Letter

Dark photons from light with cosmic domain walls

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Light crossing dark domain walls that source a top form coupled to gauge Chern-Simons terms mixing visible and dark $U(1)$ gauge fields generically convert into dark photons. The effect is entirely localized on the wall and requires no additional ingredients. The conversion rate is a sharp function of the photon frequency in the wall rest frame, vanishing above the ultraviolet cutoff of the top form sector. Partial cloaking may also induce a rotation of the polarization of transmitted light of order $\Delta\vartheta \sim 10^{-3}$ radians, modify the cosmic microwave background power spectrum, and violate Etherington's reciprocity relation at low frequencies. These effects can impact cosmological determinations of the Hubble rate.

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We show that standard electromagnetism can couple to dark domain walls through a Chern-Simons interaction that mixes it with a dark 4-form (also known as top form) and a second $U(1)$ gauge field (a paraphoton). As a result, bubbles of evanescent dark energy [1], which mediate the discrete discharge of the top form flux, can induce both conversion of light into dark photons and detectable polarization rotation, accompanied by $O(1)$ variations of the dark energy density. A universe containing such dark walls would not be perfectly transparent: a fraction of photons with frequencies below the cutoff of the top form sector convert into paraphotons and are lost. We analyze the implications for localized cosmological sources and the cosmic microwave background (CMB), and outline possible consequences for determining the cosmic expansion rate.

We consider the following gauge- and Lorentz-invariant low-energy theory in the decoupling limit of gravity, $M_{\rm Pl} \to \infty$, in order to focus on gauge field dynamics [1–4]:

$$
\begin{aligned}
S \ni \int d^4x \Big\{ & \sum_n \left(-\frac{1}{4} (F^{(n)}_{\mu\nu})^2 - A^{(n)}_\mu J^{(n)\mu} \right) \\
& - \frac{\mathcal{H}}{4!\mathcal{M}^2} \sum_{n,m} \mathcal{Z}_{nm} \epsilon^{\mu\nu\lambda\sigma} F^{(n)}_{\mu\nu} F^{(m)}_{\lambda\sigma} \\
& - \frac{1}{2}\mathcal{H}^2 + \frac{1}{6}\epsilon^{\mu\nu\lambda\sigma}\partial_\mu \mathcal{H} \mathcal{B}_{\nu\lambda\sigma} \Big\} \\
& - \mathcal{T} \int d^3\xi \sqrt{\left| \det \left(\eta_{\mu\nu} \frac{\partial x^\mu}{\partial \xi^a} \frac{\partial x^\nu}{\partial \xi^b} \right) \right|} \\
& - \frac{\mathcal{Q}}{6} \int d^3\xi \mathcal{B}_{\mu\nu\lambda} \frac{\partial x^\mu}{\partial \xi^a} \frac{\partial x^\nu}{\partial \xi^b} \frac{\partial x^\lambda}{\partial \xi^c} \epsilon^{abc}. \qquad (1)
\end{aligned}
$$

Here $A^{(n)}_\mu$, $F^{(n)}_{\mu\nu} = \partial_\mu A^{(n)}_\nu - \partial_\nu A^{(n)}_\mu$, and $J^{(n)\mu}$ denote the visible ($n = 1$) and dark ($n = 2$) $U(1)$ vector potentials, field strengths, and conserved currents, respectively. The pseudoscalar $\mathcal{H}$ is the magnetic dual of the electric top form $U(1)$ field strength $\mathcal{G}_{\mu\nu\lambda\sigma} = 4\partial_{[\mu}\mathcal{B}_{\nu\lambda\sigma]}$, where $\mathcal{B}_{\nu\lambda\sigma}$ is its electric potential and brackets denote antisymmetrization. The top form is sourced by membranes of tension $\mathcal{T} \geq 0$ and charge $\mathcal{Q}$. The mass scale $\mathcal{M}$ sets the strong-coupling scale of the UV completion of the top form sector.

As discussed in [1–4], following [5–7], we take this UV completion to be a confining non-Abelian gauge theory that breaks chiral symmetry at a scale $\sim\mathcal{M}$, generating emergent top form dynamics [5]. Thus $\mathcal{M}$ is the cutoff of the low energy theory (1). Because the confining theory violates CP, the vacuum degeneracy is lifted by the discretely varying contribution $\mathcal{H}^2/2$, which is absent in the UV theory, because at energies above $\mathcal{M}$, the terms $\propto \mathcal{H}$ are absent. We note that the cutoff, and therefore both charges and tensions of the domain walls in the low energy theory are therefore set by a scale well below MeV, as discussed in detail in [1–4]. We will comment on this more below.

