New constraints on light vectors coupled to anomalous currents

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We derive new constraints on light vectors coupled to Standard Model (SM) fermions, when the corresponding SM current is broken by the chiral anomaly. Cancellation of the anomaly by heavy fermions results, in the low-energy theory, in Wess-Zumino type interactions between the new vector and the SM gauge bosons. These interactions are determined by the requirement that the heavy sector preserves the SM gauge groups, and lead to (energy/vector mass)\(^2\) enhanced rates for processes involving the longitudinal mode of the new vector. Taking the example of a vector coupled to baryon number, Z decays and flavour changing neutral current meson decays via the new vector can occur with (weak scale/vector mass)\(^2\) enhanced rates. These processes place significantly stronger coupling bounds than others considered in the literature, over a wide range of vector masses.

Introduction: Recent years have seen a resurgence of interest in the possibility of extending the Standard Model (SM) by including relatively light and very weakly coupled states \([1, 2]\). New light vectors are a popular candidate, having been proposed for purposes including addressing experimental anomalies at low energies \([3, 8]\), explaining puzzles such as baryon stability \([9]\), or acting as a mediator to a dark sector \([10–12]\).

In this paper, we will consider light vectors with dimension-4 couplings to SM states. Unless the SM current that a vector couples to is conserved (i.e. the electromagnetic (EM) or \(B − L\) current), there are (energy/vector mass)\(^2\) processes involving the longitudinal mode of the new vector. These make the SM + vector effective field theory (EFT) non-renormalisable, requiring a cutoff at some scale \(\propto (\text{vector mass}/\text{energy})\) at tree level, but violated by the chiral anomaly, which gives a divergence \([24]\)

\[
\partial^\mu j_\mu^\text{baryon} = \frac{A}{16\pi^2} \left( g^2 W^a_{\mu\nu} (\bar{W}^a)^{\mu\nu} - g^2 B_{\mu\nu} \tilde{B}^{\mu\nu} \right)
\]

where \(A = 3/2\), and \(\tilde{B}^{\mu\nu} \equiv \frac{1}{2}\epsilon^{\mu\nu\rho\sigma} B_{\rho\sigma}\) etc. If a new light vector \(X\) is coupled to the baryon current, then the SM + \(X\) EFT is non-renormalisable, and must be UV-completed at a scale \(\lesssim \frac{4\pi m_X}{g_X}\) \([25]\), where \(g_X\) and \(m_X\) are the coupling strength and mass of \(X\), respectively.

For models in which the SM current is broken by tree-level processes — e.g. axial currents are broken by fermion masses — such constraints have been considered in a number of works \([4, 13–15, 16]\). In this Letter, we point out they can also apply if a light vector \(X\) couples to a current that is conserved at tree level, but broken by the chiral anomaly (within the SM + \(X\) EFT), such as the SM baryon number current. These loop-level, but (energy/vector mass)\(^2\) enhanced, processes can place significantly stronger constraints on light \(X\) than existing constraints.

The only way to avoid such processes is for UV physics to introduce extra electroweak symmetry breaking into the low-energy theory, e.g. via heavy anomaly-cancelling fermions with electroweak breaking masses. This generally runs into strong experimental constraints. Conversely, cancelling the anomalies with new heavy fermions, that obtain their masses from a SM-singlet vacuum expectation value (VEV), always results in enhanced longitudinal \(X\) emission, as we show and exploit in the rest of this Letter.

Anomalous amplitudes: We will use the SM baryon number current as our prototypical example — a light vector coupled to this current has been considered in many papers over the past decades, e.g. \([3, 9, 17–23]\). Within the SM, the baryon number current is conserved at tree level, but violated by the chiral anomaly, which gives a divergence \([24]\)

\[
\partial^\mu j_\mu^\text{baryon} = \frac{A}{16\pi^2} \left( g^2 W^a_{\mu\nu} (\bar{W}^a)^{\mu\nu} - g^2 B_{\mu\nu} \tilde{B}^{\mu\nu} \right)
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In the simplest such UV completions, the anomalies are cancelled by introducing new fermions with chiral couplings to \(X\), and vectorial couplings to the SM gauge bosons. For example, the mixed anomalies can be cancelled with one weak doublet of Dirac fermions, with \((Y, X_A) = (−\frac{1}{2}, −3)\), and a weak singlet with \((Y, X_A) = (−1, 3)\), where \(Y\) and \(X_A\) are the hypercharge and axial \(X\) charge respectively \([24, 27]\). The \(XXX\) anomaly can then be cancelled by an additional SM-singlet fermion, and all of the new fermions can obtain heavy masses from a SM-singlet VEV.

