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Impact of the cosmic neutrino background on long-range force searches

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Abstract: Light bosons can mediate long-range forces. We show that light bosonic mediators interacting with a background medium, in particular, with the cosmic neutrino background (CνB), may induce medium-dependent masses which could effectively screen long-range forces from detection. This leads to profound implications for long-range force searches in e.g. the Eöt-Wash, MICROSCOPE, and lunar laser-ranging (LLR) experiments. For instance, we find that when the coupling of the mediator to neutrinos is above 3×10^{-10} or 5×10^{-13}, bounds from LLR and experiments employing the Sun as an attractor, respectively, would be entirely eliminated. Larger values of the coupling can also substantially alleviate bounds from searches conducted at shorter distances.

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Contents

1 Introduction

Among the four fundamental interactions in nature — gravity, electromagnetism, and the strong and weak interactions — two of them are long-range forces, observable at macroscopic scales. It is tempting to conceive that new interactions arising from theories beyond the standard model (SM) might generate extra long-range forces. If discovered, new long-range forces would have paradigm-shifting implications for new physics explorations and could even lead to revolutionary applications in macroscopic-scale technologies.

Long-range forces are mediated by massless or sufficiently light mediators. Some examples include light scalar or vector bosons in hidden sectors or gauge extensions [1–8], axions [9–11], majorons [12], dilatons [13–15], dark matter [16, 17], or even neutrinos [18–24]. In recent years, the interest in light, feebly-interacting particles has been persistently growing (see refs. [25, 26] for reviews), while long-range force searches (also known as *fifth-force searches* in the literature) play a crucial role in constraining ultra-light particles if their Compton wavelengths can reach macroscopic extent.

Currently, the most sensitive searches are based on torsion-balance tests of gravity [27]. New forces mediated by light bosons with generic couplings to SM fermions might cause observable deviations in the probe of the weak equivalent principle (WEP) or the inverse square law of gravity. At large distance scales comparable to the Earth's radius (corresponding to a mediator mass below 10^{-14} eV), the Eöt-Wash experiment has excluded the couplings to electrons and baryons above $10^{-23} \sim 10^{-24}$. If such a small coupling arises from high-energy theories (e.g. generated by heavy loops [3, 4]), then the extremely small value can be reinterpreted as a probe of new physics at very high energy scales. Other experimental probes including lunar laser-ranging (LLR) [28–32], the detection of gravitational waves at LIGO/Virgo [33–37], and the very recent MICROSCOPE experiment [38] have their respective advantages in long-range force searches.

Given the remarkably high sensitivity of these experiments to new long-range forces, it is important to ask: to what extent are these bounds valid and what factors may significantly alter the bounds?

Taking the example of electromagnetism, the Coulomb potentials of electrons and nuclei always cancel each other at large scales in electrically neutral matter, and electromagnetic waves can often be screened by metals. Consequently, the photon and photon-like particles[1] are unable to cause observable effects in these experiments. For axions and majorons which are pseudoscalars, there are similar cancellations between particles of opposite spin directions in unpolarized matter. For light scalar or vector bosons, in general the forces on individual particles can be added coherently without cancellation. However, such bosons could still evade the aforementioned constraints if their propagation in a medium behaves differently from that in the vacuum.

In this study, we investigate potential medium effects that may exert a significant influence on the searches for long-range forces. We consider a light scalar mediator ϕ that is generically coupled to all SM fermions including neutrinos, and show that the propagation of ϕ in a medium such as normal matter or the cosmic neutrino background (CνB) could be modified by coherent forward scattering with medium particles, generating a medium-dependent mass for ϕ. Such a mass could shorten the interaction range and alleviate some experimental bounds on long-range forces. We examine this effect and find that the medium effect caused by normal matter of the Earth or the Sun is negligible in experiments employing them as the attractor. On the other hand, the CνB which is ubiquitous and cannot be evacuated in any experiments, can cause a considerably strong medium effect due to the lightness of neutrinos. We show that within a broad range of the neutrino coupling allowed by cosmology, existing bounds on long-range forces can be altered by the CνB very significantly.

This work is organized as follows. In section 2, we briefly review the formalism for generating long-range forces in the vacuum and elucidate the medium effect to be taken into account in the formalism. The medium effect can be included by solving a modified field equation, which is solved in section 3 assuming a spherically symmetric configuration. Then we inspect to what extent the experimental bounds on long-range forces can be altered by the CνB in section 4 and draw our conclusions in section 5. Some minus sign issues are addressed in appendix A and a derivation of the medium effect is presented in appendix B.

