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Search for a heavy magnetic monopole at the Fermilab Tevatron and CERN LHC

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If a heavy Dirac monopole exists, the light to light scattering below the monopole production threshold is enhanced due to the strong coupling of monopoles to photons. This effect could be observable in the collision of virtual photons at proton colliders. At the Fermilab Tevatron it will be seen as pair production of photons with energies 200–400 GeV and roughly compensated transverse momenta 100–400 GeV/c. This effect could be seen at monopole masses of about 1–2.5 TeV at the upgraded Tevatron and 7.4–19 TeV at the CERN Large Hadron Collider depending on the monopole spin. [S0556-2821(98)50511-5]

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I. INTRODUCTION

The magnetic charge (monopole) was introduced into particle theory by Dirac [1] (see also [2]) to restore symmetry among electricity and magnetism. The idea of a monopole is very attractive in order to explain the mysterious quantization of the electric charge of particles. Therefore, the search for a monopole is essential despite the fact that there is no place for it in the standard description of our world.

The Dirac-Schwinger monopoles, discussed here, are *pointlike* particles. They differ strongly from nonlocal monopoles in the context of gauge theories, first considered by Polyakov and 't Hooft [3]. Below we assume that a few monopoles exist in the Universe and they are not yet observed due to their high mass.

Our basic idea is simple: *The existence of monopoles provides for a $\gamma\gamma$ elastic scattering at large angles that is sufficiently strong below the monopole production threshold. This effect is observable at colliders with energies smaller than the monopole masses.* Such a method for discovering monopoles has been suggested in Refs. [4],[5]. A similar idea for the process $e^+e^-\to Z\to 3\gamma$ has been considered recently [6,7] and tested at the CERN e^+e^- collider LEP [8]. Our paper deals with detailed calculations of this effect at hadron colliders. Throughout the paper we denote the monopole mass by M and its spin by J_M (we assume a definite spin of the monopole), ω is the characteristic photon energy [typically—the $\gamma\gamma$ c.m. system (c.m.s.) energy].

A theory with two point-like charges, electric and magnetic, cannot be standard QED.[1] According to Refs. [1], [2] (see also [6]), the electromagnetic field in such a theory is described by a vector potential having the Dirac string or some of its surrogate. To have unambiguous results, the elementary electric and magnetic charges e and g ought to be quantified so that[2]

$$g=\frac{2\pi n}{e},\quad n=\pm 1,\pm 2,\ldots \tag{1}$$

with $\alpha\equiv e^2/(4\pi)=1/137\Rightarrow\alpha_g\equiv g^2/(4\pi)=n^2/(4\alpha)$. Unfortunately, the explicit form of such a theory is still not yet known.

We are interested in the energy region below the monopole production threshold. Here the interactions of photons via virtual monopoles are considered as main effect. The effective expansion parameter is of the order of

$$g_{\rm eff}=\frac{g\omega}{\sqrt{4\pi}M}\equiv\frac{n\omega}{2\sqrt{\alpha}M}<1. \tag{2}$$

In this region general considerations like gauge invariance, threshold behavior, etc., together with a perturbative approach allow us to believe that a perturbation theory analogous to standard QED can be applied (certainly, in lowest nontrivial order only). In that case QED-like calculations should be valid, both in tree approximation and at one loop level. For one loop we can assume a Wick rotation into Euclidean region. As it is well known, the integration over the loop momentum p is convergent due to gauge invariance. Therefore, the integration region is limited by virtualities $p^2\lesssim\omega^2$, where the effective expansion parameter is small, and QED results are valid.

The estimate of the effective coupling constant is supported by estimating qualitatively the cross section of $\gamma\gamma\to\gamma\gamma$ scattering via a monopole loop (Fig. 1) in the $\gamma\gamma$ c.m.s. (without fixing the form of the interaction). Because of gauge invariance, each photon leg of the matrix element yields a factor ω. On dimensional grounds this factor has to be divided by the monopole mass M. Additionally, the mag-

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[1] According to [9], this theory even could violate $U(1)$ local gauge invariance of standard QED. We are thankful to Dr. He for clarification of this point.