Anomaly cancellation ensures that the theory is well-behaved at very high energies. However, as reviewed in \([16, 28]\), triangle diagram amplitudes have both a fermion-mass-independent ‘anomalous’ piece, and a piece that depends on the mass of the fermions in the loop.

1 In a companion paper \([16]\), we identify processes which place stronger constraints on vectors coupling to tree-level non-conserved SM currents.
The mass-dependent parts of longitudinal triangle amplitudes are proportional to the fermion’s axial coupling; since \( X \) has vectorial couplings to SM fermions, we obtain

\[-(p+q)_\mu \mathcal{M}^{\mu\nu\rho} = \frac{1}{2\pi^2} \epsilon^{\nu\rho\lambda\sigma} p_\lambda q_\sigma g_x g^2 \times \sum_f 2m_f^2 I_{00}(m_f, p, q) X_A f Y_f^2\]

\[\mathcal{M}^{\mu\nu\rho} \equiv \sum_{f,f_{SM}} X_\mu \Gamma_{\eta\xi} \Gamma_{\mu\nu\rho} f \to B_\nu, f \to B_\rho,\]

where \( f \) (\( f_{SM} \)) runs over heavy (SM) fermions; the ‘anomalous’ parts have cancelled between the new fermions and the SM fermions, and the mass-dependent ‘scalar integral’ term \( I_{00} \) is [28]

\[I_{00}(m_f, p, q) \equiv \int_0^1 dx \int_0^{1-x} dy \frac{1}{D(x, y, p, q)}, \]

\[D \equiv y(1-y)p^2 + x(1-x)q^2 + 2xy p \cdot q - m_f^2\]

For \( m_f^2 \gg p^2, q^2, p \cdot q \) we have \( I_{00} \approx -1/2m_f^2 \). Anomaly cancellation requires that \( 2 \sum_f Y_f^2 = \mathcal{A} \). Consequently, if the external momenta on the triangle are small compared to the masses of the new heavy fermions, then we have a total longitudinal amplitude of

\[-(p+q)_\mu \mathcal{M}^{\mu\nu\rho} \simeq \frac{1}{4\pi^2} \epsilon^{\nu\rho\lambda\sigma} p_\lambda q_\sigma g_x g^2 A\]

up to a relative error \( \mathcal{O}(\{p^2, q^2, p \cdot q \}/m_f^2) \).

This result is independent of the particulars of the heavy fermion sector, and is precisely the result we would have obtained by adopting a ‘covariant’ SM-gauge-group-preserving regularisation scheme within the SM + \( X \) EFT. As we review below, this is because the lack of extra electro-weak symmetry breaking (EWSB) in the UV theory determines the behaviour in the EFT.

The amplitudes for \( XWW \) triangles will have similar behaviour, with \( g^2 \) replaced by \( g \). An additional feature is that, since \( SU(2)_L \) is non-abelian, there are anomalous \( XWWW \) box diagrams. These have an analogous story of fermion mass dependence in the UV theory.

Low-energy theory and UV completions: Other classes of UV completions can give different results for low-energy triangle amplitudes. This is corresponds to the fact that the SM + \( X \) EFT generically includes dimension-4 Wess-Zumino (WZ) terms,

\[\mathcal{L} \supset C_B g_x g^2 \epsilon^{\mu\nu\rho\sigma} X_\mu B_\nu \partial_\rho B_\sigma + C_W g_x g^2 \epsilon^{\mu\nu\rho\sigma} X_\mu (W_\nu^a \partial_\rho W^a_\sigma + \frac{1}{3} g e^{abc} W^b_\nu W^c_\rho W^a_\sigma),\]

(5)

Since \( X \) has purely vectorial couplings to SM fermions, we must have \( C_B = -C_W \equiv C_{WZ} \) to avoid breaking the EM gauge symmetry. The coefficient of the WZ terms is determined by the UV theory (with the appropriate numerical value also determined by the regularisation scheme chosen for the anomalous diagrams [13]). For example, in a ‘SM-covariant’ regularisation scheme for the EFT, \( C_{WZ} = 0 \) corresponds to the UV theory introducing no extra EWSB, as per the example above. The key point is that there is no choice for \( C_{WZ} \) that preserves both \( U(1)_X \) and the SM gauge groups [25].