The higher-rank sector in Eq. (1) originates in a dark sector, as proposed in [1,4]. The visible and dark $U(1)$ vector fields could mix with each other and with the top form through anomaly terms involving fractionally charged dark matter, with both their θ parameters linked to the same dark sector top form [8–12]. The interaction is parametrized by a dimensionless matrix $\mathcal{Z}_{nm}$, which generalizes the coupling ζ used in [1–4]. Since $\epsilon^{\mu\nu\lambda\sigma} F^{(n)}_{\mu\nu} F^{(m)}_{\lambda\sigma} = \epsilon^{\mu\nu\lambda\sigma} F^{(m)}_{\mu\nu} F^{(n)}_{\lambda\sigma}$ and the operator is real, $\mathcal{Z}_{nm}$ is real and symmetric.

We could also introduce bulk $U(1)$ mixing through kinetic bilinears, gauge boson masses, or fractional dark matter charges, as originally considered in [13–16]. Such

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effects have been extensively studied and constrained by data [17]. For clarity we neglect them here and treat them as subleading corrections. Combining them with the present mechanism would be of interest in a more comprehensive analysis. The one point we stress is that the anomalous coupling strengths that enter $\mathcal{Z}_{nm}$ are to be understood as scale dependent, $\mathcal{Z}_{nm} \sim \mathcal{Z}_{nm}(0)/(1 + k^2/\mathcal{M}^2)$, since they arise due to integrating out the anomalous phases of the theory. This in turn cuts off the mixing with paraphotons in the UV, since the effective couplings are screened as $\mathcal{Z}_{nm} \sim \mathcal{Z}_{nm}(0)\mathcal{M}^2/k^2$.

We now analyze the theory (1). Away from the walls that source $\mathcal{H}$, the bulk dynamics reduces to conventional Maxwell electrodynamics in both the visible and dark sectors. Varying the action with respect to $A^{(n)}_\mu$ yields

$$\partial_\mu F^{(n)\mu\nu} = J^{(n)\nu} - \frac{1}{6\mathcal{M}^2}\epsilon^{\mu\nu\lambda\sigma}\sum_m \mathcal{Z}_{nm}\partial_\mu\left(\mathcal{H}F^{(m)}_{\lambda\sigma}\right). \quad (2)$$

Varying (1) with respect to $\mathcal{B}_{\nu\lambda\sigma}$ gives

$$n^\mu\partial_\mu\mathcal{H} = \mathcal{Q}\delta(r - r(t)), \quad (3)$$

where, as in [4], $r(t)$ denotes the wall trajectory, n^μ the outward-pointing normal, and r the coordinate along it. Away from the wall, $\mathcal{H}$ is constant and therefore $\frac{\mathcal{H}}{6\mathcal{M}^2}\epsilon^{\mu\nu\lambda\sigma}\partial_\mu F^{(n)}_{\lambda\sigma} = 0$ in the absence of visible and dark magnetic monopoles, which we assume. This follows from the Bianchi identities $\partial_{[\mu}F^{(n)}_{\lambda\sigma]} = 0$ in both sectors. Finally, varying Eq. (1) with respect to $\mathcal{H}$ and Hodge-dualizing yields

$$\mathcal{H}_{\mu\nu\lambda\sigma} = \mathcal{H}\epsilon_{\mu\nu\lambda\sigma} - \frac{1}{\mathcal{M}^2}\sum_{n,m}\mathcal{Z}_{nm}F^{(n)}_{[\mu\nu}F^{(m)}_{\lambda\sigma]}, \quad (4)$$

where $\mathcal{H}_{\mu\nu\lambda\sigma} = 4\partial_{[\mu}\mathcal{B}_{\nu\lambda\sigma]}$ is the spectator electric top form, whose flux is fixed by the CP-violating membrane sources. Thus, off the wall the theory consists of two decoupled $U(1)$ gauge theory sectors with charged matter and a derivative of the mixed Chern-Simons term feeding the spectator top form $\mathcal{H}_{\mu\nu\lambda\sigma}$.

The presence of walls qualitatively changes the situation. The jump of $\mathcal{H}$ across a wall by one unit of charge obstructs a global extension of the gauge field strengths $F^{(n)}_{\mu\nu}$ across it, as noted in [4]. The Chern-Simons mixing encoded by $\mathcal{Z}_{nm}$ modifies this structure, and complicates the matters slightly. However, this matrix can be diagonalized straightforwardly. Since $\mathcal{Z}_{nm}$ is real and symmetric, there exists a two-dimensional rotation matrix $\mathcal{R}(\theta)$ such that $\mathcal{R}^{\mathsf{T}}(\theta)\mathcal{Z}\mathcal{R}(\theta) = \mathcal{Z}_d$, where $\mathcal{Z}_d$ is diagonal with eigenvalues ζ_n, and T denotes transposition. Because the gauge kinetic terms are canonical, the kinetic matrix is unity in the $O(2)$ adjoint notation, and this rotation leaves it invariant.

