The Effective Lagrangian of QED with a Magnetic Charge and Dyon Mass Bounds

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Abstract

The effective Lagrangian of QED coupled to dyons is calculated. The resulting generalization of the Euler-Heisenberg Lagrangian contains non-linear $P$ and $T$ noninvariant terms corresponding to the virtual pair creation of dyons. The corresponding $P$ and $T$ violating part of the matrix element for light-by-light scattering is considered. This effect induces an electric dipole moment for the electron, of order $M^{-2}$, where $M$ is the dyon mass. The current limit on the electric dipole moment of the electron yields the lower dyon mass bound $M > 1$ TeV.
I. INTRODUCTION

Very precise measurements achieved during the last decade have opened up for a new approach in elementary particle physics. According to this, evidence of new particles can be extracted from indirect measurements of their virtual contribution to processes at energies which are too low for direct production. For example, the top quark mass as predicted from precision electroweak data [1] agrees to within 10% with direct experimental measurements [2].

This approach has recently been applied [3] for the estimation of possible virtual monopole contributions to observables at energies below the monopole mass. One-loop dyon-induced quantum corrections to the QED Lagrangian were discussed in [4]. Taking into account the violation of parity (and time-reversal symmetry) in a theory with monopoles [5], the emergence of an electric dipole moment was first pointed out by Purcell and Ramsey [6]. More recently, the effect due to monopole loop contributions has been discussed [7,8].

The calculation of quantum corrections due to the virtual pair creation of dyons is a very difficult problem because the standard diagram technique is not valid in this case. The difficulty is connected both to the large value of the magnetic charge of the dyon and the lack of a consistent local Lagrangian formulation of electrodynamics with two types of charge (see e.g. [9] and references therein). So, there is no possibility to use a perturbation expansion in a coupling constant. But one can apply the loop expansion which is just an expansion in powers of the Planck constant $\hbar$.

II. EFFECTIVE LAGRANGIAN

It is known (see, e.g. [10]) that the one-loop quantum correction to the QED Lagrangian can be calculated without the use of perturbation methods. The correction is just the change in the vacuum energy in an external field. Let us review the simple case of weak constant parallel electric and magnetic fields $\mathbf{E}$ and $\mathbf{H}$. We impose the conditions $e|\mathbf{E}|/m^2 \ll 1$ and
$e|H|/m^2 \ll 1$ such that the creation of particles is not possible. In this case the one-loop correction can be calculated by summing the one-particle modes — the solutions of the Dirac equation in the external electromagnetic field — over all quantum numbers [10], [11]. For example, if there is just a magnetic field, $H = (0, 0, H)$, the corresponding equation is

$$[i\gamma^\mu(\partial_\mu + ieA_\mu) - m]\psi(x) = 0$$

(1)

where the electromagnetic potential is $A^\mu = (0, -Hy, 0, 0)$. The solution to this equation gives the energy levels of an electron in a magnetic field [12,13]

$$\varepsilon_n = \sqrt{m^2 + eH(2n - 1 + s) + k^2}$$

(2)

where $n = 0, 1, 2\ldots$, $s = \pm 1$, and $k$ is the electron momentum along the field. In this case the correction to the Lagrangian is [10,12]

$$\Delta L_H = \frac{eH}{2\pi^2} \sum_{n=1}^{\infty} \int_0^\infty dk \left( (m^2 + k^2)^{1/2} + 2 \sum_{n=1}^{\infty} (m^2 + 2eHn + k^2)^{1/2} \right)$$

$$= -\frac{1}{8\pi^2} \int_0^\infty ds e^{-ms} \left( (esH) \coth(esH) - 1 - \frac{1}{3}e^2 s^2 H^2 \right),$$

(3)

where the terms independent of the external field $H$ are dropped and a standard renormalization of the electron charge has been made [10].