At the other extreme, the UV theory may preserve \( U(1)_X \) by breaking the SM gauge group — in the EFT, this corresponds to the WZ term cancelling longitudinal \( X \) amplitudes from SM fermion triangles. For example, we could cancel the anomalies by introducing new, heavy SM-chiral fermions, which obtain their masses through large Yukawa couplings with the SM Higgs. Once the new fermions are integrated out, this introduces extra EWSB into the low-energy theory [29] [30], analogously to integrating out the top quark in the SM (after which the photon remains massless, even though the fermion content is anomalous). As reviewed in [16], this possibility is strongly constrained by electroweak precision tests and collider experiments. If the current LHC run sees no deviations from the SM, it would be fairly robustly ruled out. Variations on this scenario, employing an enlarged EWSB sector, may be slightly more viable, but also inevitably introduce dangerous new physics at the electroweak scale.

Intermediate scenarios, in which the EFT breaks \( U(1)_X \) and the SM EW group, are also possible. If, in the UV theory, the SM-breaking contributions to the anomaly-cancelling heavy fermion masses are small compared to their total mass, then the WZ coefficient in the low-energy EFT will be approximately that expected from a SM-preserving theory, up to \( (m_{EWSB}/m_t)^2 \) corrections [16]. Conversely, if the new fermions obtain most of their mass from a EWSB-breaking VEV, there will be strong experimental constraints, analogous to those mentioned above for new SM-chiral fermions.

It should be noted that such constraints rely on the existence of new, SM-chiral states, which have effects (such as electroweak precision observables) unsuppressed by the small coupling \( g_x \). There may be more exotic UV completions, without anomaly-cancelling fermions, which are experimentally viable; within the low-energy theory, the effects of the SM-breaking WZ terms are all suppressed by \( g_x \), and if this is small enough, may not be problematic. While the rest of this Letter will work under the assumption of a SM-preserving low-energy EFT, this caveat should be kept in mind.

Another possible complication is that the new ‘UV’ degrees of freedom do not necessarily have to be heavier than all of the SM states. For example, if the anomalies are cancelled by SM-vector-like fermions, then collider constraints only require that they have masses \( \gtrsim 90 \text{GeV} \) [27] (for a baryon number vector). As per
Axion-like behaviour: By the usual Goldstone boson equivalence arguments, the $1/m_X$-enhanced parts of amplitudes involving longitudinal $X$ are approximately equal to those for the corresponding Goldstone (pseudo)scalar, $\varphi$. In our case, the processes which are not suppressed by $m_X$ all come from the anomalous couplings computed above. In the $\varphi$ theory, we can integrate by parts to write the interactions within the low-energy theory as

$$\frac{A}{16\pi^2} \frac{g^2 \varphi}{m_X} \left( g^2 (W^+ W^- - W^- W^+) + g g' (\cot \theta_w - \tan \theta_w) Z \bar{Z} + 2 g g' Z \bar{F} \right)$$

$$-i e g^2 F_{\mu\nu} (W^+_{\mu} W^-_{\nu} - W^+_{\nu} W^-_{\mu}) + \ldots$$

where we have suppressed indices, and the dots correspond to further terms of the form $AW^+ W^-$ and $ZW^+ W^-$.\

Since there is no two-photon anomalous coupling, longitudinal emission processes involving sub-EW-scale momenta are suppressed. Consequently, the relatively most important effects of the anomalous couplings arise either in high-energy collisions — for example, on-shell $Z$ decay at LEP — or in virtual processes which can be dominated by large loop momenta, such as rare meson decays.

$Z \rightarrow \gamma X$ decays: If $m_X < m_Z$, then the $\varphi Z \bar{F}$ coupling in (6) gives rise to $Z \rightarrow \gamma X$ decays, with width

$$\Gamma_{Z \rightarrow \gamma X} \simeq \frac{A^2}{384 \pi^5} \frac{g^2}{3} \frac{g'^2}{2} \frac{m_Z^3}{m_X}$$

(7)

corresponding to a branching ratio

$$\frac{\Gamma_{Z \rightarrow \gamma X}}{\Gamma_Z} \simeq 10^{-7} A^2 \left( \frac{\text{TeV}}{m_X/g_X} \right)^2$$

(8)

If $X$ decays invisibly, then LEP searches for single photons at half the $Z$ energy [22, 34] limit this branching ratio to be $\lesssim 10^{-6}$. The bounds for SM decays of $X$ are less stringent [34, 36], though the large number of $Z$ bosons produced at hadron colliders should allow enhanced sensitivity to rare $Z$ decays, as we discuss later.