We therefore change from the "interaction" basis $F^{(n)}_{\mu\nu}$ in (1) to the diagonal "propagation" basis $f^{(n)}_{\mu\nu}$ that crosses the wall. Using column vectors,

$$\begin{pmatrix} F^{(1)}_{\mu\nu} \\ F^{(2)}_{\mu\nu} \end{pmatrix} = \mathcal{R}(\theta)\begin{pmatrix} f^{(1)}_{\mu\nu} \\ f^{(2)}_{\mu\nu} \end{pmatrix}. \quad (5)$$

In terms of these fields, (1) becomes

$$\begin{aligned} S \ni &\int d^4x\Big\{\sum_n\Big(-\frac{1}{4}(f^{(n)}_{\mu\nu})^2 - a^{(n)}_\mu(\mathcal{R}^{\mathsf{T}}(\theta)J)^{(n)\mu} \\ &- \frac{\zeta_n\mathcal{H}}{4!\mathcal{M}^2}\epsilon^{\mu\nu\lambda\sigma}f^{(n)}_{\mu\nu}f^{(n)}_{\lambda\sigma}\Big) - \frac{1}{2}\mathcal{H}^2 + \frac{1}{6}\epsilon^{\mu\nu\lambda\sigma}\partial_\mu\mathcal{H}\mathcal{B}_{\nu\lambda\sigma}\Big\} \\ &- \mathcal{T}\int d^3\xi\sqrt{\left|\det\left(\eta_{\mu\nu}\frac{\partial x^\mu}{\partial\xi^a}\frac{\partial x^\nu}{\partial\xi^b}\right)\right|} \\ &- \frac{\mathcal{Q}}{6}\int d^3\xi\,\mathcal{B}_{\mu\nu\lambda}\frac{\partial x^\mu}{\partial\xi^a}\frac{\partial x^\nu}{\partial\xi^b}\frac{\partial x^\lambda}{\partial\xi^c}\epsilon^{abc}, \end{aligned} \quad (6)$$

which describes two independent copies of the wall-crossing $U(1)$ sector studied in [4]. They are mutually decoupled when the charge currents $J^{(n)}_\mu$ are switched off.

Let us consider the fate of light crossing the wall. Since $f^{(n)}_{\mu\nu}$ are linear combinations of $F^{(n)}_{\mu\nu}$, both propagation eigenstates are excited even if the incident wave is a single interaction eigenstate. Thus, an ordinary electromagnetic wave has two propagation channels across the wall, appearing as an admixture of $f^{(n)}_{\mu\nu}$ even when paraphotons are initially absent.

As in [4], we treat the walls as ultrathin and macroscopically large, approximating them as infinite planes. The wall-induced conversion in the propagation basis $f^{(n)}_{\mu\nu}$ is controlled by the small dimensionless parameters $\frac{\zeta_n\Delta\mathcal{H}}{\mathcal{M}^2} = \frac{\zeta_n\mathcal{Q}}{\mathcal{M}^2}$, which we treat perturbatively in the "sudden" approximation. In this respect, the wall behaves analogously to an interface between dielectrics [18].

Since the walls are infinite, we go to their rest frame at $z = 0$ and, using the boost invariance of the wall in the z–t plane [19,20], move to the frame where incident waves propagate along the normal. Away from the wall, the wave dispersion relation is $\omega^2 = \vec{k}^2$ on both sides, so waves strike the wall at right angles, propagating in vacuum before and after.

To determine wall-induced conversion of vector gauge fields, we define $e^{(n)}_i = f^{(n)}_{0i}$ and $b^{(n)i} = \frac{1}{2}\epsilon^{ijk}f^{(n)}_{jk}$. It is convenient to use the redefined field strengths as in [4], which, with $\sigma_n = \frac{\zeta_n\mathcal{H}}{6\mathcal{M}^2}$, are

$$\vec{\tilde{e}}^{(n)} = \vec{e}^{(n)} - \sigma_n\vec{b}^{(n)}, \qquad \vec{\tilde{b}}^{(n)} = \vec{b}^{(n)} + \sigma_n\vec{e}^{(n)}. \quad (7)$$

The Maxwell equations for $J_\mu^{(n)} = 0$ read

$$\vec{\nabla}\cdot\vec{\tilde{e}}^{(n)} = \vec{\nabla}\cdot(\sigma_n\vec{b}^{(n)}),$$
$$\vec{\nabla}\times\vec{\tilde{b}}^{(n)} - \partial_t\vec{\tilde{e}}^{(n)} + \vec{\nabla}\times(\sigma_n\vec{e}^{(n)}) + \partial_t(\sigma_n\vec{b}^{(n)}) = 0, \tag{8}$$

$$\vec{\nabla}\cdot\vec{\tilde{b}}^{(n)} = \vec{\nabla}\cdot(\sigma_n\vec{e}^{(n)}),$$
$$\vec{\nabla}\times\vec{\tilde{e}}^{(n)} + \partial_t\vec{\tilde{b}}^{(n)} + \vec{\nabla}\times(\sigma_n\vec{b}^{(n)}) - \partial_t(\sigma_n\vec{e}^{(n)}) = 0. \tag{9}$$