[2] In the Schwinger theory [2] the quantity n should be even, $n=\pm 2,\pm 4,\ldots$.

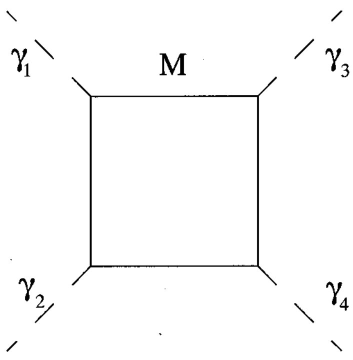

FIG. 1. $\gamma\gamma\to\gamma\gamma$ via monopole loop.

netic charge g has to be associated to each vertex. Therefore, the amplitude $\mathcal{M}\propto(g\omega/M)^4$ and the cross section is[3] $\sigma\propto(1/\omega^2)(g\omega/M)^8$.

The corresponding effective Lagrangian (of Heisenberg-Euler type) is

$$\mathcal{L}_{\rm eff}=-\frac{F_{\mu\nu}F^{\mu\nu}}{4}+\frac{g^4}{36(4\pi)^2M^4}\left(\frac{\beta_+ + \beta_-}{2}(F^{\mu\nu}F_{\mu\nu})^2 + \frac{\beta_+ - \beta_-}{2}(F^{\mu\nu}\tilde{F}_{\mu\nu})^2\right)+\cdots \tag{3}$$

with the electromagnetic field strength tensor $F_{\mu\nu}$ and $\tilde{F}^{\mu\nu}=\epsilon^{\mu\nu\alpha\beta}F_{\alpha\beta}/2$. An additional gauge fixing term has to be introduced to invert the photon propagator. The constants $\beta_\pm$ depend on the monopole spin (a numerical coefficient is introduced to simplify the final expressions). Based on arguments mentioned earlier (related to Wick rotation), we use here the coefficients $\beta_\pm$ obtained in QED. Their values have been found for different values of the monopole spin J_M in [10] ($J_M=1/2$), [11] ($J_M=0$), [12,13] ($J_M=1$). We use their combination $P=\beta_+^2+2\beta_-^2$.

[3]If some monopole—antimonopole bound state with mass$\ll 2M$ exists, it is wide enough (since $g\gg 1$), and it gives additional more strong light to light scattering. We do not speculate about this opportunity in detail.

Let us denote the standard Mandelstam variables for the $\gamma\gamma$ scattering by $\hat{s},\hat{t},\hat{u}$ with $\hat{s}+\hat{t}+\hat{u}=0$ [$\hat{s}$ is the effective mass of produced $\gamma\gamma$ system squared and $\hat{t}$ varies within the interval $(0,-\hat{s}/2)$]. Furthermore, θ is the scattering angle in c.m.s. of the photons with $0<\theta<\pi/2$. In our case $\hat{s}=4\omega^2$, $\hat{t}=-2\omega^2(1-\cos\theta)$.

The cross section for $\gamma\gamma$ scattering via a monopole loop at $\hat{s}\ll M^2$ is given by the expressions

$$\sigma^{\rm tot}=R\left(\frac{\hat{s}}{4}\right)^3, \quad R=\frac{28P}{405\pi}\left(\frac{n}{2\sqrt{\alpha}M}\right)^8, \tag{4}$$

$$d\sigma=\frac{5(3+\cos^2\theta)^2}{56}\sigma^{\rm tot}d\cos\theta\equiv\frac{5}{7}\left(\frac{\hat{s}^2+\hat{t}^2+\hat{u}^2}{\hat{s}^2}\right)^2\sigma^{\rm tot}\frac{d\hat{t}}{\hat{s}}, \tag{5}$$

$$P=\beta_+^2+2\beta_-^2=\begin{cases}0.085, & J_M=0,\\ 1.39, & J_M=1/2,\\ 159, & J_M=1.\end{cases} \tag{6}$$

Note that the result strongly depends on the monopole spin J_M.

