty]\\&\quad\times S_F^{m_\infty}(x''',x)d^4x'\,d^4x''\,d^4x'''\\&+\cdots.\end{aligned}\tag{A2}$$

The contribution of the first term to (A1) is simply equal to m_∞. The second term corresponds to the Feynman diagram Fig. 2(a). The fermion is scattered by the effective potential $M(x')\equiv m(x')-m_\infty$ and returns to its original position. The third and the fourth terms correspond to the diagrams 2(b) and 2(c), respectively. All the other terms converge when Λ goes to infinity and can be neglected as explained in Sec. II. By the Fourier transform

$$M_S(x)\equiv m_S(x)-m_\infty=\frac{1}{(2\pi)^4}\int e^{-iqx}M_S(q)d^4q\,,\tag{A3}$$

Eq. (A1) becomes

$$M_S(q)=M_S^{(2)}(q)+M_S^{(3)}(q)+M_S^{(4)}(q)\,,\tag{A4}$$

where $M_S^{(2)}(q)$, $M_S^{(2)}(q)$, and $M_S^{(4)}(q)$ are the contributions from the diagrams 2(a), 2(b), and 2(c), respectively. The results of the calculations are

$$M_S^{(2)}(q)=M_S(q)[1+L'(2m_\infty^2-\tfrac{1}{2}q^2)]\,,\tag{A5}$$

where

$$L'=8gi\int\frac{d^4p}{(2\pi)^4}\frac{1}{(p^2-m_\infty^2)^2}\,,\tag{A6}$$

$$\begin{aligned}M_S^{(3)}(q)=\Big[&3m_\infty\int M_S(q')M_S(q-q')\frac{d^4q'}{(2\pi)^4}+m_\infty\int M_P(q')M_P(q-q')\frac{d^4q'}{(2\pi)^4}\\&+i\int(q+q')_\nu M_P(q')M_A^\nu(q-q')\frac{d^4q'}{(2\pi)^4}-2m_\infty\int M_A^\nu(q')M_A^\nu(q-q')\frac{d^4q'}{(2\pi)^4}\Big]L'\,,\end{aligned}\tag{A7}$$

where

$$m_P(x)=\int\frac{d^4q}{(2\pi)^4}M_P(q)e^{-iqx}\,,\tag{A8}$$

and

$$m_A^\nu(x)=\int\frac{d^4q}{(2\pi)^4}M_A^\nu(q)e^{-iqx}\,.\tag{A9}$$

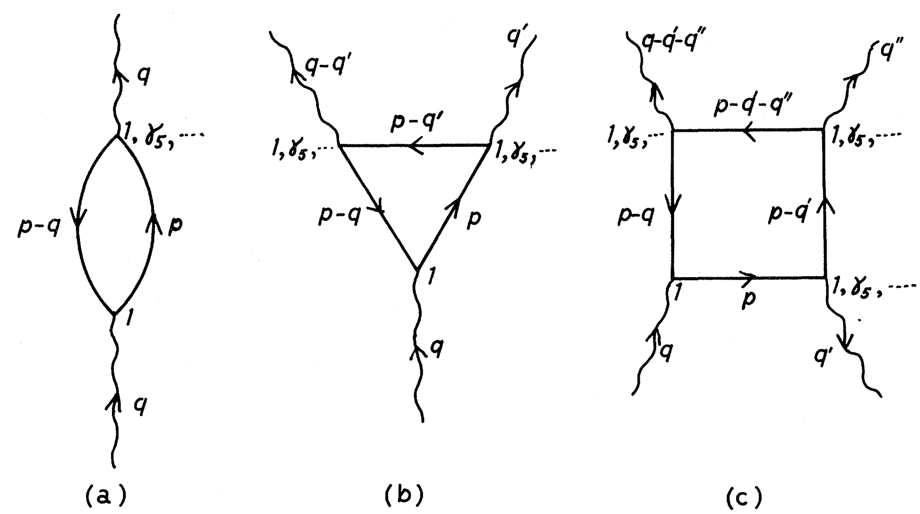

FIG. 2. Feynman diagrams used in the calculation of mass equations.

Finally

$$M_S^{(4)}(q)=\left[\iint M_S(q')M_S(q'')M_S(q-q'-q'')\frac{d^4q'}{(2\pi)^4}\frac{d^4q''}{(2\pi)^4}+\iint M_S(q')M_P(q'')M_P(q-q'-q'')\frac{d^4q'}{(2\pi)^4}\frac{d^4q''}{(2\pi)^4}\right.$$
$$\left.-2\iint M_S(q')M_A^{\nu}(q'')M_{A\nu}(q-q'-q'')\frac{d^4q'}{(2\pi)^4}\frac{d^4q''}{(2\pi)^4}\right]L'. \qquad \text{(A10)}$$

Substituting (A5), (A7), and (A10) into (A4) and transforming it back to x space we obtain Eq. (12) of Sec. II. We note that the first term of (A5) is exactly equal to the left-hand side of (A4).

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Strings, monopoles, and gauge fields

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The Nielsen-Olesen interpretation of dual strings as Abrikosov flux lines is extended to the case of open-ended strings by adapting Dirac's description of magnetic monopoles to a London-type theory. The mathematical formalism turns out to be similar to that of Kalb and Ramond. Translated to hadron physics, it implies that the quarks will act as carriers of magnetic charge, permanently bound in pairs by the string bonds. However, massive axial-vector gluons can be created by hadrons.

I. INTRODUCTION

In a very interesting paper[1] Nielsen and Olesen have pointed out a parallelism between the Higgs model of broken gauge invariance and the Landau-Ginzburg theory of superconductivity on the one hand and the dual string model and the Abrikosov flux lines in type II superconductors on the other. According to their suggestion, a dual string is nothing but a mathematical idealization of a magnetic flux tube in equilibrium against the pressure of the surrounding charged superfluid (Higgs-scalar field) which it displaces. Only strings with no ends (infinite strings or loops) were considered by them. It is known that a closed string could be a candidate for the Pomeron. But what will happen if the string is open-ended? Obviously the magnetic flux will terminate at the end points, thus creating a pair of magnetic charges.[2] In the dual quark model ordinary hadrons are viewed as being made up of quarks bound by dual strings, or, from the string's point of view, as open strings having