If we consider simultaneously magnetic ($H$) and electric ($E$) homogeneous fields, then equation (1), as well as its classical analogue can be separated into two uncoupled equations, each in two variables [13]. Indeed, in this case we can take $A^\mu = (Ez, -Hy, 0, 0)$ and the interactions of an electron with the fields $E$ and $H$ are determined independently. For such a configuration of electromagnetic fields the correction to the Lagrangian is (see [10], p. 787)

$$\Delta L = \frac{eH}{2\pi^2} \sum_{n=1}^{\infty} \int_0^\infty dk \varepsilon_n^{(E)}(k).$$

(4)

Here $\varepsilon_n^{(E)}$ is the correction to the energy of an electron in the combined external magnetic and electric fields, which is in the first order proportional to $e^2 E^2$.

So, the total Lagrangian is $L = L_0 + \Delta L$, where $L_0 = (E^2 - H^2)/2$ is just the Lagrangian of the free electromagnetic field in the tree approximation, and can be written as
\[ L = \left( 1 + \frac{\alpha}{3\pi} \int_0^\infty \frac{ds}{s} e^{-m^2 s} \right) \frac{E^2 - H^2}{2} + \Delta L'. \] (5)

The logarithmic divergency can be removed by the standard renormalization of the external fields and the electron charge:

\[ E_\text{reg} = Z_{-1/2}^{-3} E; \quad H_\text{reg} = Z_{-1/2}^{-3} H; \quad e_\text{reg} = Z_{1/2}^{3} e, \] (6)

where \( Z_{-1}^{3} = 1 + \left( \frac{\alpha}{3\pi} \right) \int_0^\infty \frac{ds}{s} e^{-m^2 s} \) is the usual QED renormalization factor. Thus the finite part of the correction to the Lagrangian \( \Delta L' \) can be written in terms of physical quantities as (see [10], p. 790)

\[ \Delta L' = -\frac{1}{8\pi^2} \int_0^\infty \frac{ds}{s^3} e^{-m^2 s} [(esE)(esH) \cot(esE) \coth(esH) - 1], \] (7)

which in the limit \( E = 0 \) reduces to the renormalized form of (3).

The series expansion of (7) in terms of the parameters \( eE/m^2 \ll 1, eH/m^2 \ll 1 \) yields the well known Euler-Heisenberg correction [14]:

\[ \Delta L' \approx \frac{e^4}{360\pi^2 m^4} \left[ (H^2 - E^2)^2 + 7(HE)^2 \right], \] (8)

where \( e^2 = \alpha \).

Let us consider how the situation changes if we consider the virtual pair creation of dyons in the external electromagnetic field. Using an analogy with the classical Lorentz force on a dyon of velocity \( v \) with electric (\( Q \)) and magnetic (\( g \)) charges [9]

\[ F = QE + gH + v \times (QH - gE), \] (9)

we shall assume that the wave equation for this particle in an external electromagnetic field can be expressed as [16,4]

\[ (i\gamma^\mu D_\mu - M)\psi(x) = 0, \] (10)

where \( M \) is the dyon mass, and \( iD_\mu \) a generalized momentum operator, with \( D_\mu = \partial_\mu + iQA_\mu + igB_\mu \).
The potential \( A_\mu \) and its dual \( B_\mu \) are defined by \( F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu = \varepsilon_{\mu\nu\rho\sigma} \partial^\rho B^\sigma \) where \( F_{\mu\nu} \) is the electromagnetic field strength tensor\(^1\) and \( \varepsilon_{0123} = 1 \). The potentials in the case of constant parallel electric and magnetic fields can be expressed as

\[
A^\mu = (E_z, -H_y, 0, 0), \quad B^\mu = (H_z, E_y, 0, 0). \tag{11}
\]

It is easily seen that the solution to the equation of motion for a dyon in an external electromagnetic field can be obtained from the solution to the equation for an electron (1) by the following substitution

\[
eE \toQE + gH; \quad eH \toQH - gE. \tag{12}
\]

Using the same substitution as in Eqs. (5) and (7), we obtain the following expression for the quantum correction to the Lagrangian, due to the vacuum polarization caused by dyons:

\[
L = \left( 1 + \frac{Q^2}{12\pi^2} \int_0^\infty \frac{ds}{s} e^{-M^2 s} - \frac{g^2}{12\pi^2} \int_0^\infty \frac{ds}{s} e^{-M^2 s} \right) \frac{E^2 - H^2}{2} + \Delta L', \tag{13}
\]

where a total derivative has been dropped.

