Breaking up the Proton: An Affair with Dark Forces

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Deep inelastic scattering of $e^{\pm}$ off protons is sensitive to contributions from “dark photon” exchange. Using HERA data fit to HERA’s parton distribution functions, we obtain the model-independent bound $\epsilon \lesssim 0.02$ on the kinetic mixing between hypercharge and the dark photon for dark photon masses $\lesssim 10$ GeV. This slightly improves on the bound obtained from electroweak precision observables. For higher masses the limit weakens monotonically; $\epsilon \lesssim 1$ for a dark photon mass of 5 TeV. Utilizing PDF sum rules, we demonstrate that the effects of the dark photon cannot be (trivially) absorbed into re-fit PDFs, and in fact lead to non-DGLAP (Bjorken $x_B$-independent) scaling violations that could provide a smoking gun in data. The proposed $e^\pm p$ collider operating at $\sqrt{s} = 1.3$ TeV, LHeC, is anticipated to accumulate $10^3$ times the luminosity of HERA, providing substantial improvements in probing the effects of a dark photon: sensitivity to $\epsilon$ well below that probed by electroweak precision data is possible throughout virtually the entire dark photon mass range, as well as being able to probe to much higher dark photon masses, up to 100 TeV.

Introduction.

Are there new gauge interactions in Nature? A new, massive abelian vector boson (“dark photon”) can, at the renormalizable level, mix kinetically with the Standard Model hypercharge boson $\tilde{A}$:

$$\mathcal{L} \supset \frac{\epsilon}{2\cos\theta_W} F_{\mu\nu}^A B^{\mu\nu}. \quad (1)$$

Kinetic mixing with the hypercharge gauge boson becomes, after electroweak symmetry breaking, mixing of the dark photon with the neutral weak boson of the Standard Model (sm). We denote the unmixed dark photon by $A'_0$ and the unmixed neutral weak boson by $Z$. Diagonalizing the Lagrangian kinetic terms and gauge boson mass matrix results in three physical vectors that couple to SM fermions: the massless photon $\gamma$, and the mass eigenstates $Z$ and $A_D$.

Numerous searches for the dark photon have been undertaken by directly producing it, in which case the signature depends on its decay mode. In the minimal setup where the only relevant couplings come from Eq. (1), the dark photon decays back into charged standard model states, e.g. lepton pairs, offering striking signatures. However, the coupling in Eq. (1) may serve as our portal to a hidden sector that contains the particle species of the enigmatic dark matter [2] [3], and in this case the dark photon might decay invisibly or in a more complicated way depending on the structure of the hidden sector. It is therefore desirable to have “decay-agnostic” bounds that are independent of the details of a hidden sector.

In this study, we investigate one such decay-agnostic process: deep inelastic scattering (DIS) of $e^{\pm}$ off protons. As seen in the Feynman diagram in Fig. 1, DIS in the presence of kinetic mixing is mediated by the photon, the $Z$, and the dark photon $A_D$. $A_D$ exchange leads to distinct non-DGLAP scaling violations that may be constrained by existing data and may also be the smoking gun of a dark photon in future experiments.

Dark photon decay-agnostic limits on kinetic mixing were obtained in Refs. [4] [5] from electroweak precision observables (EWPO), driven mainly by the 0.1% precision $Z$ pole-mass measurements at LEP. The main effect is a shift of the $Z$ mass relative to $m_W/\cos\theta_W$, and using a global fit to EWPO, a bound of $\epsilon \lesssim 0.03$ was obtained for $m_{A_D} \ll m_Z$. We show that DIS measurements at the $e^\pm p$ collider HERA can improve on this bound. With a net luminosity of 1 fb$^{-1}$, HERA achieved 1% (systematics-limited) precision, however multiple measurements at

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FIG. 1. Deep inelastic scattering of $e^{\pm}$ on the proton, mediated by the Standard Model photon and $Z$ boson, and a dark photon arising from kinetic mixing with an abelian hidden sector. Measurements of this process at HERA and LHeC probe the mixing parameter and dark photon mass without relying on any assumptions about the production and decay properties of $A_D$. 