[1]By "photon-like", we mean that the effective couplings to electrons and nuclei are proportional to their electric charges. The dark photon and the $L_\mu - L_\tau$ gauge boson with loop-induced couplings to electrons and baryons are in this category.

2 Long-range forces and the medium effect

2.1 Long-range forces in the vacuum

Let us briefly review the generation of long-range forces from light bosons, mainly following the formalism in refs. [39, 40]. We start by considering the following Lagrangian:

$$\mathcal{L} = \frac{1}{2}(\partial\phi)^2 - \frac{1}{2}m_\phi^2\phi^2 + \overline{\psi}i\partial\!\!\!/\psi - m_\psi\overline{\psi}\psi - y_\psi\phi\overline{\psi}\psi\,, \tag{2.1}$$

where ϕ is a light scalar and ψ denotes a generic fermion, with their respective masses m_ϕ and m_ψ. The Lagrangian leads to the following equations of motion (EOMs):

$$i\partial\!\!\!/\psi - (m_\psi + y_\psi\phi)\psi = 0\,, \tag{2.2}$$
$$(\partial^2 + m_\phi^2)\phi + y_\psi\overline{\psi}\psi = 0\,. \tag{2.3}$$

Eq. (2.2) implies that in the presence of a background ϕ field, the mass of ψ is essentially shifted to $m_\psi + y_\psi\phi$. For nonrelativistic ψ particles, this is equivalent to shifting its energy by $y_\psi\phi$. Thus one can introduce an effective potential to account for the influence of ϕ on ψ:

$$V = y_\psi\phi\,. \tag{2.4}$$

When ϕ has a nonzero gradient ($\nabla\phi \neq 0$), it causes a force, $F = y_\psi\nabla\phi$, acting on ψ particles. For relativistic ψ particles, the mass shift is not fully equivalent to the energy shift, causing a reduction of the Yukawa force. We refer to ref. [41] for detailed discussions on this issue.

For a bulk of ψ particles statically distributed in space, we replace $\overline{\psi}\psi$ in eq. (2.3) with its ensemble average $\langle\overline{\psi}\psi\rangle \approx n_\psi$ where n_ψ is the number density of ψ particles. Since it is static, we can neglect the temporal derivative in eq. (2.3) and write it as

$$\left[\nabla^2 - m_\phi^2\right]\phi = y_\psi n_\psi \text{ (non-relativistic)}. \tag{2.5}$$

Assuming the distribution is spherically symmetric and using spherical coordinates, we have $\nabla^2\phi = \frac{\partial^2\phi}{\partial r^2} + 2\frac{\partial\phi}{r\partial r} = \frac{u''(r)}{r}$ where $u(r) \equiv r\phi(r)$. This allows us to further rewrite eq. (2.5) as

$$\left[\frac{d^2}{dr^2} - m_\phi^2\right]u(r) = r y_\psi n_\psi(r)\,, \tag{2.6}$$

which can be solved analytically for some simple forms of $n_\psi(r)$, including e.g. constant or piecewisely constant n_ψ, or the exponential form $n_\psi \propto e^{-r\kappa}$ [39]. In particular, if $n_\psi = N_\psi\delta^3(\mathbf{r})$, the solution is $u(r) = -y_\psi N_\psi e^{-m_\phi r}/(4\pi)$, resulting in the well-known Yukawa potential:

$$V = -\frac{y_\psi^2 N_\psi}{4\pi r}e^{-m_\phi r}\,. \tag{2.7}$$

The force arising from eq. (2.7), $F = \nabla V$, is considered long-range if $m_\phi \ll r^{-1}$, in which case the exponential suppression can be neglected. As is well known, the potential is negative, implying that the Yukawa force is attractive. For vector interactions, one can derive a similar potential which is positive, corresponding to a repulsive force.

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2.2 The medium effect

In a medium, the propagation of the ϕ field could be significantly altered, similar to various known medium effects on electromagnetic waves. For instance, in plasma, due to interactions with free charged particles, the photon acquires an effective mass known as the plasmon mass. In cold (non-ionized) gas such as the air, where electrons and nuclei are bound into atoms or molecules, the response of atomic or molecular electric dipoles to electromagnetic waves slightly slows down the propagation of light by a factor of $1/n$ with n the refraction index of the medium.

