

Gap in the black-body spectrum at low temperature

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ABSTRACT: We postulate that the $U(1)_Y$ factor of the Standard Model is an effective manifestation of $SU(2)$ gauge dynamics being dynamically broken by nonperturbative effects. The modified propagation properties of the photon at low temperatures and momenta are computed. As a result of strong screening, the presence of a sizable gap in the spectral power of a black body at temperatures $T = 5 \cdots 20$ K and for low frequencies is predicted: A table-top experiment should be able to discover this gap. If the gap is observed then the Standard Model's mechanism for electroweak symmetry-breaking is endangered by a contradiction with Big-Bang nucleosynthesis. Based on our results, we propose an explanation for the stability of cold, old, dilute, and large clouds of atomic hydrogen in between spiral arms of the outer galaxy.

KEYWORDS: Duality in Gauge Field Theories, Spontaneous Symmetry Breaking.

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1. Introduction

The concept that electromagnetic waves, in analogy to the propagation of distortions in fluids, require a medium to travel - the ether -, was employed by Maxwell to deduce his famous equations [1]. The ether was abandoned with the development of Special Relativity (SR) which is based on two empirically founded postulates [2]: relativity of uniform motion and constancy of the speed of light c . SR implies that in observing the propagation of a monochromatic light wave (photon) no inertial frame of reference is singled out. This situation changes when thermalized radiation is considered: The very process of thermalization proceeds by interactions between the photon and electrically charged, massive matter. The latter's center of inertia, however, defines a preferred rest frame. This, at least, is the standard notion of how a temperature T emerges in a photon gas.

On a microscopic level, photon interactions are described by quantum mechanical transitions. The underlying and very successful field theory, Quantum Electrodynamics (QED) [3–5], is based on a $U(1)$ gauge group. In the present Standard Model of particle physics (SM) a progenitor of this symmetry is called $U(1)_Y$. On a thermodynamical level, the microscopic details of the emission and absorption processes are averaged away. As a consequence, thermalized photons exhibit a universal black-body spectrum whose shape solely depends on T and the two constants of nature c and \hbar .

In this Letter we explore the possibility that the $U(1)_Y$ factor of the SM's gauge group is not fundamental [6–9]. In this context, QED may break down under exceptional conditions. Such an exception would take place in a thermalized photon gas at temperatures not far above 2.73 Kelvin (K): The notion of a gas of interaction-free photons then would require revision. Namely, embedding $U(1)_Y$ into the fundamental, nonabelian gauge group $SU(2)$, invokes a mass scale $\Lambda \sim 10^{-4}$ electron volts (eV) [6, 10–12] which affects black-body spectra. The $SU(2)$ gauge symmetry implies the existence and relevance of the Yang-Mills scale Λ on the quantum level [13, 14]. We will discuss below why we use the name $SU(2)_{\text{CMB}}$ (CMB for *Cosmic Microwave Background*) for the fundamental gauge symmetry also describing photon propagation at low temperatures.

Let us now list some bulk properties of SU(2) gauge theory in four dimensions [6, 12, 15–22]. First, three phases exist: a deconfining, a preconfining, and a confining one (in order of decreasing temperature). Only the deconfining phase is relevant for the present discussion. Upon a spatial coarse-graining this phase is described by an effective field theory [6]. Second, for T much greater than the scale Λ all thermodynamical quantities reach their Stefan-Boltzmann limits in a power-like way. Third, there is a nontrivial ground state, obeying an equation of state $P_{gs} = -\rho_{gs} = -4\pi\Lambda^3 T$, in the deconfining phase. This ground state is tied to the presence of interacting topological defects (calorons, topology-changing quantum fluctuations [6, 21]). Through interactions with the ground state two types of gauge modes (V^\pm) acquire a temperature-dependent mass while the third type remains massless (γ). It is important to note that at a critical temperature T_c (boundary between deconfining and preconfining phase) γ 's partners V^\pm acquire an infinite mass and thus decouple thermodynamically. Moreover, within the deconfining phase quantum fluctuations are severely constrained in the effective theory: The interactions of the three types of gauge bosons are very weak [12, 22].

