Massless Gauge Bosons other than the Photon

Bogdan A. Dobrescu
Fermi National Accelerator Laboratory, Batavia, Illinois 60510, USA
(Received 24 November 2004; published 22 April 2005)

Gauge bosons associated with unbroken gauge symmetries, under which all standard model fields are singlets, may interact with ordinary matter via higher-dimensional operators. A complete set of dimension-six operators involving a massless $U(1)$ field, γ', and standard model fields is presented. The $\mu \to e\gamma'$ decay, primordial nucleosynthesis, star cooling, and other phenomena set lower limits on the scale of chirality-flip operators in the 1–15 TeV range if the operators have coefficients given by the corresponding Yukawa couplings. Simple renormalizable models induce γ' interactions with leptons or quarks at two loops, and may provide a cold dark matter candidate.

DOI: 10.1103/PhysRevLett.94.151802 PACS numbers: 12.60.Cn, 13.35.Bv, 14.70.Pw

Introduction.—Experiments have established the existence of 12 gauge bosons. Although gauge invariance requires the gauge bosons to be massless, three of these ($W^\pm, Z$) are massive due to spontaneous symmetry breaking, while eight other (gluons) are confined in massive hadrons. The only massless spin-1 particle discovered so far is the photon.

This Letter addresses the possibility that massless gauge bosons other than the photon exist. If any standard model field is charged under a new *unbroken* gauge symmetry, the requirements of fermion mass generation and gauge anomaly cancellation [1] force the new symmetry to be $U(1)_{B-L}$, and the nonobservation of long-range forces other than electromagnetism and gravity imposes an extremely severe constraint on the new gauge coupling, $g_{B-L} \ll m_n/M_{\rm Pl} \approx 10^{-19}$, where m_n is the neutron mass, and $M_{\rm Pl}$ is the Planck scale (neutrino screening cannot relax the bound [2]).

It is possible, however, that standard model fields are neutral with respect to a gauge group, but nevertheless interact with the new massless gauge bosons. This follows from the existence of gauge invariant operators, of mass-dimension six or higher, which involve gauge field strengths. Such interactions have a natural decoupling limit: rather than forcing a dimensionless parameter to be extremely small, the experimental limits translate into a lower bound on the mass scale that suppresses the higher-dimensional operators.

In the case of an Abelian gauge symmetry, $U(1)_p$, under which all standard model fields are singlets, there is a single renormalizable term in the Lagrangian that involves both the $U(1)_p$ gauge boson and standard model fields: a kinetic mixing $c_0 P^{\mu\nu} B_{\mu\nu}$, where $B^{\mu\nu}$ is the $U(1)_Y$ (hypercharge) field strength, $P^{\mu\nu}$ is the $U(1)_p$ field strength, and c_0 is a dimensionless parameter. One might worry that this kinetic mixing would induce dimension-four couplings of the new gauge boson to all standard model fermions, but this is not the case [3]: the kinetic terms of the two $U(1)$ fields can be diagonalized and canonically normalized by an $SL(2, R)$ transformation, and their ensuing global $SO(2)$ symmetry allows the identification of the linear combination of $U(1)$ fields that couples to hypercharge as the new hypercharge gauge boson. The orthogonal combination, referred to as the "paraphoton" in Ref. [3] and denoted here by γ', does not have renormalizable couplings to standard model fields.

Effective interactions.—The leading interactions of the $U(1)_p$ field with standard model fields are given by dimension-six operators, so they are suppressed by two powers of the mass scale M where the operators are generated. The terms in the Lagrangian of this type which involve fermions are given by

$$\frac{1}{M^2} P_{\mu\nu}(\overline{q}_L \sigma^{\mu\nu} C_u \tilde{H} u_R + \overline{q}_L \sigma^{\mu\nu} C_d H d_R + \bar{l}_L \sigma^{\mu\nu} C_e H e_R + \text{H.c.}). \tag{1}$$

The notation is as follows: q_L, l_L are quark and lepton doublets, u_R, d_R are up- and down-type $SU(2)$-singlet quarks, e_R are electrically charged $SU(2)$-singlet leptons, and H is the Higgs doublet. An index labeling the three fermion generations is implicit. The 3×3 matrices in flavor space, C_u, C_d, C_e, have complex elements which are dimensionless parameters.