For ϕ propagating through a medium containing a large number of free ψ particles, the medium effect is similar to the plasmon mass, giving rise to the following mass correction for the ϕ field:

$$m_\phi^2 \to m_\phi^2 + y_\psi^2 \frac{\tilde{n}_\psi}{m_\psi}\,, \tag{2.8}$$

where

$$\tilde{n}_\psi \equiv \int \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{m_\psi}{E_\mathbf{p}} f(\mathbf{p})\,, \tag{2.9}$$

with $f(\mathbf{p})$ the phase space distribution function of ψ particles. We refer to $\tilde{n}_\psi$ in eq. (2.9) as *the effective number density*, in contrast to the standard number density:

$$n_\psi \equiv \int \frac{d^3\mathbf{p}}{(2\pi)^3} f(\mathbf{p})\,. \tag{2.10}$$

Eq. (2.8) can be derived from finite-temperature or finite-density field theory [40]. However, such a derivation involves the underlying assumption that the medium is in thermal and chemical equilibrium, which implies that $f(\mathbf{p})$ should take the Fermi-Dirac distribution. In fact, eq. (2.8) also holds for general forms of $f(\mathbf{p})$, not necessarily limited to thermal distributions, since it is a consequence of coherent forward scattering of ϕ with the medium particles. This is again similar to the photon, for which the plasmon mass and the refraction index are known to arise from coherent forward scattering. In appendix B, we rederive eq. (2.8) from the theory of coherent forward scattering, following similar calculations in refs. [42, 43] for photon-dark photon or photon-axion conversions. This justifies the use of eq. (2.8) for non-thermal distributions such as the CνB.

Let us note here that the mass correction in eq. (2.8) is a local rather than global effect. In a density-varying medium, one should only apply the mass correction at the EOM level, not the Yukawa potential in eq. (2.7) which is obtained by assuming constant m_ϕ. The actual potential should be determined by solving the modified EOM:

$$\left[\nabla^2 - m_\phi^2 - y_\psi^2 \frac{\tilde{n}_\psi}{m_\psi}\right] \phi = y_\psi \tilde{n}_\psi\,. \tag{2.11}$$

Note that comparing to eq. (2.5), the right-hand side has also been changed. The right-hand side is changed from $y_\psi n_\psi$ to $y_\psi \tilde{n}_\psi$ because $\langle\overline{\psi}\psi\rangle \approx n_\psi$ is only applicable to nonrelativistic distributions while $\langle\overline{\psi}\psi\rangle = \tilde{n}_\psi$ holds for general cases including relativistic distributions [41].

By definition, n_ψ and $\tilde{n}_\psi$ in eqs. (2.9) and (2.10) do not include antiparticles. For antiparticles, we define similar quantities, $n_{\overline{\psi}}$ and $\tilde{n}_{\overline{\psi}}$. When including antiparticles, there are

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a few subtle minus signs to be discussed in appendix A and the conclusion is that antiparticles contribute positively to eq. (2.8), i.e. there is no cancellation between particles and antiparticles in the medium effect. Therefore, the contribution of antiparticles can be included by

$$\tilde{n}_\psi \to \tilde{n}_\psi + \tilde{n}_{\overline{\psi}}\,. \tag{2.12}$$

In addition, eq. (2.11) can be straightforwardly generalized to include multiple fermionic species that are coupled to ϕ. Altogether, the most general equation that determines the ϕ field in the presence of multiple species of medium particles together with antiparticles reads:

$$\left[\nabla^2 - \tilde{m}_\phi^2\right]\phi = \sum_\psi y_\psi \left(\tilde{n}_\psi + \tilde{n}_{\overline{\psi}}\right), \tag{2.13}$$

where

$$\tilde{m}_\phi^2 \equiv m_\phi^2 + \sum_\psi y_\psi^2 \frac{\tilde{n}_\psi + \tilde{n}_{\overline{\psi}}}{m_\psi}\,. \tag{2.14}$$

2.3 On the effective number density $\tilde{n}_\psi$

In the nonrelativistic limit, the effective number density $\tilde{n}_\psi$ defined in eq. (2.9) simply reduces to n_ψ because $m_\psi/E_{\mathbf{p}} \approx 1$. For relativistic distributions, $\tilde{n}_\psi$ is smaller than n_ψ and their ratio $\tilde{n}_\psi/n_\psi$ is roughly suppressed by $m_\psi/\langle E_{\mathbf{p}}\rangle$ where $\langle E_{\mathbf{p}}\rangle$ denotes the mean energy of ψ particles.