Taking these interactions into account, γ 's dispersion law

$$\omega^2 = \mathbf{p}^2 \tag{1.1}$$

modifies as [12]

$$\omega^2 = \mathbf{p}^2 + G(\omega, \mathbf{p}, T, \Lambda). \tag{1.2}$$

In eq. (1.2) ω denotes the energy of a γ -mode with spatial momentum \mathbf{p} . The screening function G depends on ω , \mathbf{p} , temperature T , and the Yang-Mills scale Λ . Notice that in writing eqs. (1.1) and (1.2) we use natural units: $c = \hbar = k_B = 1$ where k_B is Boltzmann's constant. Our results indicate that for $T \gg T_c$ the function G is negative with a negligibly small modulus (antiscreening). However, for T a few times T_c and for small $|\mathbf{p}|$ the function G becomes positive and reaches sizable values ($> |\mathbf{p}|^2$). That is, the γ -mode acquires a screening mass. If emitted with the dispersion law of eq. (1.1) then the dispersion law of eq. (1.2) is violated: γ can penetrate the plasma only up to a distance $\sim G^{-1/2}$. A useful analogy is a rain-drop falling onto the surface of a lake where it is absorbed immediately.

At $T_c = \frac{\lambda_c}{2\pi}\Lambda$ ($\lambda_c = 13.87$ [6]) for the (second-order like) transition to the preconfining phase we have

$$\lim_{T \searrow T_c} G(\omega, \mathbf{p}, T, \Lambda) = 0. \tag{1.3}$$

Thus no (anti)screening is seen in γ -propagation at $T = T_c$. The above results match with observational ($T \sim T_{\text{CMB}}$) and daily ($T \gg T_{\text{CMB}}$) experience that the photon's dispersion law is the one in eq. (1.1). Hence we are led to identify γ with the photon ($U(1)_Y \subset \text{SU}(2)_{\text{CMB}}$) and conclude that $T_c = T_{\text{CMB}} \sim 2.73 \text{ K} = 2.35 \times 10^{-4} \text{ eV}$ for $\text{SU}(2)_{\text{CMB}}$. In spite of the fact that such an identification is rather unconventional we believe that it is worthwhile to pursue its consequences, be it only to falsify such a scenario.

Within a cosmological context one derives that the photon remains massless only for a finite period of time, $\Delta t \leq 2.2$ billion years, in the future [11]. This is due to the fact that the transition to the preconfining phase invokes a nonvanishing coupling of the

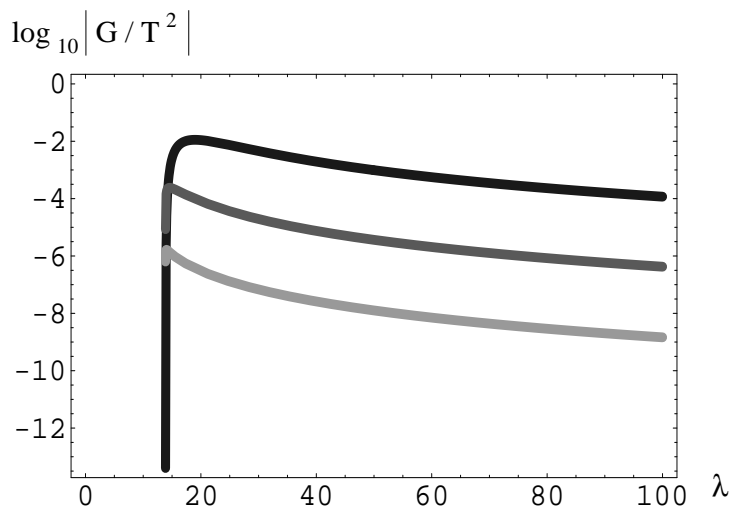


Figure 1: $\log_{10} \left| \frac{G}{T^2} \right|$ as a function of $\lambda \equiv \frac{2\pi T}{\Lambda}$ for $X = 1$ (black), $X = 5$ (dark grey), and $X = 10$ (light grey) where $X \equiv \frac{|\mathbf{p}|}{T}$.