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Signals of kinetic mixing in deep inelastic scattering.

In this section we review the basics of deep inelastic scattering, incorporating dark photon exchange (for reviews, see e.g. [9][11]). DIS is described by the Lorentz-invariant kinematic variables

\[ Q^2 = -q^2 \, , \, x_B = \frac{Q^2}{2q \cdot p} \, , \, y = \frac{q \cdot k}{k \cdot p}, \]

where \( q \) is the momentum transfer and \( p \, (k) \) is the incoming proton’s (electron’s) momentum. Using these kinematic variables, the unpolarized neutral-current differential cross section rescaled as a (dimensionless) “reduced cross section” \( \sigma_{\text{red}}^{\text{NC}} \) is:

\[ \sigma_{\text{red}}^{\text{NC}} = \frac{Q^4 x_B}{2\pi \alpha^2 [1 + (1 - y)^2]} d^2 \sigma dQ_B^2. \]  

The cross section can be expressed in terms of the parton distribution functions (PDFs) as per the QCD factorization theorem. In the Quark Parton Model (QPM), DIS proceeds via elastic scattering on point-like quarks and anti-quarks, hence their PDFs \( f_q \) alone contribute and the variable \( x_B \) becomes the momentum fraction of the proton carried by the parton in the infinite momentum frame; moreover longitudinal effects are negligible. Neglecting parity-violating effects, \( \sigma_{\text{red}} \) is equal to the structure function:

\[ \bar{F}_2 = \sum_{i,j=\gamma,Z,A_D} \kappa_i \kappa_j F_{ij}^2, \]

where \( \kappa_i = Q^2/(Q^2 + M_{i}^2) \) accounts for the propagators of vector bosons of mass \( M_i \). At leading order in \( \alpha_s \),

\[ F_{ij}^2 = \sum_q x_B f_q (C_{i,j,e}^\nu + C_{i,j,e}^{\nu d}) (C_{i,j,q}^\nu + C_{i,j,q}^{\nu d}), \]

where the summation runs over \( q = u, d, c, \bar{c}, s, \bar{s}, b, \bar{b} \), and the vector and axial couplings to fermions (in units of \( e = \sqrt{4\pi\alpha} \)) are given as follows: For the SM photon,

\[ \{C_{\gamma,e}^\nu, C_{\gamma,u}^\nu, C_{\gamma,d}^\nu\} = \left\{-1, \frac{2}{3}, -\frac{1}{3}\right\}, \quad C_{\gamma}^\nu = 0. \]

For the unmixed \( Z' \) boson,

\[ \bar{C}_{Z'}^\nu \sin 2\theta_W = T_3^f - 2 q_f \sin^2 \theta_W, \quad \bar{C}_{Z'}^\nu \sin 2\theta_W = T_3^f, \]

where \( \{T_3^e, T_3^u, T_3^d\} = \{-1/2, 1/2, -1/2\} \) is the weak isospin, \( \{q_e, q_u, q_d\} = \{-1/2, 3/2, -1/3\} \) is the electric charge, and the usual Weinberg angle \( \sin^2 \theta_W \simeq 0.23127 \).

We can now add the effects of dark photon exchange. First we diagonalize the \( A' \) mixing with hypercharge shown in Eq. (1) through the field redefinition:

\[ B_{\mu} \rightarrow B_{\mu} + \frac{\epsilon}{\cos \theta_W} A'_{\mu}, \]

and canonically normalize the resulting dark photon kinetic term through the field rescaling:

\[ A'_{\mu} \rightarrow \frac{A'_{\mu}}{\sqrt{1 - \epsilon^2 / \cos^2 \theta_W}}. \]
In the $\epsilon \to \cos \theta_W$ limit the rescaling of $A'_\mu$ results in the enhancement of the dark gauge coupling, simultaneously enhancing couplings to SM fermion currents. This in turn increases our sensitivity to large $m_{A_D}$.