The first two equations (8) correspond to $\partial_\mu\tilde{f}^{(n)\mu\nu} = 0$, while (9) are the Bianchi identities $\partial_{[\mu}f^{(n)}_{\nu\lambda]} = 0$, after adding and subtracting $\vec{\nabla}\cdot(\sigma_n\vec{e}^{(n)})$, $\vec{\nabla}\times(\sigma_n\vec{b}^{(n)})$, and $\partial_t(\sigma_n\vec{e}^{(n)})$, and then using (7) to simplify the terms under derivatives.

For $\sigma_n \ll 1$, we treat wall interactions as a perturbation of the vacuum equations. In vacuum, $\vec{\nabla}\cdot\vec{e}^{(n)} = 0$ and $\vec{\nabla}\times\vec{b}^{(n)} = \partial_t\vec{e}^{(n)}$ to leading order. With $\partial_t\sigma_n = 0$ and $\vec{\nabla}\sigma_n \| \vec{n} \| \vec{k}$, since the propagating waves are right-handed triads $(\vec{k}, \vec{e}^{(n)}, \vec{b}^{(n)})$, Eqs. (8) and (9) reduce to

$$\vec{\nabla}\cdot\vec{\tilde{e}}^{(n)} = 0, \qquad \vec{\nabla}\times\vec{\tilde{b}}^{(n)} - \partial_t\vec{\tilde{e}}^{(n)} = -\vec{\nabla}\sigma_n\times\vec{e}^{(n)},$$
$$\vec{\nabla}\cdot\vec{\tilde{b}}^{(n)} = 0, \qquad \vec{\nabla}\times\vec{\tilde{e}}^{(n)} + \partial_t\vec{\tilde{b}}^{(n)} = -\vec{\nabla}\sigma_n\times\vec{b}^{(n)}. \tag{10}$$

With these conditions, integrating (10) over Gaussian pillboxes across the wall [4], the wall crossing to linear order in σ_n imposes continuity of the redefined fields

$$\vec{n}\cdot(\vec{\tilde{e}}^{(n)}_+ - \vec{\tilde{e}}^{(n)}_-) = 0, \qquad \vec{n}\times(\vec{\tilde{b}}^{(n)}_+ - \vec{\tilde{b}}^{(n)}_-) = 0,$$
$$\vec{n}\cdot(\vec{\tilde{b}}^{(n)}_+ - \vec{\tilde{b}}^{(n)}_-) = 0, \qquad \vec{n}\times(\vec{\tilde{e}}^{(n)}_+ - \vec{\tilde{e}}^{(n)}_-) = 0. \tag{11}$$

Thus, wall crossing induces an *electromagnetic duality* transformation of the monopole-free $U(1)$ propagation eigenfields [18,21]:

$$\begin{pmatrix}\vec{e}^{(n)}_-\\ \vec{b}^{(n)}_-\end{pmatrix} = \begin{pmatrix}1 & -\Delta\sigma_n\\ \Delta\sigma_n & 1\end{pmatrix}\begin{pmatrix}\vec{e}^{(n)}_+\\ \vec{b}^{(n)}_+\end{pmatrix}, \tag{12}$$

where

$$\Delta\sigma_n = \frac{\zeta_n}{6\mathcal{M}^2}\Delta\mathcal{H} = \frac{\zeta_n}{6\mathcal{M}^2}\mathcal{Q}, \tag{13}$$

is the jump across the wall due to a single unit of $\mathcal{H}$ flux discharge.

In the final step we rewrite this result in terms of the gauge fields in the "interaction" basis, by combining (5) and (12). Using the four-component column vector $(\vec{E}^{(1)}, \vec{B}^{(1)}, \vec{E}^{(2)}, \vec{B}^{(2)})^\mathsf{T}$, we obtain

$$\begin{pmatrix}\vec{E}^{(1)}_-\\ \vec{B}^{(1)}_-\\ \vec{E}^{(2)}_-\\ \vec{B}^{(2)}_-\end{pmatrix} = \mathcal{P}\begin{pmatrix}\mathcal{R}(\theta) & 0\\ 0 & \mathcal{R}(\theta)\end{pmatrix}\mathcal{P}\begin{pmatrix}\Sigma_1 & 0\\ 0 & \Sigma_2\end{pmatrix}$$
$$\times\mathcal{P}\begin{pmatrix}\mathcal{R}^\mathsf{T}(\theta) & 0\\ 0 & \mathcal{R}^\mathsf{T}(\theta)\end{pmatrix}\mathcal{P}\begin{pmatrix}\vec{E}^{(1)}_+\\ \vec{B}^{(1)}_+\\ \vec{E}^{(2)}_+\\ \vec{B}^{(2)}_+\end{pmatrix}, \tag{14}$$