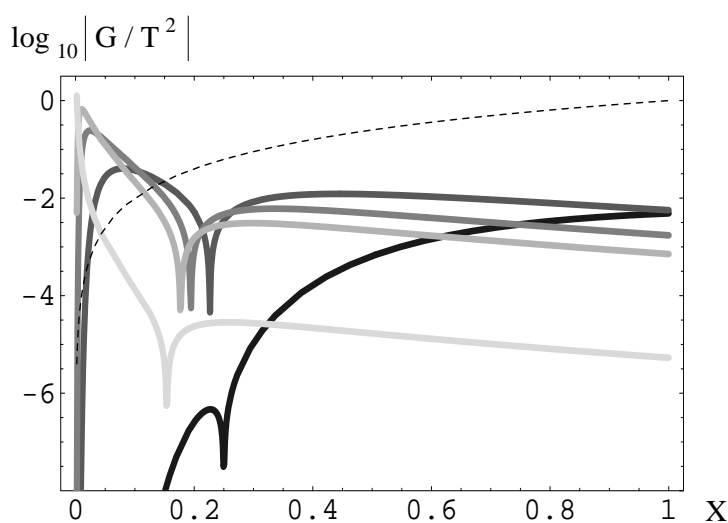


Figure 2: $\log_{10} \left| \frac{G}{T^2} \right|$ as a function of X for $\lambda = 1.12 \lambda_c$ (black), $\lambda = 2 \lambda_c$ (dark grey), $\lambda = 3 \lambda_c$ (grey), $\lambda = 4 \lambda_{c,E}$ (light grey), $\lambda = 20 \lambda_c$ (very light grey). The dashed curve is a plot of the function $f(X) = 2 \log_{10} X$. Photons are strongly screened at X -values for which $\log_{10} \left| \frac{G}{T^2} \right| > f(X)$, that is, to the left of the dashed line. The dips correspond to the zeros of G .

photon to a newly emerging ground state: The remaining gauge symmetry $U(1)_Y$ is then broken dynamically (in contrast to the deconfining phase); a familiar effect in macroscopic superconductivity [23, 24].

2. Modified black-body spectra at low temperatures and low frequencies

In [12] we have calculated the function G appearing in eq. (1.2) from the photon's polarization tensor. We have made the assumption that $\omega = |\mathbf{p}|$ on the right-hand side of eq. (1.2). (For $\omega = |\mathbf{p}|$ G is *real*.) This is a consistent approximation as long as $\frac{|G|}{\omega^2} \ll 1$, see below. In figures 1 and 2 the function $\log_{10} \left| \frac{G}{T^2} \right|$ is depicted in dependence of (dimensionless) temperature $\lambda \equiv \frac{2\pi T}{\Lambda}$ and (dimensionless) momentum $X \equiv \frac{|\mathbf{p}|}{T}$, respectively.

As figure 1 shows, at fixed values of X the function $|G|$ falls off in a power-like way at large temperatures. Equidistance of the curves for equidistant values of $X \geq 1$ indi-

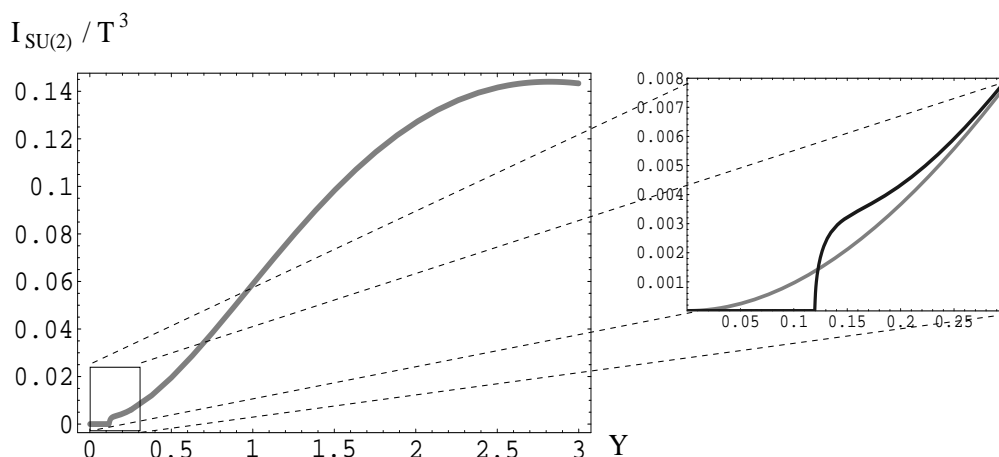


Figure 3: Dimensionless black-body spectral power $\frac{I_{\text{SU}(2)}}{T^3}$ as a function of the dimensionless frequency $Y \equiv \frac{\omega}{T}$. The black curve in the magnified region depicts the modification of the spectrum as compared to $\frac{I_{\text{U}(1)}}{T^3}$ (grey curve) for $T = 10$ K.