The dark photon and $Z$ squared mass matrix becomes

$$M^2 = \tilde{m}_Z^2 \begin{bmatrix} 1 & -\epsilon_W \\ -\epsilon_W & \epsilon_W^2 + \rho^2 \end{bmatrix},$$

where

$$\epsilon_W = \frac{\epsilon \tan \theta_W}{\sqrt{1 - \epsilon^2 \cos^2 \theta_W}},$$

$$\rho = \frac{\tilde{m}_{A'} \bar{m}_2}{\sqrt{1 - \epsilon^2 / \cos^2 \theta_W}}.$$  \hspace{1cm} (10)

and the $\bar{Z}$-$A'$ mixing angle is given by

$$\tan \alpha = \frac{1}{2\epsilon_W} \left[ 1 - \epsilon_W^2 - \rho^2 \\ -\text{sign}(1 - \rho^2) \sqrt{4\epsilon_W^2 + (1 - \epsilon_W^2 - \rho^2)^2} \right].$$  \hspace{1cm} (11)

The physical $Z$ couplings are

$$C_Z^\nu = (\cos \alpha - \epsilon_W \sin \alpha) C_Z^\nu + \epsilon_W \sin \alpha \cot \theta_W C_Z^\nu,$$

$$C_Z^\bar{\nu} = (\cos \alpha - \epsilon_W \sin \alpha) C_Z^\bar{\nu}$$ \hspace{1cm} (12)

while those of the physical $A_D$ are

$$C_{A_D}^\nu = -(\sin \alpha + \epsilon_W \cos \alpha) C_{A_D}^\nu + \epsilon_W \cos \alpha \cot \theta_W C_{A_D}^\nu,$$

$$C_{A_D}^\bar{\nu} = -(\sin \alpha + \epsilon_W \cos \alpha) C_{A_D}^\bar{\nu}.$$

The masses of the physical states are

$$m_{Z, A_D}^2 = \frac{\tilde{m}_Z^2}{2} \begin{bmatrix} 1 + \epsilon_W^2 + \rho^2 \\ \pm \text{sign}(1 - \rho^2) \sqrt{(1 + \epsilon_W^2 + \rho^2)^2 - 4\rho^2} \end{bmatrix}.$$  \hspace{1cm} (14)

Note that for fixed $\epsilon$ and any value of $\tilde{m}_{A'}/\tilde{m}_2$ the difference between the $Z$ and $A_D$ masses is always finite, $m_Z^2 - m_{A_D}^2 \geq 2 |\epsilon_W| m_Z^2$. This “eigenmass repulsion”, a well-known property of real symmetric matrices (including $M^2$) that implies that there are regions of the $\epsilon$-$m_{A_D}$ plane that cannot be realized.

Note that the cross section in Eq. (14) is invariant under $\epsilon \to -\epsilon$. This arises from requiring $A'$ to couple to both quark and lepton currents to be observable at DIS, so that deviations from the SM cross section arise first at $O(\epsilon^2)$.

For $Q^2 \ll m_Z^2$, the short-distance $Z$-exchange is negligible, and $A_D$ modifies the DIS cross section mainly through its constructive interference with the SM photon. Thus it effectively rescales the cross section by a factor of $[1 + \epsilon^2 Q^2/(Q^2 + m_{A_D}^2)]^2$ in this regime. We illustrate this in Fig. 2 where we plot, for a representative $x_B = 10^{-2}$, $\sigma_{NC}$ versus $Q$ for $\epsilon = 0$ as a band covering 2$\sigma$ uncertainties, and for $(\epsilon, m_{A_D}/\text{GeV}) = (0.1, 5), (0.3, 5)$, and $(0.3, 25)$. Clearly larger $\epsilon$ values produce larger effects; more subtly, $m_{A_D}$ sets the scale in $Q$ above which the effects of $A_D$ become significant. Note that the $\epsilon = 0.1$ curve lies well outside the SM band, indicating that HERA can probe $\epsilon \ll 0.1$ with a dataset spanning multiple $x_B$.