where $\mathcal{P}$ is an involutory 4×4 permutation matrix, $\mathcal{R}$ the 2×2 rotation (5), and Σ_n the wall-borne duality matrix:

$$\mathcal{P} = \begin{pmatrix}1&0&0&0\\0&0&1&0\\0&1&0&0\\0&0&0&1\end{pmatrix}, \qquad \mathcal{R}(\theta) = \begin{pmatrix}\cos\theta & -\sin\theta\\ \sin\theta & \cos\theta\end{pmatrix},$$
$$\Sigma_n = \begin{pmatrix}1 & -\Delta\sigma_n\\ \Delta\sigma_n & 1\end{pmatrix}. \tag{15}$$

We now consider the implications when the incident wave is purely the standard electromagnetic wave from a distant source, crossing the wall en route to the observer. In this case, the initial dark components vanish, $\vec{E}^{(2)}_+ = \vec{B}^{(2)}_+ = 0$. Substituting into Eq. (14) and performing the matrix multiplications, we find

$$\vec{E}^{(1)}_- = \vec{E}^{(1)}_+ - \frac{\mathcal{Q}(\zeta_2\sin^2\theta + \zeta_1\cos^2\theta)}{6\mathcal{M}^2}\vec{B}^{(1)}_+, \qquad \vec{B}^{(1)}_- = \vec{B}^{(1)}_+ + \frac{\mathcal{Q}(\zeta_2\sin^2\theta + \zeta_1\cos^2\theta)}{6\mathcal{M}^2}\vec{E}^{(1)}_+,$$
$$\vec{E}^{(2)}_- = \frac{\mathcal{Q}\sin\theta\cos\theta(\zeta_2 - \zeta_1)}{6\mathcal{M}^2}\vec{B}^{(1)}_+, \qquad \vec{B}^{(2)}_- = -\frac{\mathcal{Q}\sin\theta\cos\theta(\zeta_2 - \zeta_1)}{6\mathcal{M}^2}\vec{E}^{(1)}_+. \tag{16}$$

The first line shows that the transmitted electromagnetic wave experiences polarization rotation by

$$\Delta\vartheta \simeq \frac{\mathcal{Q}(\zeta_2 \sin^2\theta + \zeta_1 \cos^2\theta)}{6\mathcal{M}^2} \ll 1, \tag{17}$$

generalizing [4] by including contributions from the dark $U(1)$ channel. Applying this to the CMB, polarization rotation analyses imply that $\Delta\vartheta$ is at most a degree of angle, or so [22].

The second line of Eq. (16) shows that a fraction of the visible wave converts into dark $U(1)$ during crossing. The fractional loss is estimated by comparing the transmitted dark intensity with the initial visible intensity. Since the waves are right-handed triads, their intensities are reliably approximated by our equations to quadratic order in σ_n. To this order, the transmitted dark power fraction is

$$\gamma \simeq \left(\frac{\mathcal{Q}}{6\mathcal{M}^2}\right)^2 \sin^2\theta \, \cos^2\theta (\zeta_2 - \zeta_1)^2. \tag{18}$$

We ignore the small apparent amplification of the visible wave, suggested by the formulas in (16), that indicate a slight increase of wave intensity upon crossing. This behavior is a known pathology in wave scattering, where the first-order perturbation fails by suggesting that the probability that an initial state scatters into itself exceeds unity. Because the wave fields in (16) form right-handed triads, this excess however is $\mathcal{O}(\sigma^4)$ and can be safely ignored. The issue is fully resolved in the Brillouin-Wigner approach to perturbation theory, by imposing unitarity, calculating off-diagonal transitions, and then subtracting those rates from unity.

From Eq. (18), if ζ_1 and ζ_2 have the same sign, visible-to-dark transitions are constrained by polarization bounds, $\gamma \sim \Delta\vartheta^2 \lesssim 10^{-6}$. An interesting possibility arises if ζ_1 and ζ_2 have opposite signs and $\tan\theta \simeq \pm\sqrt{|\zeta_1/\zeta_2|}$, suppressing $\Delta\vartheta$ while enhancing the conversion:

$$\gamma \simeq \left(\frac{\mathcal{Q}}{6\mathcal{M}^2}\right)^2 \zeta_2^2 \tan^2\theta \simeq \left(\frac{\mathcal{Q}}{6\mathcal{M}^2}\right)^2 |\zeta_1 \zeta_2|, \tag{19}$$

which can be appreciably larger than $\Delta\vartheta^2$. In this regime, sources of light behind the wall could appear significantly dimmed, potentially with $\Delta\vartheta \sim 10^{-3}$ and γ of order a few percent.