For the renormalization of this expression we can introduce the renormalization factors [16]

\[
Z_e^{-1} = 1 + \frac{Q^2}{12\pi} \int_0^\infty \frac{ds}{s} e^{-M^2 s}, \quad Z_g^{-1} = 1 - \frac{g^2}{12\pi} \int_0^\infty \frac{ds}{s} e^{-M^2 s}, \tag{14}
\]

which are generalizations of the definition \( Z_3 \) of Eq. (6). In this case the fields and charges are renormalized as \([16,17]\]

\[
E_{\text{reg}}^2 = Z_e^{-1} Z_g^{-1} E^2; \quad H_{\text{reg}}^2 = Z_e^{-1} Z_g^{-1} H^2; \quad e_{\text{reg}}^2 = Z_e Z_g e^2; \quad g_{\text{reg}}^2 = Z_e^{-1} Z_g^{-1} g^2. \tag{15}
\]

This relation (15) means that the vacuum of electrically charged particles shields the external electromagnetic field but the contribution from magnetically charged particles antishields it. This agrees with the results of \([18]\) and \([19]\).

\[^1\]This definition is consistent only if \( \Box A_\mu = \Box B_\mu = 0 \), i.e., for constant electromagnetic fields or for free electromagnetic waves.
Considering now the case of weak electromagnetic fields, the finite part of the Lagrangian \( \Delta L' \), can, by analogy with (8), be written as

\[
\Delta L' = \frac{1}{360\pi^2 M^4} \{ [(Q^2 - g^2)^2 + 7Q^2g^2](H^2 - E^2)^2 + [16Q^2g^2 + 7(Q^2 - g^2)^2](HE)^2 \\
+ 6Qg(Q^2 - g^2)(HE)(H^2 - E^2) \}.
\]

(16)

The expressions (8) and (16) describe nonlinear corrections to the Maxwell equations which correspond to photon-photon interactions. The principal difference between the formula (16) and the standard Euler-Heisenberg effective Lagrangian consists in the appearance of \( P \) and \( T \) non-invariant terms proportional to \((HE)(H^2 - E^2)\). It should however be noted that this term is invariant under charge conjugation \( C \), since then both \( Q \) and \( g \) would change sign.

If we consider separately the virtual creation of dyon pairs, then because of invariance of the model under a dual transformation (see, e.g. [9]) the physics is determined not by the values \( Q \) and \( g \) separately, but by the effective charge \( \sqrt{Q^2 + g^2} \). In the same way the operations of \( P \) and \( T \) inversions are modified. However we will consider simultaneously the contributions from vacuum polarization by electron-positron and dyon pairs. In this case it is not possible to reformulate the theory in terms of just one effective charge by means of a dual transformation. Moreover the Dirac charge quantization condition connects just the electric charge of the electron and the magnetic charge of a dyon: \( e g = n/2 \) whereas the electric charge \( Q \) is not quantized.

It is widely believed, based both on experimental bounds and theoretical predictions [20] that the dyon mass would be large, \( M \gg m \), where \( m \) is the electron mass. Thus, in the one-loop approximation the first non-linear correction to the QED Lagrangian from summing the contributions (8) and (16) can be written as

\[
\Delta L' \approx \frac{e^4}{360\pi^2 m^4} [(H^2 - E^2)^2 + 7(HE)^2] + \frac{Qg(Q^2 - g^2)}{60\pi^2 M^4}(HE)(H^2 - E^2),
\]

(17)

where the \( P \) and \( T \) invariant terms corresponding to vacuum polarization by dyons have been dropped because they are suppressed by factors \( M^{-4} \). Thus, their contribution to
the effective Lagrangian will be of the same order as that of the ordinary QED multiloop amplitudes which we neglect.