For $Q \gg m_Z$ DIS probes the regime of unbroken electroweak symmetry, where the SM process transpires effectively via massless $B$ exchange. As we will see in the next section, here too the effect of kinetic mixing is to rescale the cross section by $[1 + \epsilon^2 Q^2/(Q^2 + m_{A_D}^2)]^2$.

**HERA constraints and LHeC sensitivities.**

To constrain kinetic mixing with DIS, we use the combined datasets of Runs I and II at HERA [12] over the ranges

$$0.15 \leq Q^2/\text{GeV}^2 \leq 3 \times 10^4, \quad 5 \times 10^{-6} \leq x_B \leq 0.65.$$  \hspace{1cm} (15)

In principle, our constraints must be obtained fitting the HERA data simultaneously to both the dark photon parameters ($\epsilon, m_{A_D}$) and the PDFs $f_q$. In practice, however, we only fit to the dark photon parameters and use the HERAPDF2.0 LO PDF set derived in Ref. [12] (importing it into Mathematica via ManeParse2.0 [14]).

We believe the bounds we obtain from this simplified approach are, in fact, robust against performing a simultaneous fit. Let us justify this. First consider the region $m_{A_D} \ll Q_{\text{min}}$, where $Q_{\text{min}}$ is the smallest $Q$ probed at HERA. The cross sections in Eq. (14) are rescaled by $(1 + \epsilon^2)^2$ with respect to the SM, as discussed in the previous section. This implies that the PDFs could absorb this rescaling of the cross section by a simultaneous rescaling of all quark flavors, c.f. Eq. (14). However the normalization of $f_q$ is constrained by PDF sum rules. The quark-number sum rules

$$\int dx_B \left[ f_q(x_B) - f_{\bar{q}}(x_B) \right] = \begin{cases} 2, & q = u, \\ 1, & q = d, \\ 0, & q = s, c, b, \end{cases}$$

and the momentum sum rule

$$\int dx_B \left[ \sum_q f_q(x_B) + f_{\bar{q}}(x_B) \right] = 1,$$  \hspace{1cm} (15)

applied over the HERA $Q$ range are satisfied to $O(10^{-4})$ precision. Using additional data in ranges of $Q^2$ outside HERA’s, such as from beam dumps and hadron colliders, the sum rules can be further constrained once DGLAP evolution is accounted for.

\footnote{Our procedure can be viewed as estimating the expected sensitivities on our model using a second dataset with the same ($Q^2, x_B$) grid as HERA, with PDFs fitted under the Standard Model null hypothesis. This is precisely what we do when we estimate the sensitivity of LHeC below.}
In deriving our bounds, we use the \((Q^2, x_B)\) grid used for the HERA run involving \(e^+\) scattering with \(\sqrt{s} = 318\) GeV and 0.5 fb\(^{-1}\) luminosity, as given in Table 10 of Ref. \[12\]. This grid, containing 485 points, covers most of the \((Q^2, x_B)\) used in the other runs involving \(e^+\) scattering at smaller \(\sqrt{s}\) and luminosity; although data from all these runs were used for fitting PDFs, we do not use these other grids to avoid oversampling. We derive the 95\% c.l. limit by locating values of \((\epsilon, m_{A_D})\) for which

\[
\chi^2 = \sum_{\text{grid}} \frac{(\sigma_{\text{red}}^{\text{NC}} - \sigma_{\text{red}}^{\text{NC}}|\epsilon=0)^2}{(\delta \sigma_{\text{red}}^{\text{NC}})^2} = 5.99 ,
\]