If "right-handed neutrinos" (gauge singlet fermions) are present, there is an additional dimension-six operator involving $P^{\mu\nu}$, as well as a lepton-number violating operator of dimension five: $P_{\mu\nu} \overline{N^c}_R \sigma^{\mu\nu} N_R$. Otherwise, the leading γ' interactions with neutrinos are given by a dimension-seven operator, $P_{\mu\nu} H \overline{l^c}_L \sigma^{\mu\nu} H l_L$.

Equation (1) displays a complete set of dimension-six operators involving $U(1)_p$ gauge fields and standard model fermions. For example, operators which involve the dual field strength, $\tilde{P}_{\mu\nu} \equiv \epsilon_{\mu\nu\lambda\tau} P^{\lambda\tau}$, can be reduced to Eq. (1) using the identities

$$\begin{aligned} \tilde{P}_{\mu\nu}\sigma^{\mu\nu} &= -2 P_{\mu\nu}\sigma^{\mu\nu}\gamma_5, \\ \tilde{P}_{\mu\nu}\gamma^\mu D^\nu &= P_{\mu\nu}(\sigma^{\mu\nu} \not{D} - 2i\gamma^\mu D^\nu)\gamma_5, \end{aligned} \tag{2}$$

and field equations such as $i\not{D}q_L = \lambda_u \tilde{H} u_R + \lambda_d H d_R$, where $\lambda_{u,d}$ are Yukawa coupling matrices, and the covariant derivatives D^ν contain only standard model gauge fields. Furthermore, using field equations and integration by parts, and ignoring operators of dimension higher than six, one can show that any chirality-preserving operator can be written in terms of operators (1). For example, the following identity is valid up to a total derivative:

$$iP_{\mu\nu}\overline{q}_L \gamma^\mu D^\nu q_L = \frac{-i}{4} P_{\mu\nu}\overline{q}_L \sigma^{\mu\nu}(\lambda_u \tilde{H} u_R + \lambda_d H d_R) + \text{H.c.} \tag{3}$$

In addition to the interactions with quarks and leptons shown in Eq. (1), the $U(1)_p$ field has purely bosonic interactions described by dimension-six operators:

$$\frac{1}{M^2} H^\dagger H (c_1 B_{\mu\nu} + \tilde{c}_1 \tilde{B}_{\mu\nu} + c_2 P_{\mu\nu} + \tilde{c}_2 \tilde{P}_{\mu\nu}) P^{\mu\nu}. \tag{4}$$

These renormalize the $U(1)_Y \times U(1)_p$ gauge couplings, and include vertices with paraphotons, Higgs bosons, photons, and Z bosons. The dimensionless parameters $c_{1,2}$, $\tilde{c}_{1,2}$ are real. Other operators vanish (e.g., $P_{\mu\nu} D^\mu H^\dagger D^\nu H$ and $P_{\mu\nu} G^\mu_\rho G^{\nu\rho}$, where $G^{\mu\nu}$ is the gluon field strength) or are total derivatives (e.g., $P_{\mu\nu} G^\mu_\rho \tilde{G}^{\nu\rho}$).

The dimensionless coefficients of the dimension-six operators can in principle have any value consistent with an effective theory description, namely, below $\sim 4\pi$. However, the operators (1) flip chirality, and although the origin of flavor structure in the standard model is not yet known, the elements of C_u, C_d, and C_e are expected to be of the order of or smaller than the corresponding Yukawa couplings.

Equation (1) is written in the weak eigenstate basis, where the standard model Yukawa couplings are flavor nondiagonal. In the mass-eigenstate basis, obtained by acting with different unitary 3×3 matrices on the left- and right-handed fields, the chirality-flipping operators change their flavor dependence: $C_f \to C'_f = U^f_L C_f U^{f\dagger}_R$, where U^f_L and U^f_R are the unitary matrices that diagonalize the masses of the $f = e, u, d$ fermions. Note that $U^{u\dagger}_L U^d_L$ is the Cabibbo-Kobayashi-Maskawa matrix, whereas the relation between the neutrino mixing angles and U^e_L is looser. The interactions of mass-eigenstate fermions, f', with the $U(1)_p$ field appear in the Lagrangian as follows:

$$\frac{v_h}{M^2} P_{\mu\nu} [\overline{f}' \sigma^{\mu\nu} (\text{Re} C'_f + i\,\text{Im} C'_f \gamma_5) f']. \tag{5}$$

These terms proportional to the vacuum expectation value of the Higgs doublet, $v_h \approx 174$ GeV, represent magnetic- and electriclike dipole moment operators.