III. PHOTON-PHOTON SCATTERING

Expression (17) yields the matrix element for low-energy photon-photon scattering. In order to determine it, we substitute into (17) the expansion

$$ F_{\mu\nu}(x) = \frac{i}{(2\pi)^4} \int d^4q \ (q_{\mu} A_{\nu} - q_{\nu} A_{\mu}) \ e^{iqx}. $$

(18)

Corresponding to the second term of (17) we find

$$ \frac{Qg(Q^2 - g^2)}{480\pi^2 M^4} \epsilon_{\mu\rho\sigma} F^{\mu\nu} F^{\rho\sigma} F_{\alpha\beta} F^{\alpha\beta} $$

$$ = \frac{1}{(2\pi)^{12}} \int d^4q_1 d^4q_2 d^4q_3 d^4q_4 \delta(q_1 + q_2 + q_3 + q_4) A_{\mu}(q_1) A_{\nu}(q_2) A_{\rho}(q_3) A_{\sigma}(q_4) \tilde{M}^{\mu\nu\rho\sigma}, $$

(19)

where

$$ \tilde{M}^{\mu\nu\rho\sigma} = \tilde{M}^{\mu\nu\rho\sigma}(q_1, q_2, q_3, q_4) = \frac{Qg(Q^2 - g^2)}{60\pi^2 M^4} \epsilon_{\alpha\beta}^{\mu\nu} q_1^{\alpha} q_2^{\beta} [q_3^{\rho} q_4^{\sigma} - q_3^{\sigma} q_4^{\rho}]. $$

(20)

Symmetrizing this pseudotensor one obtains the $P$ and $T$ violating part of the matrix element for light-by-light scattering. With all momenta flowing inwards, $k_1 + k_2 + k_3 + k_4 = 0$, the matrix element takes the form

$$ M'_{\mu\nu\rho\sigma} = \frac{1}{6} \{ \tilde{M}_{\mu\nu\rho\sigma}(k_1, k_2, k_3, k_4) + \tilde{M}_{\mu\rho\nu\sigma}(k_1, k_3, k_2, k_4) + \tilde{M}_{\mu\sigma\nu\rho}(k_1, k_4, k_2, k_3) $$

$$ + \tilde{M}_{\nu\mu\rho\sigma}(k_2, k_3, k_1, k_4) + \tilde{M}_{\nu\rho\mu\sigma}(k_2, k_4, k_1, k_3) + \tilde{M}_{\rho\sigma\mu\nu}(k_3, k_4, k_1, k_2) \} $$

$$ = \frac{Qg(Q^2 - g^2)}{60\pi^2 M^4} (\epsilon_{\alpha\beta}^{\mu\nu} k_1^{\alpha} k_2^{\beta} k_3^{\sigma} k_4^{\rho} + \epsilon_{\alpha\beta}^{\mu\rho} k_1^{\alpha} k_2^{\sigma} k_3^{\beta} k_4^{\nu} + \epsilon_{\alpha\beta}^{\mu\sigma} k_1^{\alpha} k_2^{\rho} k_3^{\nu} k_4^{\beta} $$

$$ + \epsilon_{\alpha\beta}^{\nu\rho} k_1^{\alpha} k_2^{\sigma} k_3^{\nu} k_4^{\rho} + \epsilon_{\alpha\beta}^{\nu\sigma} k_1^{\alpha} k_2^{\rho} k_3^{\nu} k_4^{\beta} + \epsilon_{\alpha\beta}^{\rho\sigma} k_1^{\alpha} k_2^{\mu} k_3^{\nu} k_4^{\beta} $$

$$ - \epsilon_{\alpha\beta}^{\mu\nu} g^{\rho\sigma}(k_3 k_4) k_1^{\alpha} k_2^{\beta} - \epsilon_{\alpha\beta}^{\mu\rho} g^{\nu\sigma}(k_2 k_4) k_1^{\alpha} k_3^{\beta} - \epsilon_{\alpha\beta}^{\mu\sigma} g^{\rho\nu}(k_2 k_3) k_1^{\alpha} k_4^{\beta} $$

$$ - \epsilon_{\alpha\beta}^{\nu\rho} g^{\mu\sigma}(k_1 k_4) k_2^{\alpha} k_3^{\beta} - \epsilon_{\alpha\beta}^{\nu\sigma} g^{\mu\rho}(k_1 k_3) k_2^{\alpha} k_4^{\beta} - \epsilon_{\alpha\beta}^{\rho\sigma} g^{\mu\nu}(k_1 k_2) k_3^{\alpha} k_4^{\beta}). $$