These expressions are correct up to contributions of order $\mathcal{O}(g_{\rm eff}^2)$. We expect that with growing $g_{\rm eff}$, the increase of the cross section as a function of energy stops and the production of a larger (even) number of photons becomes essential. The effective parameter here is expected to be less than $g_{\rm eff}^2$. For example, in QED with spin 1/2 fermions, the ratio of coefficients of the 6γ to 4γ operators in the effective Lagrangian for $\gamma\gamma$ interactions is about $0.3g_{\rm eff}^2$ [10]. Besides, one can expect that below the threshold ($\omega<M$) the expressions (4) correctly describe the sum of cross sections of processes $\gamma\gamma\to 2\gamma$, 4γ, 6γ, ... with increasing multiple photon production at higher $g_{\rm eff}$.

II. $\gamma\gamma$ PRODUCTION VIA A MONOPOLE LOOP AT HADRON COLLIDERS

The cross section for the production of two photons via a virtual heavy monopole loop in pp or $p\bar{p}$ collisions is a convolution of the $\gamma\gamma$ cross section (5) with the photon fluxes arising from the colliding protons

$$\sigma_{pp\to\gamma\gamma X}=\left(\frac{\alpha}{\pi}\right)^2\int n_1(\omega_1,Q_1^2)n_2(\omega_2,Q_2^2)\frac{d\omega_1}{\omega_1}\frac{d\omega_2}{\omega_2}\frac{dQ_1^2}{Q_1^2}\frac{dQ_2^2}{Q_2^2}d\sigma_{\gamma\gamma\to\gamma\gamma}. \tag{7}$$

Since in the dominant integration region $Q_i^2\ll\hat{s}$, the transverse motion of initial photons can be neglected with good accuracy.

The relevant scale for the virtuality dependence of the $\gamma\gamma\to\gamma\gamma$ subprocess cross section is given by the unique inner parameter of monopole loop—the monopole mass. Therefore, this dependence appears in the form of the quantity Q^2/M^2. Since $Q^2\ll\hat{s}\ll M^2$, it is safely neglected below.

Taking these approximations into account, the cross section (7) can be written in factorized form (valid for $Q_i^2\ll\hat{s}$). The photon flux densities n_i depend on their ''own'' photon variables only

$$n_i\equiv n_i\left(y_i=\frac{\omega_i}{E},Q_i^2\right), \quad \hat{s}=4\omega_1\omega_2 \tag{8}$$

and we have for the cross section

$$\sigma_{pp\to\gamma\gamma X}=\left(\frac{\alpha}{\pi}\right)^2\int\frac{dy_1}{y_1}f(y_1)\frac{dy_2}{y_2}f(y_2)d\sigma_{\gamma\gamma\to\gamma\gamma}$$
$$\equiv\left(\frac{\alpha}{\pi}\right)^2 RE^6N^2(E), \tag{9}$$

$$N(E)=\int y^2f(y)dy,\quad f(y)=\int n(y,Q^2)\frac{dQ^2}{Q^2}. \tag{10}$$

The photon flux density arising from one proton is a sum over elastic and inelastic contributions:

$$n(y,Q^2)=(D_{\rm el}+D_{\rm in})+\frac{y^2}{2}(C_{\rm el}+C_{\rm in}). \tag{11}$$

In the individual contributions the quantity Q^2 is limited kinematically from below (m_p is the proton mass):

$$Q^2>Q^2_{\rm min}=(M_X^2-m_p^2)\frac{y}{1-y}+m_p^2\frac{y^2}{1-y}. \tag{12}$$

M_X is the effective mass of the system produced in the virtual γ^*p collision, $M_X=m_p$ in the elastic case.

The elastic contribution is written via standard proton form factors G_E and G_M:

$$C_{\rm el}(Q^2)=G_M^2(Q^2),$$
$$D_{\rm el}(Q^2)=(1-y)\frac{4m_p^2G_E^2(Q^2)+Q^2G_M^2(Q^2)}{4m_p^2+Q^2}\left(1-\frac{Q^2_{\rm min}}{Q^2}\right). \tag{13}$$

The integral over Q^2 in Eq. (10) rapidly converges at the upper boundary due to the form factors. The integral is saturated at Q^2 values given by the form factor scale. This scale is much lower than the other parameters of the problem. Therefore, the upper integration limit can be extended to infinity and the elastic contribution to $N(E)$ becomes energy independent.