To estimate the visible power loss in wall crossing, Eq. (18), we ignored contributions from incident dark waves. This is reasonable for sources composed of the Standard Model particles, which are not prolific emitters of dark light. One might wonder about what happens with the CMB, since in other models [13–16] visible-dark mixing may induce distortions of the CMB spectrum, which can be constrained by observations [17].

As an illustrative example, consider photon-axion mixing [23] affecting Type Ia supernovae used to probe dark energy, reviewed in [24]. There, ultralight axionlike fields mix with light in the presence of large scale, randomly oriented cosmic magnetic fields. Because of the randomness, the mixing is not coherent but stochastic, leading to flavor equilibration. The effect is cumulative, frequency-dependent, and requires very specific environmental conditions–nano-Gauss magnetic fields, very dilute extragalactic plasma, and absence of axion background–to be significant. Even under optimal conditions, the mixing may induce quasar color shifts or CMB distortions [24]. The limits on these distortions indicate that while the photon-axion mixing can affect the supernovae observation results, it cannot be the dominant source of supernova dimming. Other mixing mechanisms in the bulk experience similar limitations.

In contrast, the wall-localized Chern-Simons mixing of visible and dark electromagnetism is much stealthier. A wall with sub-MeV tension does not produce significant gravitational anisotropies [25], nor does it lead to problems with overclosing the Universe [26]. In our case, the scales of the charges and tensions are parametrically far below MeV, and so the walls are cosmologically safe. They are not relics of a conventional high-temperature symmetry-breaking phase transition. Rather, they arise dynamically during the late-time top-form flux discharge process discussed in [1,4]. They are produced very late, with small tensions, and quickly collide and decay, facilitating rapid percolation of the true vacuum around the present epoch. The wall-wrapped bubbles will not be produced in very large electromagnetic fields because the UV cutoff is low [4]. The mixing is completely localized to the wall; it is neither a continuous flavor oscillation nor a frequency-dependent cumulative effect. A single wall produces a fixed power loss incurred only as light crosses it, in the frequency range bounded from above by the cutoff $\mathcal{M}$ of the top form sector.

For frequencies above the cutoff $\mathcal{M}$ the wall becomes transparent. As we noted in the Introduction, at scales above the cutoff the anomalous couplings are screened as $\mathcal{Z}_{nm} \sim \mathcal{Z}_{nm}(0)\mathcal{M}^2/k^2$, and so the mixing with paraphotons shuts down. The visible-to-dark conversions at high frequencies do not occur since those waves propagate through the wall like gamma rays through ordinary matter. This will come naturally in discretely evanescent dark energy [1], where the UV completion is an asymptotically free gauge theory with degenerate vacua and no Chern-Simons walls above the cutoff. Below the cutoff, Chern-Simons flavor-mixing terms $\sim\vec{E}^{(1)} \cdot \vec{B}^{(2)}$ and the like (as needed to maintain gauge symmetry and Lorentz symmetry) arise after particles with masses at the cutoff are integrated out. These terms intertwine the incoming states once they hit the wall. Off the wall, these terms are total derivatives and do not affect local fields.

On a wall with Chern-Simons terms, the jump in $\mathcal{H}$ enforces a single, localized mixing event. In fact, even polarization rotation can be viewed as a flavor mixing of photon helicities. If polarization rotation is sufficiently small to satisfy the CMB bounds [22],

$$\Delta\vartheta \sim 10^{-3}\,\text{radians}, \tag{20}$$

the visible-to-dark mixing could yield observable effects. As long as such walls are rare and have not percolated or collided extensively at small redshifts, low-tension walls could still be around, having evaded detection to date [27].

In [1] we considered a cutoff $\mathcal{M} \sim \text{milli} - \text{eV}$ for simplicity. In that framework, the cloaking of the visible light would only affect frequencies $\omega \lesssim \text{milli} - \text{eV}$; this means a fraction of the CMB would be converted to dark light, while optical range sources would remain unaffected. Minor modifications of the model in [1] allow $\mathcal{M} \sim \text{eV}$, coinciding with the scale of early dark energy [28]. In such a scenario, the wall can mix visible and dark waves with optical frequencies ~eV and below, while remaining transparent to ultraviolet waves and beyond. In such a framework one could also simultaneously incorporate both early and late dark energy if the late vacuum energy was not fully cancelled, e.g., for $\sqrt{\mathcal{Q}}/\mathcal{M} < 1$. In this work, we will be agnostic to which option is preferred.