(21)

Since the interaction contains an $\epsilon$ tensor, the coupling between two of the photons is different from that involving the other two, and the familiar pairwise equivalence of the six
terms does not hold. The matrix element satisfies gauge invariance (with respect to any of the four photons),

$$k_1^\mu M'_{\mu\nu\rho\sigma}(k_1, k_2, k_3, k_4) = 0, \quad \text{etc.}$$

(22)

We note that the above contribution to the matrix element is proportional to the fourth power of the inverse dyon mass, $M'_{\mu\nu\rho\sigma} \propto M^{-4}$. However, this result is only valid at low energies, where the photon momenta are small compared to $M$, being obtained from an effective, non-renormalizable theory.

Thus, as a result of interference between two one-loop diagrams corresponding to loops with dyons and those with simply electrically charged particles there is an asymmetry between the processes of photon splitting and photon coalescence [4]. The physical effect of this asymmetry will depend on the photon spectrum and the directions of the photon momenta with respect to the magnetic field. In particular, the asymmetry vanishes when these are perpendicular, i.e. for $\cos \theta = 0$. Furthermore, the asymmetry is linear in the product of the dyon charges, and proportional to the fourth power of the electron to dyon mass ratio.

**IV. ELECTRIC DIPOLE MOMENT**

The contribution of this matrix element (21) breaks the $P$ and $T$ invariance of ordinary electrodynamics. Thus, among the sixth-order radiative corrections to the electron-photon vertex there are terms containing this photon-photon scattering subdiagram with a dyon loop contribution (see Fig. 1), that induce an electric dipole moment of the electron [8].

Indeed, one can write the contribution of this diagram to the electron-photon vertex as:

$$\Lambda_\mu(p', p) = \frac{e^2}{(2\pi)^8} \int d^4k_1d^4k_3 \frac{1}{k_1^2 + i\epsilon} \frac{1}{k_2^2 + i\epsilon} \frac{1}{k_3^2 + i\epsilon} M_{\alpha\beta\gamma\mu}(k_1, k_2, k_3, k)$$

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2This has been noted by I.B. Khriplovich [21] — see also a recent paper by Flambaum and Murray [7].

3Of course, there are more diagrams.
where \( M_{\alpha\beta\gamma\mu}(k_1, k_2, k_3, k) \) is the polarization pseudotensor representing the dyon box diagram contribution to the photon-photon scattering amplitude, the low-energy limit of which is given by the pseudotensor \( M'_{\alpha\beta\gamma\mu} \) of Eq. (21).

In order to extract the electric dipole moment from the general expression (23) it is convenient, according to the approach by [22] to exploit the identity

\[
M'_{\alpha\beta\gamma\mu}(k_1, k_2, k_3, k) = -k^\nu \frac{\partial}{\partial k^\mu} M_{\alpha\beta\gamma\nu}(k_1, k_2, k_3, k),
\]

which can be obtained upon differentiating the gauge invariance condition of the polarization tensor [cf. Eq. (22)] with respect to \( k^\mu \).

Substituting (24) into (23) we can write the \( ee\gamma \) matrix element as

\[
M_{ee\gamma}(p', p, k) = e^\mu(k) \bar{\psi}(p') \Lambda_{\mu
u}(p', p) u(p) = e^\mu(k) k^\nu \bar{\psi}(p') \Lambda_{\mu
u}(p', p) u(p)
\]

where \( e^\mu(k) \) is the photon polarization vector and

\[
\Lambda_{\mu\nu}(p', p) = -\frac{e^2}{(2\pi)^8} \int d^4k_1 d^4k_3 \frac{1}{k_1^2 + i\epsilon} \frac{1}{k_2^2 + i\epsilon} \frac{1}{k_3^2 + i\epsilon} \frac{\partial}{\partial k^\mu} M_{\alpha\beta\gamma\nu}(k_1, k_2, k_3, k) \times \gamma^\alpha \frac{\hat{p}'}{\hat{p} + k_1 + m} \gamma^\beta \frac{\hat{p} - k_3 + m}{(p - k_3)^2 - m^2 + i\epsilon \gamma}. \tag{26}
\]

Since the matrix element (25) is already proportional to the external photon momentum \( k \), one can put \( k = 0 \) in \( \Lambda_{\mu\nu} \) after differentiation to obtain the static electric dipole moment.