**FCNCs:** The couplings of $X$ to quarks, and the anomalous $XWW$ coupling, lead to flavour changing neutral current (FCNC) interactions between quarks. These effects can be summarised by integrating out EW-scale states to obtain an effective interaction,

$$\mathcal{L} \supset g_{Xd,d_j} X_{\mu} d_\gamma^{\mu} P_L d_i + \text{h.c.} + \ldots$$

(9)

where we have taken a down-type FCNC for illustration, and have omitted other, higher-loop-order diagrams (as well as $X$ emission from external quark legs). The solid $XWW$ vertex indicates the sum of $WZ$ terms and fermion triangles (within a UV theory, it would simply be the sum over triangles). If $X$ is coupled to a fully-conserved current, then $g_{Xd,d_j} = 0$, and the effective interaction is higher-dimensional; if $X$ is coupled to a tree-level conserved current (as we consider here), then only the anomalous $XWW$ coupling contributes to $g_{Xd,d_j}$.

In the calculation of $g_{d,d_j,X}$, while each individual diagram is divergent, these divergences cancel in the sum over virtual quark generations, by CKM unitarity. As a result, the integral is dominated by momenta $\sim m_t$, and higher-dimensional couplings suppressed by the cut-off scale will give sub-leading contributions (in the UV theory, the masses of the UV fermions in triangles will be much larger than the external momenta of these triangles). The coefficient of the effective vertex is

$$g_{Xd,d_j} = -\frac{3 g^4 A}{(16 \pi^2)^2} g_X \sum_{\alpha \in \{u,c,t\}} V_{\alpha i} V_{\alpha j}^* F \left( \frac{m_2^2}{m_W^2} \right) + \ldots$$

$$\simeq -\frac{3 g^4 A}{(16 \pi^2)^2} g_X V_{ti} V_{tj}^* F \left( \frac{m_2^2}{m_W^2} \right) + \ldots,$$

(10)

where

$$F(x) \equiv \frac{x (1 + x (\log x - 1))}{(1 - x)^2} \simeq x \quad \text{(for } x \ll 1 \text{)}$$

(11)

Compared to these effective FCNC vertices, other effective flavour-changing operators are higher-dimensional, and so are suppressed by more powers of $g_X/m_X$ and/or $1/m_W$. Thus, despite equation (10) representing a 2-loop contribution (within the UV theory), it is able to dominate over 1-loop $d_i d_j X$ processes. For example, in the $B \rightarrow K X$ decay, we have

$$M^2_{\text{loop}}/M^4_{\text{1-loop}} \propto g^2/(16 \pi^2) \times (m_t/m_X)^2$$

(12)

which, for $m_X$ light enough to be emitted in the decay, is $\gg 1$. Competing SM FCNC processes are also

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2 the $WWWW$ terms from $W_{\mu\nu}^a (\bar{W}^a)^{\mu\nu}$ cancel, reflecting the lack of pentagon anomalies for an abelian vector [31]

3 The $\propto m_X^2$ (rather than $\propto m_X$) relative suppression of 1-
loop emission comes from angular momentum conservation in the pseudoscalar → pseudoscalar + vector decay; for $B \to K^*X$ decays, we would have $\mathcal{M}^{\text{loop}} / \mathcal{M}^{\text{tree}} \propto m_t^2/(m_X m_b)$ instead.

suppressed; for example, the $bs\gamma$ vertex is of the form $\propto m_b F_{\mu \nu} \bar{b} \gamma^\mu s_L$ (since the photon couples to a conserved current), while 4-fermion vertices are suppressed by at least $G_F$.

If $m_X$ is light enough, then FCNC meson decays via an on-shell longitudinal $X$ become possible, and are enhanced by $(\text{energy}/m_X)^2$, in addition to being lower-dimensional than other effective flavour-changing processes. Most directly, the $bsX$ and $sdX$ vertices result in $B \to K^{(*)}X$ and $K \to \pi X$ decays, giving new flavour-changing meson decays that can place strong constraints on the coupling of $X$. This is in exact analogy to the FCNC processes discussed in [18], for axion-like particles with a coupling to $W^a\tilde{W}^a$.