The inelastic contribution is written via the proton structure functions F_1, F_2. $C_{\rm in}$ and $D_{\rm in}$ are integrals over the squared effective mass of the produced system M_X^2.

We basically start from Eq. (D.4) and Table 8 of Ref. [14]:

$$C_{\rm in}(Q^2)=\frac{2}{Q^2}\int dM_X^2F_1(M_X^2,Q^2),$$
$$D_{\rm in}(Q^2)=\frac{(1-y)}{Q^2}\int dM_X^2xF_2(M_X^2,Q^2)\left(1-\frac{Q^2_{\rm min}}{Q^2}\right). \tag{14}$$

In order to use the standard representations of the structure functions, we change the integration variable from M_X^2 to the Bjorken variable x using the relation $M_X^2=m_p^2+Q^2(1-x)/x$. Next, from inequality (12) the lower limit in x is found,

$$x_{\rm min}=\frac{y}{1-y^2m_p^2/Q^2}, \tag{15}$$

yielding the useful relation

$$(1-y)x\left(1-\frac{Q^2_{\rm min}}{Q^2}\right)=y\left(\frac{x}{x_{\rm min}}-1\right).$$

The upper x value is reached at the minimal value for the quantity $L^2=M_X^2-m_p^2\approx 2m_pm_\pi$:

$$x_{\rm max}=\frac{Q^2}{Q^2+L^2}. \tag{16}$$

With these transformations we present $C_{\rm in}(Q^2)$ and $D_{\rm in}(Q^2)$ in the form

$$C_{\rm in}(Q^2)=2\int_{x_{\rm min}}^{x_{\rm max}}\frac{dx}{x^2}F_1(x,Q^2),$$
$$D_{\rm in}(Q^2)=y\int_{x_{\rm min}}^{x_{\rm max}}\frac{dx}{x^2}\left(\frac{x}{x_{\rm min}}-1\right)F_2(x,Q^2). \tag{17}$$

Concerning the upper Q^2 limit for the inelastic contribution, one should keep in mind that the basic representation with photon flux factorization is valid at $Q_i^2\ll\hat{s}$ only. Near this bound the original integrand is less than that obtained using our factorization, but its contribution to the total cross section is small. Therefore, the integration region can be restricted from above by $Q_i^2\ll\hat{s}=y_1y_2s$ with sufficient accuracy. Since the virtual photon energies in both fluxes are roughly equal, we use

$$Q^2_{i\,\rm max}\approx y_i^2s \tag{18}$$

as upper limit for the inelastic contribution without further reducing the accuracy. With this choice the factorization remains valid.

In the y integration of Eq. (10), $y_{\rm max}=1$ can be used as upper bound. The inaccuracy in the quantities that enter (structure functions) is inessential, since contributions at $y_i\approx 1$ can be safely neglected.

III. TOTAL CROSS SECTION AND MASS LIMITS

Elastic contribution. The proton form factors are written in the standard dipole approximation

$$\frac{G_M(Q^2)}{\mu_p}=G_E(Q^2)=\frac{1}{(1+Q^2/Q_0^2)^2},$$
$$\mu_p=2.79285,\qquad Q_0^2=0.71\ {\rm GeV}^2. \tag{19}$$

Using dimensionless variables $a=4m_p^2/Q_0^2$ and $z=(a/4)\times y^2/(1-y)$, the energy distribution is found[4] as follows:

$$\frac{dN_{\rm el}}{dy}=\frac{y^2(1-y)}{a}\Bigg((a+z(1+\mu_p^2+4a))I(z,0)$$
$$+(a+z)(\mu_p^2-1)I(z,a)-\frac{a}{(1+z)^3}\Bigg), \tag{20}$$

where

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[4]Compare [14], Eq. (D.7) with a sign misprint corrected and upper Q^2 limit neglected.

$$I(z,a) = -\frac{1}{(1-a)^4}\left(\log\frac{a+z}{1+z} + \sum_{k=1}^{3}\frac{1}{k}\left(\frac{1-a}{1+z}\right)^k\right).$$

With the numerical values for μ_p and Q_0, the remaining y integration leads to the energy independent constant $N_{\rm el} = 0.017672$ as elastic contribution to $N(E)$.