Then, following [22], we note that due to Lorentz covariance of \( \Lambda_{\mu\nu} \), it can be written in the form

\[
\Lambda_{\mu\nu}(p', p) = \left( \tilde{A} g_{\mu\nu} + \tilde{B} \sigma_{\mu\nu} + \tilde{C} P_\mu \gamma_\nu + \tilde{D} P_\nu \gamma_\mu + \tilde{E} P_\mu P_\nu \right) \gamma_5 + \ldots \tag{27}
\]

where we have omitted terms that do not violate parity, as well as those proportional to \( k_\mu \), and where \( \sigma_{\mu\nu} = (\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu)/2 \), and \( P_\mu = p_\mu + p'_\mu \).

Substituting this expression into the matrix element \( M_{ee\gamma}(p', p, k) \) of Eq. (25) one can see that there are two contributions to the \( P \) violating part, arising from the \( \tilde{B} \) and \( \tilde{C} \) terms.
In order to project out the dipole moment from (25), one has to compare Eq. (27) with the phenomenological expression for the electric dipole moment \( d_e \) [23]:

\[
M_{ee\gamma}(p', p, k) = e^\mu(k)k'^\nu \bar{u}(p') \frac{d_e}{2m} \gamma_5 \sigma_{\mu\nu} u(p),
\]

(28)

In the non-relativistic limit it corresponds to the interaction Hamiltonian \(-\frac{d_e}{2m}\vec{\sigma}\vec{E}\).

Thus, multiplying (27) by \( \sigma_{\mu\nu}\gamma_5 \) and taking the trace we have:

\[
d_e = -\frac{m}{24} \text{Tr} [\sigma_{\mu\nu}\gamma_5 \Lambda^{\mu\nu}].
\]

(29)

In order to provide an estimate of the induced electric dipole moment we need to estimate \( \Lambda^{\mu\nu} \). The first task is to evaluate the polarization pseudotensor \( M_{\alpha\beta\gamma\mu} \) corresponding to the virtual dyon one-loop subdiagram. If we were to substitute for \( M_{\alpha\beta\gamma\mu} \) the low-energy form \( M'_{\alpha\beta\gamma\mu} \) of Eq. (21) into Eq. (23), we would obtain a quadratically divergent integral.

On the other hand, straightforward application of the Feynman rules in QED with magnetic charge (see, e.g., [16]) would give for the photon-by-photon scattering subdiagram in Fig. 1:

\[
M_{\alpha\beta\gamma\mu}(k_1, k_2, k_3, k) = \frac{Qg^3}{2\pi^4} \int d^4q \text{ Tr} \left( \Gamma_{\alpha} \frac{1}{\not{q} + k_1 - M} \Gamma_{\beta} \frac{1}{\not{q} - k_3 - k - M} \right)
\times \Gamma_\gamma \frac{1}{\not{q} - k - M} \gamma_\mu \frac{1}{\not{q} - M}.
\]

(30)

Here \( \Gamma_\alpha \) represents the magnetic coupling of the photon to the dyon, which we take according to ref. [16] to be

\[
\Gamma_\mu = -i\varepsilon_{\mu\nu\rho\sigma} \frac{\gamma^\nu k^\rho n^\sigma}{(n \cdot k)}.
\]

(31)

The vertex function depends on \( k^\rho \), the photon momentum entering the vertex, and on \( n^\sigma \), a unit constant space-like vector corresponding to the Dirac singularity line. It was shown

\[\text{It should be noted that the expression (21) contains contributions from such loop diagrams with all possible combinations of either three or one magnetic-coupling vertex } \Gamma_\rho.\]
by Zwanziger [24] that although the matrix element depends on $n$, the cross section as well as other physical quantities are $n$ independent.