cates exponential suppression in X . For $T \searrow T_c$ the thermodynamical decoupling of V^\pm -modes at the phase boundary leads to a rapid drop of $|G|$. In figure 2 the low-momentum behavior of $|G|$ at fixed temperatures not far above T_c is depicted. For $\text{SU}(2)_{\text{CMB}}$ the (dimensionless) temperatures $\lambda = 1.12 \lambda_c, 2 \lambda_c, 3 \lambda_c, 4 \lambda_c$ and $\lambda = 20 \lambda_c$ convert into $T = 3.02$ K, 5.5 K, 8.2 K, 10.9 K, and $T = 55$ K, respectively. For $T = 3.02$ K $\sim T_{\text{CMB}}$ (black curve) and $T = 55$ K (very light grey curve) the regime, where photons are strongly screened, is too small to be resolved by existing low-temperature black-body (LTBB) observations and experiments [25]. For the other temperatures considered in figure 2 there is a sizable range of X -values for this effect. (We only mention here that photons do propagate again at very small momenta [12].)

The spectral power $I_{\text{U}(1)}(\omega)$ for a black body subject to the gauge symmetry $\text{U}(1)$ is given as

$$I_{\text{U}(1)}(\omega) = \frac{1}{\pi^2} \frac{\omega^3}{\exp[\frac{\omega}{T}] - 1}. \quad (2.1)$$

For $\text{SU}(2)_{\text{CMB}}$ this modifies as

$$I_{\text{U}(1)}(\omega) \rightarrow I_{\text{SU}(2)}(\omega) = I_{\text{U}(1)}(\omega) \times \frac{(\omega - \frac{1}{2} \frac{d}{d\omega} G) \sqrt{\omega^2 - G}}{\omega^2} \theta(\omega - \omega^*) \quad (2.2)$$

where ω^* is the root of $\omega^2 = G$, and $\theta(x)$ is the Heaviside step function. In figure 3 the modification of the black-body spectrum according to eq. (2.2) is depicted for $T = 10$ K: There is no spectral power at frequencies $\omega < 0.12 T$ whereas there is a (rapidly decreasing) excess at frequencies $\omega > 0.12 T$.

Let us now investigate how reliable the approximation $\omega = |\mathbf{p}|$ is when evaluating the function G . In figure 4 a plot of $\frac{G}{\omega^2}$ is shown as a function of $Y \equiv \frac{\omega}{T}$ for $T = 5$ K (black curve) and for $T = 10$ K (grey curve) in the vicinity of G 's zero Y_0 . To the right of Y_0 the condition $\frac{|G|}{\omega^2} \ll 1$ is well satisfied for $\omega = |\mathbf{p}|$, to the left of Y_0 this continues to be

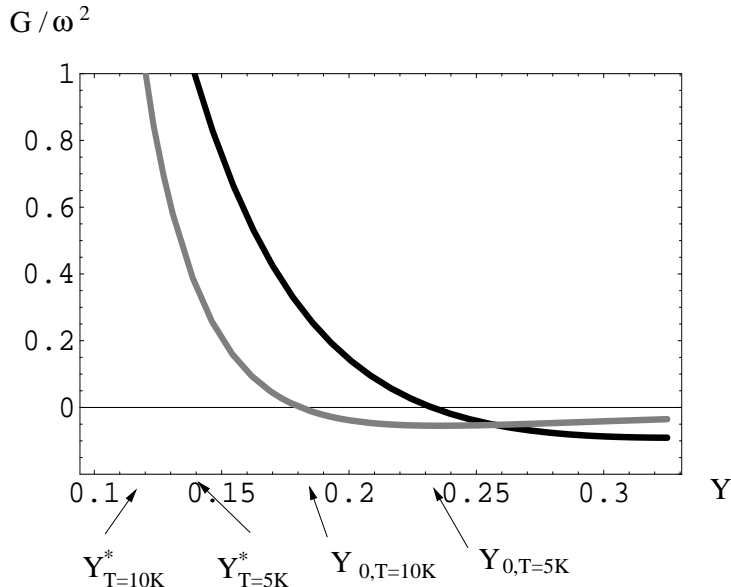


Figure 4: $\frac{G}{\omega^2}$ as a function of $Y \equiv \frac{\omega}{T}$ for $T = 5$ K (black curve) and $T = 10$ K (grey curve).