In general, by computing the integral in eq. (2.9) for a given $f(\mathbf{p})$, one obtains

$$\tilde{n}_\psi \approx \begin{cases} n_\psi & \text{(nonrelativistic)} \\ \xi m_\psi n_\psi^{2/3} & \text{(relativistic)} \end{cases}, \tag{2.15}$$

where ξ is an $\mathcal{O}(1)$ coefficient depending on the shape of $f(\mathbf{p})$. For the Fermi-Dirac (FD) distribution with zero chemical potential,

$$\xi = \frac{1}{6}\left(\frac{\pi^2}{6\zeta(3)}\right)^{2/3} g_i^{1/3} \approx 0.2054 g_i^{1/3}\,, \tag{2.16}$$

with g_i the number of internal degrees of freedom: $g_i = 2$ for electrons and $g_i = 1$ for Majorana neutrinos. For the FD distribution with zero temperature but a finite chemical potential (i.e. the degenerate limit), $\xi \approx 0.3848 g_i^{1/3}$. The Maxwell-Boltzmann distribution would lead to $\xi \approx 0.2331 g_i^{1/3}$.

Note that the standard cosmic neutrino background (CνB) with nonzero neutrino masses is not exactly in the FD distribution (see e.g. [44] for detailed discussions):

$$f_{\mathrm{C}\nu\mathrm{B}}(p) = \frac{1}{e^{p/T}+1} \neq \frac{1}{e^{E_p/T}+1}\,. \tag{2.17}$$

But for neutrino masses well below 1.9 K $\approx 1.6 \times 10^{-4}$ eV (possible for the lightest neutrino mass eigenstate), the difference is negligible and one can take $\xi \approx 0.2054 g_i^{1/3}$.

The neutrino mass squared differences in neutrino oscillations have been well measured [45]:

$$|\Delta m_{31}^2| \approx 2.5 \times 10^{-3}\ \mathrm{eV}^2\,,\ \Delta m_{21}^2 \approx 7.4 \times 10^{-5}\ \mathrm{eV}^2\,, \tag{2.18}$$

which implies that two of the mass eigenstates are heavier than $\sqrt{|\Delta m_{31}^2|} \approx 0.05\,\mathrm{eV}$ or $\sqrt{\Delta m_{21}^2} \approx 0.009\,\mathrm{eV}$. Therefore, among three neutrino mass eigenstates, at least two of them should be nonrelativistic today.

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3 The spherically symmetric solution

The differential equation (2.11) can be solved analytically if the effective number density is spherically symmetric with the following form:

$$\tilde{n}_\psi(r) = \begin{cases} \tilde{n}_{\psi 0} & \text{for } r \le R \\ 0 & \text{for } r > R \end{cases}, \tag{3.1}$$

where $\tilde{n}_{\psi 0}$ is a constant and R is the radius of the spherical distribution. By solving eq. (2.11) in the $r \le R$ and $r > R$ regimes and requiring that ϕ and $d\phi/dr$ are continuous at $r = R$, we obtain the solution

$$\phi = \begin{cases} \frac{y_\psi \tilde{n}_{\psi 0}}{\tilde{m}_\phi^2}\left[-1 + \frac{C_{\rm in}}{r}\sinh(\tilde{m}_\phi r)\right] & \text{for } r \le R \\ \frac{y_\psi \tilde{n}_{\psi 0}}{\tilde{m}_\phi^2} \cdot \frac{C_{\rm out}}{r} e^{-m_\phi (r-R)} & \text{for } r > R \end{cases}, \tag{3.2}$$

where $\tilde{m}_\phi$ denotes the effective mass within the sphere ($\tilde{m}_\phi^2 = m_\phi^2 + y_\psi^2 \tilde{n}_{\psi 0}/m_\psi$) and

$$C_{\rm in} = \frac{m_\phi R + 1}{\tilde{m}_\phi \cosh(\tilde{m}_\phi R) + m_\phi \sinh(\tilde{m}_\phi R)}, \tag{3.3}$$

$$C_{\rm out} = \frac{\sinh(\tilde{m}_\phi R) - \tilde{m}_\phi R \cosh(\tilde{m}_\phi R)}{m_\phi \sinh(\tilde{m}_\phi R) + \tilde{m}_\phi \cosh(\tilde{m}_\phi R)}. \tag{3.4}$$