As we noted above, we have neglected dark sector contributions to the incident waves from sources behind the wall, which is reasonable for sources made up of Standard Model particles. On the other hand the early Universe dynamics generically produces a relic dark radiation background due to inflation, reheating, and dark matter interactions. Cosmological constraints, particularly from BBN and late-time observations, limit these contributions to be subleading. In [1] we noted that early dark radiation in our case is generically colder than the early CMB, with $T^{\text{dark}}/T^{\text{visible}} \lesssim 0.001 - 0.01$. Consequently, the dark photon number is extremely suppressed in the CMB frequency range ($\omega \sim 0.1$–1 meV),

$$\left.\frac{N(\omega)_{\text{dark}}}{N(\omega)_{\text{visible}}}\right|_{\max} \simeq e^{-\omega/T^{\text{dark}}} < 10^{-5}, \tag{21}$$

so the only relevant mixing process is on-the-wall conversion of CMB photons into dark photons. Being frequency independent below the cutoff $\mathcal{M}$, the wall creates a dark "snapshot" of the CMB, depleting the visible CMB by $\rho = 1 - \gamma$, with γ from Eq. (18). In our approximation, where we ignored cosmic expansion, for a large wall or bubble the depletion does not induce significant corrections to anisotropies. The main effect is a rescaling of the CMB black body spectrum at wall crossing, $N(\omega)d\omega = \frac{\omega^3 d\omega}{4\pi^3}(e^{\omega/T} - 1)^{-1}$, by a factor of ρ. Because of the form invariance of the CMB photon number $N(\omega)_{\text{visible}}d\omega$ under rescaling of frequency, the factor ρ can be absorbed into the measured temperature inside the bubble,

$$T_-^{\text{visible}} = \rho^{1/4} T_+^{\text{visible}}. \tag{22}$$

This shifts the last scattering surface of the CMB as seen by a terrestrial observer. The wall-induced depletion of the CMB power thus allows the CMB to reach the observed temperature more efficiently, needing less time and space to "cool" down to the observed temperature. If we combine this with cosmic expansion, the CMB temperature would decrease in less time, which could increase H_0, due to the scaling invariance described above. The scaling remains an approximate degeneracy between H_0 and T_0^{visible} in an expanding universe [29].

After crossing the wall, the lost CMB photons convert into the dark ones, replicating the CMB spectrum at the crossing. Their power is given by γ times the original CMB power. We can interpret this as the generation of the dark thermal background on our side of the wall, with the temperature $T_-^{\text{dark}} = \gamma^{1/4} T_+^{\text{visible}}$, similarly to (22). Observational constraints on the T_{CMB} evolution [30] are precise to ~0.5% at low redshifts and a few percent at $z \sim 0.65$, with the fitting of evolution done relative to Λ cold dark matter. While these bounds are pretty tight, a mere 2% variation of T_{CMB} yields a power reduction of $(0.98)^4 \sim 0.92$, or ~8%, induced by the wall-borne Chern-Simons mixing. Correspondingly the dark background temperature right after the wall crossing would be $T_-^{\text{dark}}/T_+^{\text{visible}} \simeq \gamma^{1/4}$, which could be as much as 0.5. This might contribute to ΔN_{eff}, but only at very low redshifts, much later than recombination.

In the case where $\mathcal{M} \sim \text{few eV}$, the depletion of visible light by conversion to dark photons by our mechanism would also affect optical range sources. The light-to-dark conversion could reduce optical flux from sources behind the wall by the same fraction as the CMB, due to its frequency independence below the cutoff $\mathcal{M}$. This could additionally impact the inferred Hubble rate from the CMB versus low-redshift sources [29,31].

Such attenuation may also lead to violations of the Etherington reciprocity relation, which relates luminosity and angular diameter distances via $D_L(z)/D_A(z) = (1+z)^2$ [32,33]. While most tests probe smooth, cumulative violations of this ratio, steplike luminosity changes at specific redshifts have been considered recently [34] as a possible way to address the H_0 tension. In our scenario, the violation would be steplike location-wise, at a specific redshift with a correlated change in vacuum energy due to the top form flux discharge. It would also only affect the light with frequencies below the cutoff $\mathcal{M}$ of the dark top form sector, rapidly shutting off at higher frequencies.

In summary, Chern-Simons terms mixing visible and dark $U(1)$ vector gauge fields on discretely evanescent dark energy walls can alter the apparent brightness of sources behind the wall. The jump in luminosity is correlated with a vacuum energy jump. The luminosity jump could be interpreted as a shift of the last scattering surface, after redefining the CMB temperature on the far side of the wall. The mechanism is stealthy, potentially evading most current probes. Observations of the CMB polarization rotation [22] and the H_0 tension [31] may offer the first hints of such dark sector physics, motivating further investigation of these exotic phenomena.