Inelastic contribution. The structure functions $F_{1,2}$ are parametrized using the next to leading order quark distributions of [15] and the parton model relation $F_2 = 2xF_1$. To test the sensitivity of our results on this particular parametrization, we have repeated some of the calculations using the structure functions of [16]. The difference is small enough.

The Q^2 dependence of the $F_{1,2}$ parametrizations is valid above some low input value $Q^2_{\rm low}$. Fortunately, the inelastic contributions to $N(E)$ should be almost insensitive to the behavior of structure functions at small Q^2. To test this statement, we consider two extrapolations for these functions below $Q^2_{\rm low}$: $F_{1,2}(x,Q^2) = 0$ or $F_{1,2}(x,Q^2) = F_{1,2}(x,Q^2_{\rm low})$. At $E = 0.9$ TeV in Glück-Reya-Vogt (GRV) parametrization [15] the results coincide within 1% accuracy.

Numerical estimates. At 0.9 TeV we obtain $N(E) = 0.0410$ in the GRV and $N(E) = 0.0338$ in the Martin-Roberts-Ryskin-Stirling (MRRS) approximation [16] for this factor of Eq. (9). The quantity $N(E)$ depends only weakly on the proton energy E. The ratio $N(E)/N(0.9\ \text{TeV})$ varies from 0.966 at $E = 0.5$ TeV to 1.006 at 1 TeV and 1.102 at $E = 7$ TeV using GRV parametrization.

For a monopole with $J_M = 1/2$ and $M/n = 1$ TeV and a proton energy of 1 TeV, we obtain the total cross section

$$\sigma_{pp\to\gamma\gamma X} = 150\ \text{fb}. \tag{21}$$

It is useful to rewrite this photon production cross section for different proton collider energies and different kinds of monopoles:

$$\sigma_{pp\to\gamma\gamma X}(E,M,P,n) = 108P\left(\frac{nE}{M}\right)^8\left(\frac{N(E)}{N(1\ \text{TeV})}\right)^2 \times\left(\frac{1\ \text{TeV}}{E}\right)^2\ \text{fb}. \tag{22}$$

Let us consider a luminosity integral of 2 fb^{-1} and a beam energy of 0.9 TeV (Tevatron). If we assume 10 events to be sufficient to detect the discussed effect, the following mass limits can be reached (for different spins J_M):

$$M < n \otimes \begin{cases} 0.998\ \text{TeV}, & J_M = 0, \\ 1.42\ \text{TeV}, & J_M = 1/2, \\ 2.56\ \text{TeV}, & J_M = 1. \end{cases} \tag{23}$$

Taking 100 fb^{-1} and $E = 7$ TeV (LHC) we obtain the following mass limits:

$$M < n \otimes \begin{cases} 7.40\ \text{TeV}, & J_M = 0, \\ 10.5\ \text{TeV}, & J_M = 1/2, \\ 19.0\ \text{TeV}, & J_M = 1. \end{cases} \tag{24}$$

The obtained limiting quantities correspond to a $\gamma\gamma \to \gamma\gamma$ subprocess cross section [calculated as $R\langle\omega\rangle^6$ with an estimate of $\langle\omega\rangle$ taken from Eq. (26) below] which is roughly 500 pb for the Fermilab Tevatron and 10 pb for the CERN Large Hadron Collider independent on monopole spin and n. These values are much higher than the cross section of the main background process $\gamma\gamma\to\gamma\gamma$ via a W boson and a t quark loop which is about 20–30 fb [13].