Here we would like to discuss an interesting limit: $m_\phi \to 0$, i.e. ϕ is massless in the vacuum but acquires a mass within the sphere due to the medium. In this limit, eq. (3.2) reduces to

$$\phi = \begin{cases} -\frac{y_\psi \tilde{n}_{\psi 0}}{\tilde{m}_\phi^2}\left[1 - \frac{\sinh(\tilde{m}_\phi r)}{\tilde{m}_\phi r \cosh(\tilde{m}_\phi R)}\right] & \text{for } r \le R \\ -\frac{y_\psi \tilde{n}_{\psi 0}}{\tilde{m}_\phi^2}\left[1 - \frac{\tanh(\tilde{m}_\phi R)}{\tilde{m}_\phi R}\right]\frac{R}{r} & \text{for } r > R \end{cases}, \quad (m_\phi \to 0). \tag{3.5}$$

Obviously, the $r > R$ part of the solution behaves like an infinitely long-range force, as can be seen from its $1/r$ dependence. If the medium effect is weak ($\tilde{m}_\phi R \ll 1$), one can further expand it in terms of $x \equiv \tilde{m}_\phi R$:

$$\phi = \begin{cases} y_\psi \tilde{n}_{\psi 0}\left[\frac{r^2 - 3R^2}{6} + \frac{x^2\left(r^2 - 5R^2\right)^2}{120 R^2} + \mathcal{O}(x^4)\right] & \text{for } r \le R \\ -\frac{y_\psi \tilde{n}_{\psi 0} R^3}{3r}\left[1 - \frac{2x^2}{5} + \mathcal{O}(x^4)\right] & \text{for } r > R \end{cases}, \quad (m_\phi \to 0). \tag{3.6}$$

Here the first terms in squared brackets correspond to known results of the Coulomb potential, while the second terms represent the leading-order medium effect.

In the presence of multiple species of medium particles coupled to ϕ, one needs to solve the more general differential equation (2.13). Note that due to the mass correction term in eq. (2.14), eq. (2.13) cannot be solved by linearly adding the solutions in eq. (3.2) for each species. Unlike eq. (2.5) for which the solution scales linearly with respect to the number density, eqs. (2.11) and (2.13) exhibit nonlinearity when the medium effect is significant.

Nevertheless, eq. (3.2) can still be practically useful under certain approximations and assumptions. Consider for instance the Earth as a spherically symmetric object. Assuming that its matter density and chemical composition are homogeneous, the above solution can be

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used by replacing $y_\psi \tilde{n}_{\psi 0} \to \sum_\psi y_\psi (\tilde{n}_{\psi 0} + \tilde{n}_{\overline{\psi}0})$ in eq. (3.2) and taking $\tilde{m}_\phi^2 = m_\phi^2 + \sum_\psi y_\psi^2 (\tilde{n}_{\psi 0} + \tilde{n}_{\overline{\psi}0})/m_\psi$, where the summation $\sum_\psi$ goes over all ingredients of the Earth. Furthermore, adding the CνB to this problem can also be solved analytically, assuming that the CνB is homogeneous and extends infinitely in space. Under this assumption, the medium effect of CνB is a constant shift added to m_ϕ^2.

Another example is the Yukawa force between two celestial bodies. Assuming that each celestial body is a spherically symmetric object with a homogeneous matter density and the two bodies are well separated, one can first solve their respective ϕ field equations:

$$\left[\nabla^2 - \tilde{m}_{\phi 1}^2\right] \phi_1 = y_\psi \tilde{n}_{\psi 1} \, , \tag{3.7}$$

$$\left[\nabla^2 - \tilde{m}_{\phi 2}^2\right] \phi_2 = y_\psi \tilde{n}_{\psi 2} \, , \tag{3.8}$$

where ϕ_i ($i = 1,\, 2$) denotes the field strength generated by the i-th celestial body and $\tilde{m}_{\phi i}^2 = m_\phi^2 + y_\psi^2 \tilde{n}_{\psi i}/m_\psi$. Here for simplicity we assume that they all consist of a single species of ψ particles. Summing the two solutions ϕ_1 and ϕ_2 together cannot generate an exact solution of the combined system, but $\phi = \phi_1 + \phi_2$ satisfies

$$\left[\nabla^2 - \tilde{m}_\phi^2\right] \phi = y_\psi \tilde{n}_{\psi 1} + y_\psi \tilde{n}_{\psi 2} + y_\psi^2 \frac{\tilde{n}_{\psi 2}\phi_1 + \tilde{n}_{\psi 1}\phi_2}{m_\psi} \, , \tag{3.9}$$

where $\tilde{m}_\phi^2 = m_\phi^2 + y_\psi^2(\tilde{n}_{\psi 1} + \tilde{n}_{\psi 2})/m_\psi$ and the last term in eq. (3.9) comes from cross terms in $\tilde{m}_\phi^2 \phi$. Compared to $y_\psi \tilde{n}_{\psi 1}$ or $y_\psi \tilde{n}_{\psi 2}$, the last term is suppressed by a factor of $y_\psi \phi_i / m_\psi$ which is $\ll 1$ as long as the mass shift of ψ [see the discussion above eq. (2.4)] is small. Hence $\phi = \phi_1 + \phi_2$ can be viewed as an approximate solution of the field equation for the two-body system.