Experimental limits.—There are various phenomenological constraints on the γ' interactions. First, the successful predictions of big-bang nucleosynthesis (BBN) limit the number Δg_* of effective relativistic degrees of freedom contributed by new particles that are in thermal equilibrium at a temperature $T_{\text{BBN}} \approx 1$ MeV. The maximum value of Δg_* is often expressed in terms of the maximum number of additional neutrino species allowed by the data: $\Delta g^{\max}_* = (7/4)\Delta N^{\max}_\nu$. The data on light element abundances exhibit some inconsistencies which translate into an uncertainty on $\Delta N^{\max}_\nu$. In Ref. [4], it is found that $\Delta N^{\max}_\nu = 0.6$ at the 2σ level, while in Ref. [5] the 2σ and 3σ contours in the N_ν versus baryon density plane extend to $\Delta N^{\max}_\nu \approx 0.2$ and 0.5, respectively. At any rate, the bound does not allow the 2 degrees of freedom of a paraphoton in thermal equilibrium. Thus, the paraphoton must decouple at $T_{\gamma'} > T_{\text{BBN}}$, so that [6]

$$\Delta g_*(T_{\text{BBN}}) = 2\left[\frac{g_*(T_{\text{BBN}})}{g_*(T_{\gamma'})}\right]^{4/3}. \tag{6}$$

Given that $g_*(T_{\text{BBN}}) = 43/4$, the lower bound on the number of relativistic degrees of freedom at freeze-out is $g_*(T_{\gamma'}) > 11.9 \times (\Delta N^{\max}_\nu)^{-3/4}$. In the standard model, $g_* = 247/4$ just above the temperature of the QCD phase transition, $T_{\text{QCD}} = 150$–180 MeV [7], and $g_* = 69/4$ just below T_{QCD}. Hence, $T_{\gamma'} = 180$ MeV is allowed as long as $\Delta N^{\max}_\nu > 0.11$. Note that the bound on $\Delta N^{\max}_\nu$ from a fit to the cosmic microwave background does not apply to the paraphoton, because at the time of recombination $g_* = 2$ rendering the effective number of γ' degrees of freedom negligible [smaller by a factor of $(43/8)^{4/3}$ than at T_{BBN}].

The interaction rate of the paraphoton with the thermal bath is given by

$$\Gamma(T) = \frac{2\zeta(3)}{\pi^2} T^3 \langle \sigma_{\gamma'} \rangle, \tag{7}$$

where $\zeta(3) \approx 1.202$, and $\langle \sigma_{\gamma'} \rangle$ is the thermally averaged cross section for γ' interactions with the standard model particles that are in thermal equilibrium. At temperatures just above T_{QCD}, the thermal bath includes light quarks (u, d, s), leptons (e, μ), photons, gluons, and neutrinos. The dominant γ' interactions involve the heaviest fermions, namely μ and s. A detailed study of the γ' decoupling would entail solving a set of coupled Boltzmann equations, but for the purpose of estimating the γ' couplings in terms of $T_{\gamma'}$ it suffices to impose that $\Gamma(T_{\gamma'})$ equals the expansion rate of the universe at freeze-out,

$$\Gamma(T_{\gamma'}) \approx \frac{T^2_{\gamma'}}{M_{\text{Pl}}} \left(\frac{2\pi^3}{45} g_*(T_{\gamma'})\right)^{1/2}. \tag{8}$$

The γ' annihilation due to interactions with muons proceeds through the processes shown in Fig. 1. The parametric dependence of the annihilation cross section is

$$\langle \sigma_{\gamma'} \rangle \sim \frac{\alpha c^2_\mu m^2_\mu}{M^4}, \tag{9}$$

where $c_\mu = |(C'_e)_{22}| v_h / m_\mu \lesssim O(1)$ sets the strength of the μ-γ' interaction, and α is the fine structure constant. Thus, the constraint on $T_{\gamma'}$ results in a limit on the effective mass scale of the μ-γ' interaction,

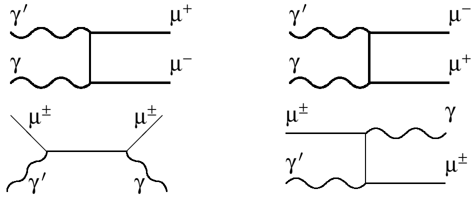

FIG. 1. Paraphoton annihilation via muon pair production ($\gamma'\gamma \to \mu^+\mu^-$) and Compton-like processes ($\gamma'\mu^\pm \to \gamma\mu^\pm$).