where the summation is over the \((Q^2, x_B)\) grid mentioned above. The resulting limits are displayed in Fig. 3. We also show the decay-agnostic limits from the \((g - 2)_\mu\) measurement at the E821 experiment, requiring 5\(\sigma\) deviation from the central value using the calculations in Ref. \[3\], as well as the limits from \textsc{ewpo} derived in Ref. \[4\]. Our bounds are driven largely by about 25 data points in the \((Q^2, x_B)\) grid where the cross section is obtained with a maximum precision of 0.3\%-0.4\%. This is the origin of why our limits are (slightly) stronger than \textsc{ewpo} for \(m_{A_D} \lesssim 10\) GeV \(\ll m_Z\). In this regime the observable correction at both LEP Z pole measurements and DIS scales as \(c^2\), and while LEP operated at a precision of 0.1\%, which appears better than our precision, our bound actually benefits from 25 independent measurements, effectively diminishing our uncertainty by a statistical factor of \(\sqrt{25}\). We also display a hypothetical bound obtained by simply rescaling the SM cross sections by \(1 + 2c^2Q^2/(Q^2 + m_{A_D}^2)\), seen to trail the actual bound with amusing proximity. As discussed earlier, such a rescaling amounts to accounting only for \(A_D\) interference with the \(B\) exchange amplitude: for \(Q \ll m_Z\) this effectively rescales photon exchange, and for higher \(Q\) it rescales both photon and \(Z\) exchange.

In this figure we also show the 95\% c.l. sensitivity of the future \(e^+p\) collider \textsc{lhec} \[10\] derived by using the \((Q^2, x_B)\) grid for \(e^+\) scattering over the range

\[
5 \leq Q^2/\text{GeV}^2 \leq 10^6 , \quad 5 \times 10^{-6} \leq x_B \lesssim 0.8 .
\]

The \textsc{lhec} is anticipated to obtain \(10^3\) times the integrated luminosity of HERA, thus gaining in statistical precision by a factor of about 30. We are interested in characterizing the maximal sensitivity that \textsc{lhec} could achieve with this increased precision. This is a different objective from obtaining the best-fit PDFs across all datasets. Therefore, to estimate \textsc{lhec} sensitivity, we use \textsc{pdf4lhc15}\_\textsc{nlnlo}\_\textsc{lhec} PDFs fitted to pseudo-data Ref. \[17\], but then rescale the fractional uncertainties to match with \(1/10^3\) times the fractional uncertainties of \textsc{herapdf2.0} LO, optimistically assuming that systematic errors can be kept below this level. We have checked that rescaling the \(Q^2\) values of the \textsc{hera} grid by a factor of \(5/0.15 = 10^6/(3 \times 10^4)\), the envelopes of smallest uncertainties (as a function of \(Q^2\)) for either PDF set are well-aligned. We see from the figure that \textsc{lhec} exceeds

![FIG. 3. 95\% c.l. limits on the kinetic mixing parameter vs dark photon mass from deep inelastic scattering (DIS) measurements. Shown are limits from HERA derived using HERAPDF2.0 LO sets, and future sensitivities at \textsc{lhec}, using PDF4LHC15\_\textsc{nlnlo}\_\textsc{lhec} PDFs. For comparison are shown other decay-agnostic limits from measurements of electroweak precision observables (\textsc{ewpo}) and the muon \(g - 2\). Also shown are hypothetical limits obtained by rescaling SM DIS cross sections by a factor of \([1 + \epsilon^2Q^2/(Q^2 + m_{A_D}^2)]^2\), amounting to accounting only for interference between \(A_D\) and \(B\) boson exchange. In the gray-shaded region there is no physical value of \(m_{A_D}\) in the neighborhood of \(m_Z = 91.1876\) GeV due to repulsion of eigenmasses. The change in slope shift of the \textsc{hera} and \textsc{lhec} sensitivity curves at large \(\epsilon \gtrsim 0.7\) and \(m_{A_D} \gtrsim 1\) TeV occurs due to a factor of \(1/\sqrt{1 - \epsilon^2/\cos \theta_{W}}\) enhancement in the dark photon-fermion coupling. See text for further details.](image-url)
HERA in the entire $m_{A_D}$ range constrained by the latter, and indeed reaches $m_{A_D}$ up to 100 TeV thanks to probing the proton at very high Q.