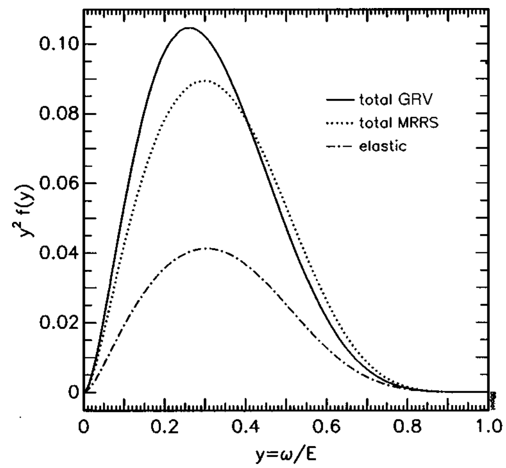

FIG. 2. Energy distribution of virtual photons from the proton, $y^2 f(y)$ at 0.9 TeV.

IV. ENERGY AND MOMENTUM DISTRIBUTIONS

Energy distribution for virtual photons. The photon fluxes in Eq. (9) decrease with increasing photon energies. On the other hand, the $\gamma\gamma$ subprocess cross section rapidly increases with $\hat{s} = 4\omega_1\omega_2$. Therefore, the main contribution to the cross section is given by region of intermediate ω_i. As already mentioned, the dependence of the subprocess cross section on Q^2 can be neglected since the characteristic values of virtuality $Q^2 \ll \hat{s} \ll M^2$. Therefore, the energy distribution for virtual photons is given by [compare Eqs. (9) and (10)]

$$\frac{d^2\sigma}{dy_1 dy_2} = \left(\frac{\alpha}{\pi}\right)^2 RE^6 y_1^2 f(y_1) y_2^2 f(y_2). \tag{25}$$

Figure 2 shows the energy distribution $y^2 f(y)$ for photons arising from one proton (the distribution is identical for each photon) using the two structure function parametrizations. The virtual photon energy distribution varies only weakly with the energy E, this weak energy dependence manifests itself in the weak E dependence of $N(E)$ mentioned above. At 0.9 TeV the average energy of the colliding photons and their energy spread are

$$\langle\omega\rangle = 0.314E, \qquad \langle\Delta\omega\rangle = 0.149E. \tag{26}$$

$\gamma\gamma$ differential cross sections. As noticed at the beginning of Sec. II, the transverse motion of incident (virtual) photons can be neglected with reasonable accuracy. Therefore, the transverse momenta of the produced photons are balanced: $p_{T3} \approx -p_{T4} \equiv p_T$. The 4-momenta of these photons in the c.m.s. of the protons can be written in two forms:

$$p_{3,4} = (\varepsilon_{3,4}, \pm p_T, 0, p_{L3,L4}) \equiv p_T(\cosh\eta_{3,4}, \pm 1, 0, \sinh\eta_{3,4}). \tag{27}$$

Here transverse and longitudinal momenta and rapidities of photons are introduced. Using these notations we have

$$\sin\theta = \frac{p_T}{\sqrt{\omega_1\omega_2}}, \quad y_{1,2} = \frac{\omega_{1,2}}{E} = \frac{p_T}{2E}(e^{\pm\eta_3} + e^{\pm\eta_4}). \tag{28}$$

With the standard transformation

$$\frac{\partial}{\partial y_1 \partial y_2 \partial \cos\theta} \equiv \frac{\varepsilon_3\varepsilon_4 \partial}{p_T \partial p_{L3} \partial p_{L4} \partial p_T}$$

the integrand of Eq. (7) (after integrating over Q_i^2) can be considered as the distribution over the momenta of produced photons. Then the differential cross section of $\gamma\gamma$ production via monopole loop can be presented in the form

$$\begin{aligned} \varepsilon_3\varepsilon_4 \frac{d^6\sigma}{d^3p_3 d^3p_4} &= \left(\frac{\alpha}{\pi}\right)^2 y_1^2 f(y_1) y_2^2 f(y_2) \frac{5RE^4}{112\pi} \\ &\quad \times \delta^{(2)}(\mathbf{p}_{T3} + \mathbf{p}_{T4})\Phi, \\ \Phi &= \left(4 - \frac{p_T^2}{\omega_1\omega_2}\right)^2 \\ &\equiv \left(4 - \frac{1}{\cosh^2[(\eta_3 - \eta_4)/2]}\right)^2. \end{aligned} \tag{29}$$