Calculations using this technique are very complicated and can only be done in a few simple situations [16], for example, in the case of the charge-monopole scattering problem [25]. We will here avoid this approach.

While the integration over $q$ in Eq. (30) is logarithmically divergent (the magnetic couplings in (30) are dimensionless), after renormalization the sum of such contributions must in the low-energy limit reduce to the form given in Eq. (21). We also note that the substitution of (30) into Eq. (26) yields a convergent integral. Thus, the following method for evaluating $\Lambda_{\mu\nu}$ suggests itself. We divide the region of integration into two domains: (i) the momenta $k_1$ and $k_3$ are small compared to $M$, and (ii) the momenta are of order $M$ (or larger).

In the first region, the form (21) can be used, but since the integral is quadratically divergent, the integral will be proportional to $M^2$. Together with the over-all factor $M^{-4}$ this will give a contribution $\propto M^{-2}$. For large values of the photon momenta, the other form, Eq. (30) can be used. This gives a convergent integral, and dimensional arguments determine the scale to be $M^{-2}$. It means that

$$|\Lambda_{\mu\nu}| \sim \frac{e^2 Q g (Q^2 - g^2)}{(4\pi^2)^3 M^2}.$$  (32)

The numerical coefficient has been estimated as $1/(4\pi^2)^3$, one factor $1/4\pi^2$ from each loop, and the $1/24$ of Eq. (29) is assumed cancelled by a combinatorial factor from the number of diagrams involved. This is of course a very rough assessment.

Now we can estimate the order of magnitude of the electron dipole moment generated by virtual dyons. It is obvious from Eqs. (29) and (32), that in order of magnitude one can write

$$d_e \sim \frac{e^2 Q g (Q^2 - g^2) \ m \ m}{(4\pi^2)^3 M^2}.$$  (33)

This estimate can be used to obtain a new bound on the dyon mass. Indeed, recent experimental progress in the search for an electron electric dipole moment [27] gives a rather
strict upper limit:  \( d_e < 9 \cdot 10^{-28} e \) cm. If we suppose that \( Q \sim e \), then from (33) one can obtain \( M \geq 2 \cdot 10^6 m \approx 10^3 \) GeV. This estimate shows that the dyon mass belongs at least to the electroweak scale.

The above estimate coincides with the bound obtained by De Rújula [3] for monopoles, from an analysis of LEP data, but it is weaker than the result given in [7], where the limit \( M \geq 10^5 \) GeV was obtained. The authors of ref. [7] used the hypothesis that a radial magnetic field could be induced due to virtual dyon pairs. In order to estimate the effect, they used the well-known formula for the Ueling correction to the electrostatic potential, simply replacing the electron charge and mass with those of the monopole. But the Ueling term is just a correction to the scalar Coulomb potential due to vacuum polarization and cannot itself be considered as a source of a radial magnetic field. Indeed, there is only one second order term in the effective Lagrangian that can violate the \( P \) and \( T \) invariance of the theory, namely \( \Delta L' \propto E\overline{E} \). But in the framework of QED there is no reason to consider such a correction because it is just a total derivative. The reference to the \( \theta \)-term, used in [7] to estimate the electric charge of the dyon, is only relevant in the context of a non-trivial topology (e.g., in the 't Hooft-Polyakov monopole model) where their limit applies. In this case there are arguments in favour of stronger limits on the monopole (dyon) mass (see, e.g., [20]).

One should note that the dyon loop diagram considered above can also contribute to the neutron electric dipole moment. The experimental value \( d_n < 1.1 \cdot 10^{-25} e \) cm [28] will in the naive quark model with \( m \approx 10 \) MeV allow us to obtain an estimate of the dyon lower mass bound which is similar to the one obtained for the electron.

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FIG. 1. Typical three-loop vertex diagram. The closed line represents a dyon loop.