a reasonable approximation almost down to $Y^* \equiv \frac{\omega^*}{T}$ because of the large negative slope of the function $\frac{G}{\omega^2}$ in the vicinity of Y^* : Although our approximation is doomed to break down at Y^* it is still valid for values of Y slightly above Y^* where the tendency towards large G is seen. For an experiment to detect the reshuffling of spectral power as indicated in figure 3 the spectrometer must not be further away from the aperture of the LTBB than $G^{-1/2}$.

Let us now discuss how sensitive the measurement of the LTBB spectral intensity $I_{\text{SU}(2)}(\omega)$ needs to be in order to detect the spectral gap setting in at ω^* . A useful criterion is determined by the ratio $R(Y^*)$ of $I_{\text{U}(1)}(\omega^*)$ and $I_{\text{U}(1)}(\omega_{\text{max}})$ where $\omega_{\text{max}} = 2.82T$ (in natural units) is the position of the maximum of $I_{\text{U}(1)}$:

$$R(Y^*) \equiv \frac{I_{\text{U}(1)}(\omega^*)}{I_{\text{U}(1)}(\omega_{\text{max}})} = \frac{1}{1.42144} \frac{(Y^*)^3}{\exp(Y^*) - 1}. \quad (2.3)$$

For $T = 80$ K, which was experimentally realized in [25], we have $R(Y^* = 0.0366) = 9 \times 10^{-4}$. To achieve such a high precision is a challenging task. To the best of the authors knowledge only the overall and not the spectral intensity of the LTBB was measured in [25]. For $T = 5$ K one has $R(Y^* = 0.14) = 1.2 \times 10^{-2}$. Thus at low temperatures the precision required to detect the spectral gap is within the 1%-range. It is, however, experimentally challenging to cool the LTBB down to these low temperatures. To the best of the authors knowledge a precision measurement of the spectral power in the low-frequency regime of a LTBB at $T = 5 \cdots 10$ K has not yet been performed. We know, however, that such an experiment is well feasible [26]: It will represent an important and inexpensive (on particle-physics scales) test of the postulate $\text{SU}(2)_{\text{CMB}} \stackrel{\text{today}}{=} \text{U}(1)_Y$.

3. Stability of dilute and cold hydrogen clouds in the outer galaxy

In [7, 27] the existence of a large (up to 2 kpc), old (estimated age ~ 50 million years), cold (mean brightness temperature $T_B \sim 20$ K with cold regions of $T_B \sim 5 \cdots 10$ K), dilute

(number density: $\sim 1.5 \text{ cm}^{-3}$) and massive (1.9×10^7 solar masses) intergalactic cloud (GSH139-03-69) of atomic hydrogen (HI) forming an arc-like structure in between spiral arms was reported. In [28] and references therein smaller structures of this type were identified. These are puzzling results which do not fit into the dominant model for the interstellar medium [27]. Moreover, considering the typical time scale for the formation of H_2 molecules out of HI of about 10^6 yr [28] at these low temperatures and low densities clashes with the inferred age of the structure observed in [7].

To the best of our knowledge there is no standard explanation for the existence and the stability of such structures. We wish to propose a scenario possibly explaining the stability based on $\text{SU}(2)_{\text{CMB}}$. Namely, at temperatures $T_B \sim 5 \cdots 10 \text{ K}$, corresponding to $T_B \sim 2 \cdots 4 T_{\text{CMB}}$, the function G for photons with momenta ranging between $|\mathbf{p}^*| \sim 0.15 T_B > |\mathbf{p}_c| > |\mathbf{p}_{\text{low}}|$ is such that it strongly suppresses their propagation, see figures 2. We mention in passing only that $|\mathbf{p}_{\text{low}}| < 0.02 T_B$ depends rather strongly on temperature [12].