$$\frac{M}{\sqrt{c_\mu}} \approx 3.9\ \mathrm{TeV} \times [g_*(T_{\gamma'})]^{-1/8}\left(\frac{T_{\gamma'}}{1\ \mathrm{GeV}}\right)^{1/4}. \qquad (10)$$

The $T_{\gamma'}^{1/4}$ dependence dampens the sensitivity of the limit to the approximations used here. Using $g_*(T_{\gamma'}) = 255/4$ and $T_{\gamma'} \gtrsim 180$ MeV, Eq. (10) gives $M/\sqrt{c_\mu} \gtrsim 1.5$ TeV. The $\gamma'\gamma \to s\bar{s}$, $\gamma' s \to \gamma s$, and $\gamma'\bar{s} \to \gamma\bar{s}$ processes at $T_{\gamma'}$ impose a limit stronger by a factor of approximately $(\sqrt{3}m_s/m_\mu)^{1/2} \approx 1.2$ on $M/\sqrt{c_s}$, where $c_s \equiv |(C'_d)_{22}|v_h/m_s$.

Another phenomenon that can be affected by paraphotons is star cooling. One may use the studies of axions in order to derive the energy loss in stars due to γ' emission [8]. The axion is similar to the paraphoton: they are bosons with derivative couplings to fermions. The axion though has spin zero rather than one, so that the energy loss from stars is twice larger for γ' emission than for axion emission when the effective couplings are equal [9]. The energy loss due to electron-γ' interactions is proportional to the square of $g_{e\gamma'} = 4c_e m_e^2/M^2$, where $c_e \equiv |(C'_e)_{11}|v_h/m_e$. The limit on γ' emission through bremsstrahlung, such as $e^- + {}^4\mathrm{He} \to e^- + {}^4\mathrm{He} + \gamma'$, from the core of red giant stars [9] requires $g_{e\gamma'}^2/(4\pi) < 2.5 \times 10^{-27}$, so that $M/\sqrt{c_e} \gtrsim 3.2$ TeV. Compton-like scattering, $\gamma e^- \to \gamma' e^-$, in horizontal-branch stars sets a slightly weaker limit, $M/\sqrt{c_e} \gtrsim 1.8$ TeV.

The neutrino signal from supernova 1987A also limits the cooling through γ' emission, setting a bound on the coupling of the paraphoton to nucleons. The leading nucleon-γ' interactions are similar to Eq. (5):

$$\frac{v_h}{M^2} P_{\mu\nu}\overline{\mathcal{N}}\sigma^{\mu\nu}(\mathrm{Re}C_{\mathcal{N}} + i\,\mathrm{Im}C_{\mathcal{N}}\gamma_5)\mathcal{N}. \qquad (11)$$

The $C_{\mathcal{N}}$ form factors are of the order of $(C'_d)_{11}$ or $(C'_u)_{11}$. For example, QCD sum rules give $\mathrm{Im}C_{\mathcal{N}} \approx 0.2\,\mathrm{Im}(4C'_d - C'_u)_{11}$ for the neutron [10]. The effective nucleon-γ' coupling $g_{\mathcal{N}\gamma'}$ is proportional to the nucleon mass: $g_{\mathcal{N}\gamma'} \approx 4c_n m_d m_n/M^2$, where $c_n \equiv |C_{\mathcal{N}}|v_h/m_d \lesssim O(1)$. Requiring that the supernova was cooled predominantly by neutrinos implies $g_{\mathcal{N}\gamma'} \lesssim 2 \times 10^{-10}$, so that $M/\sqrt{c_n} \gtrsim 7$ TeV. There is also a range of larger nucleon-γ' couplings ($2 \times 10^{-7} \lesssim g_{\mathcal{N}\gamma'} \lesssim 7 \times 10^{-5}$) allowed by the supernova 1987A signal: the paraphotons are there too strongly coupled to nucleons to escape easily from the supernova, and too weakly coupled to produce a signal in the detectors that measured the neutrino signal. However, that range is not compatible with the nucleosynthesis constraint on quark-γ' couplings.

The long-range forces induced by paraphoton exchange between chunks of ordinary matter are feeble. The interactions of nonrelativistic electrons or nucleons with the paraphoton are spin dependent. Given that the average spin of macroscopic objects is small, the lower limits on $M/\sqrt{c_e}$ or $M/\sqrt{c_n}$ imposed by measurements of long-range forces are weaker than the TeV scale [11].