Experimental constraints: The left panel of Figure 1 shows a selection of experimental bounds on the coupling of a baryon number vector ($A = 3/2$); for consistency with other literature [21], we assume a loop-suppressed kinetic mixing with the photon, $\varepsilon = e g_X/(4\pi)^2$. As the figure illustrates, the anomalous bounds, derived in this work, are significantly stronger than existing bounds across a wide mass range. In particular, they constrain couplings significantly smaller than those at which we might expect the anomaly-cancelling fermions to have been observed at colliders [27], showing that our assumption of separation of scales can be valid. These improved bounds rule out some models of phenomenological interest. For example, [8] proposes a baryon number vector model to account for the claimed evidence of a new particle in $^8\text{Be}$ decays [19], with the anomalies being cancelled by heavy fermions that are vectorial under the SM. Their fiducial parameters of $m_X \approx 17\,\text{MeV}$, $g_X \approx 6 \times 10^{-4}$ (and a large kinetic mixing $\varepsilon \approx -10^{-3}$) result in $\text{Br}(B \to K^*X) \approx 2 \times 10^{-4}$ from the anomalous $XWW$ coupling, well above the experimental bound of $\Delta\text{Br}(B \to K^*\pi^-\pi^+) \lesssim 10^{-6}$ [41].

Figure 1 (right) shows the constraints that arise if $X$ has a significant branching ratio to invisible states (e.g. light dark matter, or additional neutrino species). For example, one light Dirac fermion $\chi$ with $X$-charge of 1 and $2 m_\chi < m_X$ will result in an invisible branching fraction of $\gtrsim 30\%$. The constraints from missing energy searches are strong, and for light $X$, limit $g_X^2/m_X^2$ to be below the Fermi coupling $G_F$. This disfavours such $X$ as a mediator between dark matter and the SM, since the dark matter annihilation cross section will, in the simplest models, be well below the $\sim \text{pb}$ level required for successful thermal freeze-out.

There will also be cosmological constraints from the effects of a thermal $X$ population in the early universe (e.g. on BBN), and astrophysical constraints from the production of $X$ in stars and supernovae. We leave the

FIG. 1. Left panel: Constraints on a vector $X$ coupling to baryon number, assuming a kinetic mixing with the SM photon $\epsilon \sim e g_X/(4\pi)^2$, and no additional invisible $X$ decay channels. Colored regions with solid borders indicate constraints from visible decays, dashed borders correspond to missing energy searches, and dotted borders denote projections based on current expected sensitivities. The gray regions indicate constraints from the previous literature. The new constraints come from searches for $K \to \pi X$ (green) [37,39], $B \to K X$ (blue) [40,43], $Z \to X\gamma$ (red) [32,36], and very displaced decays at the CHARM proton beam dump experiment [44]. For the latter, the enhanced $K \to \pi X$ decays result in larger $X$ production than computed in naive analyses [15,16]. The ‘anomalon’ line shows the approximate region in which anomaly-cancelling fermions would be light enough to have been detected [27]. The other gray constraints are from $\phi$ and $\eta$ decays [21], and $\Upsilon$ decays [9] (left to right). Right panel: As above, but with the assumption that $X$ dominantly decays invisibly.
calculation of such constraints to future work, but note that supernova constraints will apply for \( m_X \lesssim 100 \text{ MeV} \), while cosmological constraints will only apply at \( m_X \lesssim 10 \text{ MeV} \), or at significantly smaller couplings than shown in Figure 1 [10].

Future searches: At \( m_X \lesssim \text{GeV} \), the enhanced rate of \( K \to \pi X \) and \( B \to K X \) decays means that future proton beam dump experiments such as SHIP [45] will be more sensitive than projected in existing analyses. At higher masses, the enhanced \( B \to K X \) rate motivates searches for bumps in the invariant mass spectrum of \( B \to K + \) hadronic decays. For example, the \( B \to K \omega \) decay is detected as a peak in the \( m_{3\pi} \) distribution of \( B \to K \pi^+ \pi^- \pi^0 \) decays, with branching ratio error \( \sim 10^{-6} \) [50]; a similar search could be performed at other invariant masses.

For \( Z \to \gamma X \) decays, the large number of \( Zs \) produced at hadron colliders would likely allow leptonic \( Z \to \gamma(X \to l^+l^-) \) decays to be probed down to \( \mathcal{O}(10^{-5}) \) branching ratios or better [51, 52]. This would be especially helpful in constraining models of other anomalous vectors — for example, those with purely right-handed couplings [7], which result in \( XZ\gamma \) anomalous couplings, but no \( XWW \) coupling.

Conclusions: In this Letter, we have pointed out the phenomenological consequences of energy-enhanced longitudinal mode production, for light vectors coupling to anomalous SM currents. Such models have been considered for a variety of purposes in previous literature, but anomalous processes were overlooked. Taking the example of a light vector coupled to baryon number, we showed that anomalous production can place stronger coupling constraints over a wide mass range. We discuss these points in more depth in a companion paper [16], and also derive improved constraints on vectors coupled to SM currents that are broken by tree-level processes.

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