4 Altering the experimental bounds on long-range forces

A variety of experiments measuring gravitational effects or testing gravitational laws can also be utilized to search for extra long-range forces, as has been conducted by a number of experimental groups [27, 46–53]. In recent years, the detection of gravitational waves from black hole or neutron star binary mergers also offers an important avenue for probing long-range forces [35–37]. We focus our analyses on experimental bounds derived from measuring gravitational effects of the most well-understood celestial bodies nearby, namely the Sun, Earth, and Moon. With these celestial bodies as attractors, new long-range forces can be severely constrained by measuring possible acceleration differences between two test bodies composed of different materials (e.g. Be, Ti, Al, Pt), known as the test of the weak equivalence principle (WEP). Such a test has been conducted by the Princeton [54] and Moscow [55] groups utilizing the Sun as the attractor and subsequently by the Eöt-Wash group [27, 47] using the Earth as the attractor. In addition to the WEP test, there is another interesting type of experimental probe — the lunar laser-ranging (LLR) experiments [28–32] which are capable to measure anomalous precession of the lunar orbit and hence sensitive to new forces that modify the inverse square law. Very recently, the MICROSCOPE experiment [38] deployed test masses orbiting the Earth in a drag-free satellite and achieved an unprecedented precision of measuring WEP violation.

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In the literature, these measurements have been used to use to set constraints on the standard Yukawa potential in eq. (2.7). Typically for m_ϕ ranging from 10^{-14} eV (corresponding to the length scale of the Earth radius) to 10^{-18} eV (corresponding to the length scale of Earth-Sun distance), the experimental bounds on y_ψ for $\psi \in \{e^-,\ p,\ n\}$ vary from 10^{-22} to 10^{-24} — see the top left panel of figure 1. Let us first estimate the medium effect caused by ϕ interacting with electrons in the medium, which leads to a mass correction of the order of

$$y_e \left(\frac{n_e}{m_e}\right)^{1/2} \sim 10^{-22}\ \text{eV} \cdot \left(\frac{y_e}{10^{-24}}\right) \cdot \left(\frac{\rho}{5\ \text{g/cm}^3}\right)^{1/2}, \tag{4.1}$$

where ρ denotes the matter density and 5 g/cm^3 is the average density of the Earth. In eq. (4.1) we have assumed $n_e = n_p \approx n_n$ so that $n_e \approx 0.5\rho/m_p$. Obviously this medium effect is negligibly small for the experiments considered here.

Next, we estimate the medium effect caused by the CνB which is ubiquitous in the entire universe and according to the standard cosmological model has the number density $56\ \nu/\text{cm}^3 + 56\ \overline{\nu}/\text{cm}^3$ for Dirac neutrinos or $112\ \nu/\text{cm}^3$ for Majorana neutrinos. The medium effect due to the CνB can be substantially enhanced by the smallness of neutrino masses. Besides, the Yukawa coupling y_ν for very light ϕ can be much greater than y_e as known bounds on y_ν are much less restrictive than those on y_e. Theoretically, $y_\nu \gg y_e$ can be well motivated from models that introduce light mediators via the neutrino portal [3, 4, 6].

For nonrelativistic cosmic relic neutrinos ($m_\nu \gg 1.6 \times 10^{-4}$ eV), the mass correction can be estimated as follows:

$$y_\nu \left(\frac{n_\nu}{m_\nu}\right)^{1/2} \sim 10^{-16}\ \text{eV} \cdot \left(\frac{y_\nu}{10^{-10}}\right) \cdot \left(\frac{n_\nu}{56/\text{cm}^3}\right)^{1/2} \left(\frac{0.1\text{eV}}{m_\nu}\right)^{1/2}. \tag{4.2}$$

For relativistic neutrinos, according to eq. (2.15), the mass correction should be independent of m_ν:

$$y_\nu \left(\frac{\tilde{n}_\nu}{m_\nu}\right)^{1/2} \sim 10^{-15}\ \text{eV} \cdot \left(\frac{y_\nu}{10^{-10}}\right) \cdot \left(\frac{n_\nu}{56/\text{cm}^3}\right)^{1/3}. \tag{4.3}$$

The mass corrections of 10^{-15} or 10^{-16} eV in eqs. (4.2) and (4.3) are comparable to the inverse of the geometrical sizes of the celestial bodies being considered. Hence the above estimates imply that the CνB medium effect can be considerably large in the aforementioned experiments.