The magnetic- and electriclike dipole moments of the t quark could be probed at the Fermilab Tevatron, CERN Large Hadron Collider (LHC), or a linear e^+e^- collider if the dimensionless coefficient $|(C_u)_{33}|$ is of order unity and the mass scale M is of order 1 TeV.

The flavor off-diagonal interactions included in Eq. (5) are constrained by various flavor-changing neutral current processes. The most severe limit comes from the $\mu \to e\gamma'$ decay, and depends on $c_{e\mu} \equiv |(C'_e)_{12}|v_h/m_\mu$. A comparison of the decay width,

$$\Gamma(\mu \to e\gamma') = c_{e\mu}^2 \frac{m_\mu^5}{8\pi M^4}, \qquad (12)$$

times the measured muon lifetime, $3.3 \times 10^9\ \mathrm{eV}^{-1}$, to the experimental limit shown in Fig. 2 of Ref. [12], $\mathrm{Br}(\mu^+ \to e^+ X) < 3 \times 10^{-5}$ for any massless particle X, leads to $M/\sqrt{c_{e\mu}} \gtrsim 15$ TeV. Hadronic flavor-changing processes are also affected by γ' emission. Note though that the $K^+ \to \pi^+\gamma'$ decay is forbidden by angular momentum conservation, while the experimental constraints on other processes such as $\Sigma^+ \to p\gamma'$ or $B \to K^*\gamma'$, where γ' carries missing energy, are quite weak.

In the presence of kinetic mixing of the $U(1)_p \times U(1)_Y$ fields, the operators (1) contribute to the magnetic and electric dipole moments of the quarks and leptons. Hence, there are constraints on the products of the kinetic mixing parameter c_0 and the coefficients of operators (1). For example, the magnetic moment of the muon is shifted by $\sim 4c_0 c_\mu (m_\mu/M)^2/e$, and requiring that this is less than 10^{-9} gives $M/\sqrt{c_0 c_\mu} \gtrsim 12$ TeV. The lower limits on the mass scales that suppress the γ' interactions with various fermions are summarized in Table I.

The bosonic operators (4) lead to nonstandard Higgs boson decays, $h \to \gamma\gamma', Z\gamma', \gamma'\gamma'$, with γ' behaving in detectors as missing transverse energy. If M is not much higher than the electroweak scale and the parameters c_1

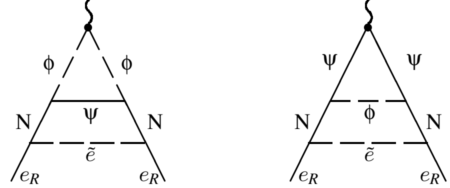

FIG. 2. Electron-γ' interaction in a renormalizable model.

TABLE I. Lower limits on the scales of various operators.

Mass scale	Limit (TeV)	Process
$M/\sqrt{c_e}$	3.2	bremsstrahlung in red giants
	1.8	Compton scattering in stars
$M/\sqrt{c_n}$	7	SN1987A cooling
	0.4	BBN
$M/\sqrt{c_{e\mu}}$	15	$\mu \to e\gamma'$
$M/\sqrt{c_\mu}$	1.5	BBN
$M/\sqrt{c_s}$	1.8	BBN
$M/\sqrt{c_0 c_\mu}$	12	$g_\mu - 2$

and $\tilde{c}_1$ are not much less than unity, a decay $h \to \gamma \not{E}_T$ could produce striking signals in collider experiments.

In the case of a non-Abelian gauge symmetry with respect to which all standard model fields are singlets, gauge invariance forces the operators describing interactions of the new gauge bosons to include two field strengths. As a result, the new gauge bosons can be produced only in groups of two or more, and most operators have mass dimension eight or higher. The only exceptions are the operators analogous to the last two terms in Eq. (4), which lead to invisible Higgs decays.

Renormalizable models.—The operators of Eq. (1) describe physics only at scales below $\sim M$. However, simple renormalizable models generate these operators after integrating out some heavy states. A generic feature of these models is the presence of fields charged under $U(1)_p$. The lightest particle of this type is stable and may be a viable dark matter candidate [6] provided it does not carry color or electric charge. It turns out though that even if that particle is a $U(1)_Y$ singlet, its electric charge cannot be exactly zero. This is because a kinetic mixing of the $U(1)_p \times U(1)_Y$ gauge bosons is induced at some loop level, and the $SL(2, R)$ and $SO(2)$ transformations that diagonalize the kinetic terms shift the electric charges of all fields carrying $U(1)_p$ charge [3]. A dark matter particle with mass of order 1 TeV must have an electric charge less than 10^{-4}, because otherwise it would have left an imprint on the cosmic microwave background [13] (other phenomena are discussed in Ref. [14]).