The HERAPDF2.0 LO PDF set is designed to fit solely the HERA DIS data. We used this set not because we believe this is the best description of the quark PDFs, but because this is the most accurate interpolated description of the HERA data. Since we are interested in the sensitivity of HERA alone to kinetic mixing, we believe this is the correct approach to obtain the most accurate estimate of the sensitivity. We point out, however, that a more wide-ranging description of PDFs require “global fits” to DIS at HERA combined with beam-dump and hadron collider experiment datasets that include complementary ranges of $Q^2$ and $x_B$. Such PDF determinations contain additional sources of uncertainty [11]: (1) a “tolerance” factor to rescale the goodness-of-fit so that tensions in fitting multiple datasets may be eased to within $1\sigma$ uncertainty, (2) parameterization uncertainties introduced by the need to use numerous parameters to fit numerous datasets. The combination of these effects significantly increases the PDF uncertainties. Indeed we find that, had we used the global PDF set CT18Q [18], our bounds on $\epsilon$ would be weakened by a factor of up to 3. We do not believe this is a fair characterization of HERA’s bounds on kinetic mixing.

Finally, we note that EWPO sensitivities on $\epsilon$ are expected to improve by $\mathcal{O}(1)$ factors (a factor of $\sim 10$) with increased sensitivities provided by future LHC (ILC in GigaZ mode) measurements [5].

**DIS-cussion.**

Deep inelastic scattering of $e^\pm$ off protons is a sensitive, model-independent probe of kinetic mixing with a dark photon up to 100 TeV masses. No assumptions need to be made regarding the dark photon’s decay modes. We find HERA data is slightly more sensitive than EWPO for dark photon masses less than about 10 GeV. The LHeC could significantly improve the sensitivity of DIS to kinetic mixing, probing values of $\epsilon$ well below the sensitivity of EWPO data.

It is intriguing to consider the possibility of discovering a dark photon’s signature in DIS data. This seems quite unlikely with existing HERA data, since EWPO leads to a stronger constraint for most of the parameter space. The main constraint from EWPO arises due to a shift of $m_Z$ relative to $m_W/\cos\theta_W$. It is possible, though unlikely, that other physics in the dark sector could compensate for this apparent contribution to custodial violation and weaken the EWPO bounds. In addition, the parameter region $m_{A_D} \lesssim 10$ GeV where DIS is slightly more sensitive is strongly constrained by model-dependent searches, especially from $B$-factories. In particular, the BaBar collaboration has searched for dark photons produced via $e^+e^- \to A_D$ assuming that $A_D$ decays visibly through its kinetic mixing [19] or invisibly into a dark sector [20]. In both cases limits on $\epsilon$ at the level of $10^{-3}$ are obtained. A similar search by LHCb in the $\mu^+\mu^-$ final state constrains dark photons to $\epsilon \lesssim \mathcal{O}(10^{-3})$ up to 70 GeV masses [21]. These limits can potentially be weakened if $A_D$ couples to a dark sector with further structure, leading to more complicated final states as in, e.g., [22].

The LHeC’s sensitivity is significantly better than EWPO, and this provides the most exciting possibility to directly search for the non-DGLAP (Bjorken $x_B$-independent) scaling violation in the cross section illustrated in Fig. 2. Maximizing the sensitivity would be best optimized by simultaneously fitting the PDFs with dark photon exchange. Nevertheless, we have emphasized that PDF sum rules provide strong constraints on “fitting away” the effects of a dark photon on PDFs, and other studies [15] have also found that PDFs do not easily fit away the polynomial scaling exhibited by massive dark photons.

In this study we have focused on the effects of a dark photon on DIS, however this can also be extended to any new force between quarks and leptons, such as mediated by a gauged $U(1)_{B-L}$ vector boson, new scalars in the Higgs sector, or other exotic force carriers. We leave these investigations to future work.

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