Indeed, as is shown in figure 1, for a rather broad range of y_ν, the CνB may significantly alter some of the experimental bounds on long-range forces. In this figure, we select three representative experiments, the Moscow experiment using the Sun as the attractor (green curves), the Eöt-Wash experiment using the Earth as the attractor (blue), and LLR which measures the varying Earth-Moon distance (orange). The experimental bounds without including medium effects are taken from ref. [27]. We recast the bounds into those presented in figure 1 with significant CνB medium effects included. Since the CνB is a homogeneous distribution, the main effect is the mass correction in eq. (2.14). The CνB could also contribute to the source term on the right-hand side of eq. (2.13), but due to the homogeneity this contribution only causes an constant shift of ϕ, corresponding to adding a constant to the effective potential. Hence it can be omitted.

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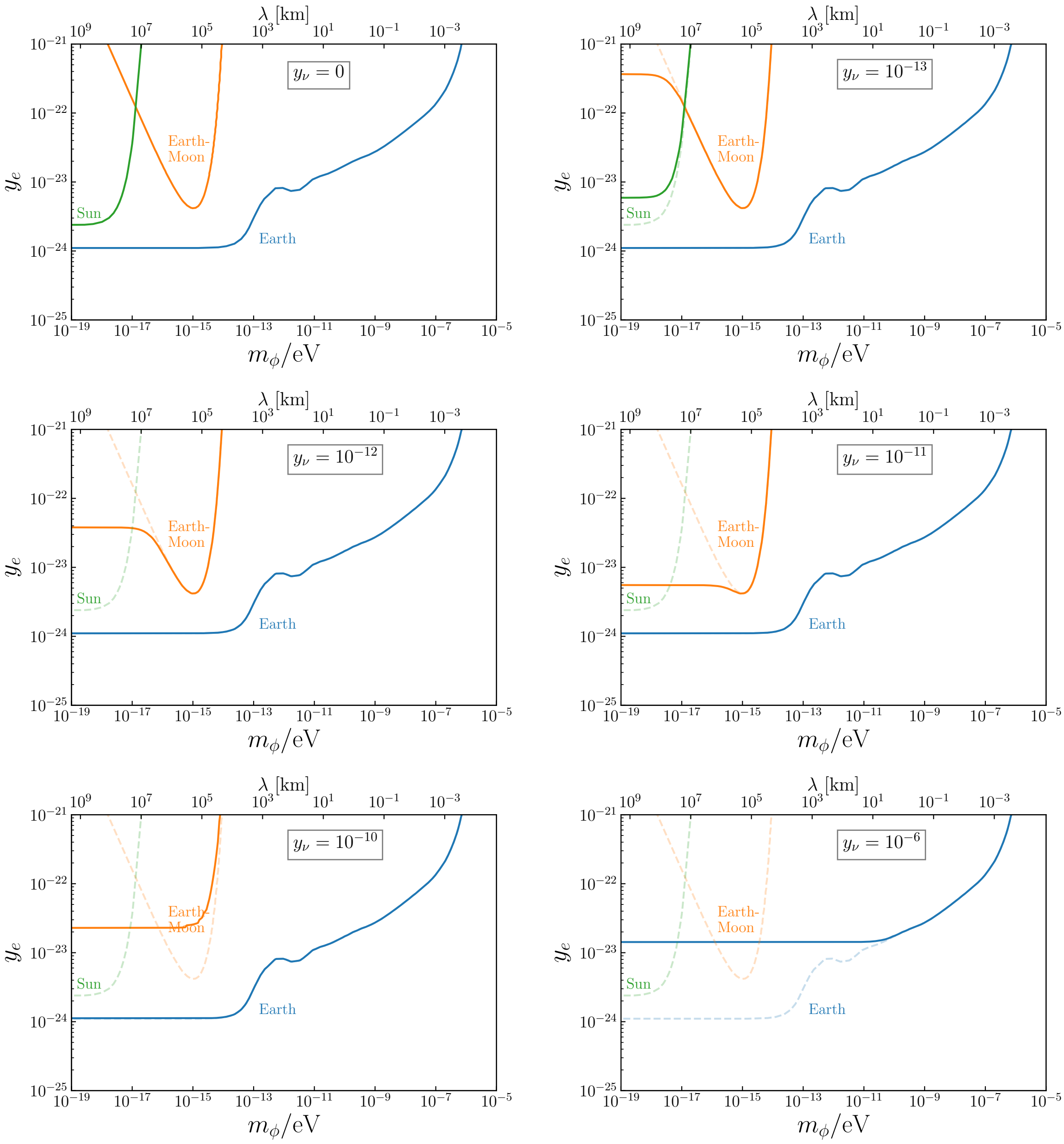