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[1] N. Kaloper, Discretely evanescent dark energy, J. Cosmol. Astropart. Phys. 11 (2025) 075.
[2] N. Kaloper, Alternative to axions, Phys. Rev. D **113**, L011701 (2026).
[3] N. Kaloper, Quantal theory of restoration of strong *CP* symmetry, Phys. Rev. D **113**, 075020 (2026).
[4] N. Kaloper, Electromagnetic couplings of dark domain walls, Phys. Rev. D **114**, L011904 (2026).
[5] M. Lüscher, The secret long range force in quantum field theories with instantons, Phys. Lett. **78B**, 465 (1978).
[6] G. Gabadadze, Modeling the glueball spectrum by a closed bosonic membrane, Phys. Rev. D **58**, 094015 (1998).
[7] G. Gabadadze and M. Shifman, QCD vacuum and axions: What's happening?, Int. J. Mod. Phys. A **17**, 3689 (2002).
[8] G. Dvali, Three-form gauging of axion symmetries and gravity, arXiv:hep-th/0507215.
[9] G. Dvali, A vacuum accumulation solution to the strong *CP* problem, Phys. Rev. D **74**, 025019 (2006).
[10] N. Kaloper and L. Sorbo, Where in the string landscape is quintessence, Phys. Rev. D **79**, 043528 (2009).
[11] N. Kaloper and L. Sorbo, A natural framework for chaotic inflation, Phys. Rev. Lett. **102**, 121301 (2009).
[12] N. Kaloper, Axion flux monodromy discharges relax the cosmological constant, J. Cosmol. Astropart. Phys. 11 (2023) 032.
[13] L. B. Okun, Limits of electrodynamics: Paraphotons?, Sov. Phys. JETP **56**, 502 (1982).
[14] H. Georgi, P. H. Ginsparg, and S. L. Glashow, Photon oscillations and the cosmic background radiation, Nature (London) **306**, 765 (1983).
[15] B. Holdom, Two $U(1)$'s and epsilon charge shifts, Phys. Lett. B **166**, 196 (1986).
[16] H. P. Nordberg and G. F. Smoot, The cosmic microwave background spectrum: An analysis of observations, arXiv: astro-ph/9805123.
[17] M. Fabbrichesi, E. Gabrielli, and G. Lanfranchi, The dark photon, arXiv:2005.01515.
[18] J. D. Jackson, *Classical Electrodynamics* (John Wiley and Sons, NY USA, 1998).
[19] M. C. Huang and P. Sikivie, The structure of axionic domain walls, Phys. Rev. D **32**, 1560 (1985).
[20] D. Harari and P. Sikivie, Effects of a Nambu-Goldstone boson on the polarization of radio galaxies and the cosmic microwave background, Phys. Lett. B **289**, 67 (1992).
[21] A. D. Shapere, S. Trivedi, and F. Wilczek, Dual dilaton dyons, Mod. Phys. Lett. A **06**, 2677 (1991).
[22] E. Komatsu, New physics from the polarized light of the cosmic microwave background, Nat. Rev. Phys. **4**, 452 (2022).
[23] C. Csaki, N. Kaloper, and J. Terning, Dimming supernovae without cosmic acceleration, Phys. Rev. Lett. **88**, 161302 (2002).
[24] A. Mirizzi, G. G. Raffelt, and P. D. Serpico, Photon-axion conversion in intergalactic magnetic fields and cosmological consequences, Lect. Notes Phys. **741**, 115 (2008).
[25] Y. B. Zeldovich, I. Y. Kobzarev, and L. B. Okun, Cosmological consequences of the spontaneous breakdown of discrete symmetry, Zh. Eksp. Teor. Fiz. **67**, 3 (1974).
[26] R. A. Battye, M. Bucher, and D. Spergel, Domain wall dominated universes, arXiv:astro-ph/9908047.
[27] S. Koren, Y. Tsai, and R. Wang, Boiling after the dust settles: Constraining first-order phase transitions during dark energy domination, Phys. Rev. Lett. **137**, 051002 (2026).
[28] V. Poulin, T. L. Smith, T. Karwal, and M. Kamionkowski, Early dark energy can resolve the Hubble tension, Phys. Rev. Lett. **122**, 221301 (2019).
[29] M. M. Ivanov, Y. Ali-Haïmoud, and J. Lesgourgues, H_0 tension or T0 tension?, Phys. Rev. D **102**, 063515 (2020).
[30] G. Hurier, N. Aghanim, M. Douspis, and E. Pointecouteau, Measurement of the T_{CMB} evolution from the Sunyaev-Zel'dovich effect, Astron. Astrophys. **561**, A143 (2014).
[31] M. Kamionkowski and A. G. Riess, The Hubble tension and early dark energy, Annu. Rev. Nucl. Part. Sci. **73**, 153 (2023).
[32] I. M. H. Etherington, On the definition of distance in general relativity, Philos. Mag. Ser. 7 **15**, 761 (1933).
[33] B. A. Bassett and M. Kunz, Cosmic distance-duality as a probe of exotic physics and acceleration, Phys. Rev. D **69**, 101305 (2004).
[34] E. M. Teixeira, W. Giarè, N. B. Hogg, T. Montandon, A. Poudou, and V. Poulin, Implications of distance duality violation for the H_0 tension and evolving dark energy, Phys. Rev. D **112**, 023515 (2025).