Integrating over one transverse momentum and azimuthal angle, the cross section is given by

$$\frac{d^3\sigma}{d\eta_3 d\eta_4 dp_T^2} = \left(\frac{\alpha}{\pi}\right)^2 y_1^2 f(y_1) y_2^2 f(y_2) \frac{5RE^4}{112}\Phi. \tag{30}$$

$\gamma\gamma$ total transverse momentum distribution. The total transverse momentum of the produced photon pair $\mathbf{k}_T \equiv \mathbf{p}_{T3} + \mathbf{p}_{T4}$ is equal to the sum of transverse momenta of virtual photons, $\mathbf{k}_T = \mathbf{q}_{T1} + \mathbf{q}_{T2}$. The latter are related to the photon virtualities by

$$\mathbf{q}_{Ti}^2 = (1 - y_i)(Q_i^2 - Q_{i\,\min}^2), \quad i = 1,2, \tag{31}$$

where $Q_{i\,\min}^2$ is the minimal value of Q_i^2 for energy fraction y_i as given in Eq. (12).

Since characteristically $Q_i^2 \ll \hat{s}$, the photon pair transverse momentum is typically much smaller than the transverse momenta of the produced photons: $k_T \ll p_{T3} \approx p_{T4}$. Therefore, the distribution in p_{Li}, p_T factorizes from that in k_T. The latter distribution can be obtained by integration over the virtual photon fluxes only (changing the order of integration). The detailed form of this dependence would be an additional test for the origin of the discussed photons. The corresponding calculations are simple but cumbersome, and one can postpone them to the time, when first events of discussed type are observed.

However, even before performing the calculation, we can conclude that this distribution is peaked near $k_T = 0$. This is evident for the elastic contribution, where the scale of the distribution is limited from above by that of the form factor. For the inelastic contribution the virtual photon flux distribution is wider, however the mean value of k_T^2 is much lower than $\hat{s}$.

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[1] P. A. M. Dirac, Proc. R. Soc. London **A133**, 60 (1931); Phys. Rev. **74**, 817 (1940).
[2] J. Schwinger, Phys. Rev. **144**, 1087 (1966).
[3] G. 't Hooft, Nucl. Phys. **B79**, 276 (1974); A. M. Polyakov, JETP Lett. **20**, 194 (1974).
[4] I. F. Ginzburg and S. L. Panfil, Sov. J. Nucl. Phys. **36**, 850 (1982).
[5] G. Belanger, F. Boudjema, and I. F. Ginzburg (unpublished).
[6] A. De Rujula, Nucl. Phys. **B435**, 257 (1995).
[7] M. Baillagreon, F. Boudjema, E. Chopin, and V. Lafage, Z. Phys. C **71**, 431 (1996).
[8] M. Acciari *et al.*, Phys. Lett. B **345**, 609 (1995).
[9] H.-J. He, Z. Qiu, and C. H. Tze, Z. Phys. C **65**, 175 (1995), and references therein.
[10] W. Heisenberg and H. Euler, Z. Phys. **38**, 714 (1936); V. Costantini, B. De Tollis, and G. Pistoni, Nuovo Cimento A **2**, 733 (1971); see also A. V. Berestetskii, E. M. Lifshits, and L. P. Pitaevskii, *Quantum Electrodynamics* (Pergamon, New York, 1982).
[11] J. Schwinger, Phys. Rev. **82**, 664 (1951).
[12] F. Boudjema, Phys. Lett. B **187**, 362 (1987); M. Baillargeon and F. Boudjema, *ibid.* **272**, 158 (1991); F.-X. Dong, X.-D. Jiang, and X.-J. Zhou, Phys. Rev. D **47**, 5169 (1993).
[13] G. Jikia and A. Tkabladze, Phys. Lett. B **323**, 453 (1994).
[14] V. M. Budnev, I. F. Ginzburg, G. V. Meledin, and V. G. Serbo, Phys. Rep., Phys. Lett. **15C**, 181 (1975).
[15] M. Glück, E. Reya, and A. Vogt, Z. Phys. C **67**, 433 (1995).
[16] A. D. Martin, R. G. Roberts, M. G. Ryskin, and W. J. Stirling, Eur. Phys. J. **C2**, 287 (1998).