In models where the dimension-six operators arise at one loop, there is always a field charged under both $U(1)_p$ and $U(1)_Y$. A loop with this field in the internal line would induce a gauge kinetic mixing, which in turn would lead to an electric charge for the dark matter particle that may be larger than 10^{-4}. This problem is avoided in some models that generate the dimension-six operators at two loops: the loop-induced kinetic mixing is negligible if all fields charged under $U(1)_p$ have zero hypercharge, and the renormalizable tree-level kinetic mixing is absent if one of the $U(1)$s is embedded in a non-Abelian group.

As an example, consider a new scalar, $\tilde{e}$, with the same gauge charges as e_R, a gauge singlet Dirac fermion N, and a $\tilde{e}\overline{e}_R N$ coupling. If in addition there is a scalar ϕ and a fermion ψ, charged only under $U(1)_p$ and a $\phi\overline{\psi}N$ coupling, then e_R couples to the paraphoton at two loops, as shown in Fig. 2. A Higgs Yukawa coupling inserted on an e_R external line generates the last operator in Eq. (1). If the new particles have masses of order M, and the new gauge and Yukawa couplings are of order unity, then c_e, c_μ, $c_{e\mu}$ are given by a two-loop factor of order $(4\pi)^{-4}$. Consequently, the limit on M is rather loose: $M \gtrsim 100$ GeV (from Table I). The only $U(1)_p$-charged particles are ϕ and ψ, and the lightest of them is a cold dark matter candidate. Further studies are necessary to determine the region of parameter space where the dark matter halo does not collapse too fast due to γ' emission.

The quark-γ' moments are induced in this model only at three loops and are negligible. Similar models, with $\tilde{e}$ replaced by a scalar having the same charges as one of the quark fields, induce at two loops only the quark-γ' moments, and the tightest limit on M in that case is set by γ' emission from supernova 1987A: $M \gtrsim 40$ GeV for $c_n \sim (4\pi)^{-4}$. In the presence of an $H^\dagger H \phi^\dagger \phi$ coupling, the last two operators in Eq. (4) are induced at one loop, and the gauge kinetic mixing arises at two loops.

It is intriguing that massless gauge bosons other than the photon may interact with ordinary matter. The rather weak bound, below the electroweak scale in perturbative models, on the scale M that suppresses such interactions, makes it possible to search in collider experiments for the underlying dynamics that generate the dimension-six operators. It would also be interesting to investigate alternative experimental methods of searching for massless gauge bosons.

I am grateful to W. Bardeen, J. Beacom, G. Bertone, S. Chivukula, and I. Mocioiu for helpful comments. This work was supported by DOE under Contract No. DE-FG02-92ER-40704.

[1] T. Appelquist, B. A. Dobrescu, and A. R. Hopper, Phys. Rev. D **68**, 035012 (2003).
[2] S. I. Blinnikov, A. D. Dolgov, L. B. Okun, and M. B. Voloshin, Nucl. Phys. **B458**, 52 (1996).
[3] B. Holdom, Phys. Lett. **166B**, 196 (1986).
[4] A. Cuoco *et al.*, Int. J. Mod. Phys. A **19**, 4431 (2004).
[5] V. Barger *et al.*, Phys. Lett. B **566**, 8 (2003).
[6] S. Davidson, B. Campbell and D. C. Bailey, Phys. Rev. D **43**, 2314 (1991).
[7] Z. Fodor and S. D. Katz, J. High Energy Phys. 03 (2002) 014.
[8] S. Hoffmann, Phys. Lett. B **193**, 117 (1987).
[9] G. G. Raffelt, *Stars As Laboratories For Fundamental Physics* (University of Chicago, Chicago, 1996), Chap. 3.
[10] M. Pospelov and A. Ritz, Phys. Rev. D **63**, 073015 (2001).
[11] B. A. Dobrescu and I. Mocioiu (to be published).
[12] D. A. Bryman and E. T. H. Clifford, Phys. Rev. Lett. **57**, 2787 (1986).
[13] S. L. Dubovsky, D. S. Gorbunov, and G. I. Rubtsov, JETP Lett. **79**, 1 (2004).
[14] B. Holdom, Phys. Lett. B **178**, 65 (1986).