Incidentally, the regime for the wavelength l_c associated with $|\mathbf{p}_c|$ is comparable to the interatomic distance $\sim 1 \text{ cm}$ in GSH139-03-69: At $T = 5 \text{ K}$ we have $l^* = 2.1 \text{ cm} \leq l_c \leq 8.8 \text{ cm} = l_{\text{low}}$, at $T = 10 \text{ K}$ we have $l^* = 1.2 \text{ cm} \leq l_c \leq 1.01 \text{ m} = l_{\text{low}}$. Thus the photons mediating the dipole interaction between HI particles practically do not propagate: the dipole force at these distances appears to be switched off. As a consequence, H_2 molecules are prevented from forming at the temperatures and densities which are typical for GSH139-03-69.

The astrophysical origin of the structure GSH139-03-69 appears to be a mystery. The point we are able to make here is that once such a cloud of HI particles has formed it likely remains in this state for a long period of time.

4. Implications for the Standard Model and conclusions

We will now argue that if $\text{SU}(2)_{\text{CMB}}$ is, indeed, realized in Nature then the SM's Higgs mechanism for electroweak symmetry-breaking is not. The key is to work out the consequences of $\text{SU}(2)_{\text{CMB}}$ for Big-Bang nucleosynthesis. The SM predicts that within this cosmological epoch the number of relativistic degrees of freedom g_* is given as

$$g_* = 5.5 + \frac{7}{4} N_\nu \tag{4.1}$$

where $1.8 \leq N_\nu \leq 4.5$ [29]. This prediction relies on the following consideration: the neutron to proton fraction n/p at freeze-out is given as $n/p = \exp[-Q/T_{\text{fr}}] \sim 1/6$ where $Q = 1.293 \text{ MeV}$ is the neutron-proton mass difference, and one has

$$T_{\text{fr}} \sim \left(\frac{g_* G_N}{G_F^4} \right)^{1/6}. \tag{4.2}$$

In eq. (4.2) G_N denotes Newton's constant, and

$$G_F = \pi \frac{\alpha_w}{\sqrt{2} m_W^2} \sim 1.17 \times 10^{-5} \text{ GeV}^{-2} \tag{4.3}$$

is the Fermi coupling at zero temperature. To use the zero-temperature value of G_F at $T_{\text{fr}} = 1 \text{ MeV}$, as it is done in eq. (4.2), is well justified by the large ‘electroweak scale’ $v = 247 \text{ GeV}$ in the SM: the vacuum expectation of the Higgs-field. Invoking $\text{SU}(2)_{\text{CMB}}$ yields additional six relativistic degrees of freedom (V^\pm with three polarizations each) at $T_{\text{fr}} = 1 \text{ MeV}$: a result which exceeds the above cited upper bound for N_ν . The value of $T_{\text{fr}} \sim 1 \text{ MeV}$ is rather reliably extracted from the primordial ${}^4\text{He}$ abundance $Y_p \sim 0.25$ and the subsequent determination of $n/p \sim 1/6$ [29]. To save this value of T_{fr} one needs to prescribe a value of G_F at $T_{\text{fr}} = 1 \text{ MeV}$ which is about 12% larger than the value of G_F in eq. (4.3). Since G_F at $T = 0$ is measured to per mille accuracy there would be a contradiction with electroweak SM physics. A larger value of $G_F(T = T_{\text{fr}})$ is, however, expected if the weak interactions are based on a higgsless $\text{SU}(2)$ gauge theory of Yang-Mills scale $\Lambda \sim 0.5 \text{ MeV}$, see the discussion in [11].

The prime physical system, for which our results are relevant, is the cosmic microwave background. Namely, at small redshift ($z < 20$) the screening effects of intermediate V^\pm bosons have the potential to explain the large-angle anomalies in the power spectra as they were reported in [8].

For the above reasons the importance of an experimental verification or falsification of the postulate $\text{SU}(2)_{\text{CMB}} \stackrel{\text{today}}{=} \text{U}(1)_Y$ involving the low-temperature, low-frequency regimes in black-body spectra is evident.

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