Figure 1. The impact of the CνB medium effect on long-range force searches. The top left panel shows the experimental bounds on y_e without the CνB medium effect. The subsequent panels demonstrate the variation of these bounds due to the CνB medium effect, with the dashed and solid lines representing the bounds before and after including such an effect, respectively. The x-axis on the top of each panel indicates the wavelength scale $\lambda \equiv 1/m_\phi$.

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More specifically, we recast the bounds as follows. For each given y_ν and m_ϕ, we can compute the effective mass $\tilde{m}_\phi$ according to eq. (2.14), assuming that the mass corrections due to y_ψ with $\psi \in \{e^-, \ p, \ n\}$ are negligible, which has been justified by eq. (4.1). Since the mass correction caused by y_ν is independent of r, the new experimental bounds on y_ψ including the medium effect can be obtained by matching $y_\psi^{(\text{new})}(m_\phi) = y_\psi^{(\text{old})}(\tilde{m}_\phi)$, where $y_\psi^{(\text{new})}(m_\phi)$ and $y_\psi^{(\text{old})}(m_\phi)$ denote the new and old bounds with and without the medium effect included, respectively. Note that this is valid only when the effective mass is independent of r. If the CνB is not homogeneous at local scales or if the mass corrections due to y_ψ with $\psi \in \{e^-, \ p, \ n\}$ becomes significant, we should numerically solve eq. (2.13) to obtain the potential, with the geometry of the actual matter distribution taken into account, and confront it with the experimental data to obtain the actual bound.

In figure 1, we gradually increase the coupling y_ν from 0 to 10^{-6} to show the impact of the CνB medium effect on long-range force searches, assuming relativistic CνB. For nonrelativistic CνB, the results are similar and the difference can be absorbed by rescaling y_ν as follows:

$$y_\nu \to 18.5 \times \left(\frac{m_\nu}{0.1 \text{ eV}}\right)^{1/2} y_\nu \, . \tag{4.4}$$

In the upper right panel, we set $y_\nu = 10^{-13}$ which according to eq. (4.3) leads to a mass correction of 10^{-18} eV, corresponding to the inverse of the Earth-Sun distance. Consequently, the green curve in the top right panel is significantly altered and becomes a weaker bound. It is interesting to note that the LLR bound (orange curve) in this panel becomes more constraining for $m_\phi \lesssim 10^{-18}$ eV. This is because the LLR bound is essentially a constraint on deviations from the inverse square law. Without the medium effect, sufficiently light mediators cannot be effectively constrained by LLR since the new force in the massless limit also follows the inverse square law. With the medium effect, the effective mass is dominated by the medium-induced mass if the vacuum mass m_ϕ is small. So LLR sets a bound on y_e independent of m_ϕ in this regime.

Further increasing y_ν, some bounds can be substantially alleviated or even entirely eliminated. For example, for $y_\nu \gtrsim 5 \times 10^{-13}$ and 3×10^{-10}, the green and orange bounds would disappear respectively, as the medium effect renders the forces short-range with respect to the relevant distances. One may ask to what extent these long-range force bounds can be alleviated by continuing increasing y_ν. In fact, when y_ν is too large, the ϕ field would thermalize via neutrino scattering in the early universe, modifying the BBN observables and the cosmological effective number of relativistic neutrino species, N_{eff}. According to refs. [56, 57], for light scalars predominantly coupled to neutrinos, the cosmological upper bound on y_ν is about 10^{-5}. Therefore, we stop increasing y_ν at 10^{-6} in figure 1. In addition to the cosmological bound, there are other bounds on y_ν from constraints on neutrino self-interactions [58–64]. Typically laboratory bounds on y_ν can reach 10^{-3} to 10^{-5}.

From figure 1, we conclude that the CνB medium effect on long-range force searches can fully eliminate the bounds derived from LLR and the Sun, and alleviate the bound from the Earth by about one order of magnitude.

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5 Conclusions

Long-range force searches generally set the most constraining bounds on ultra-light bosons that interact with normal matter. In this work, we assume that such bosons are also coupled to neutrinos and investigate whether the presence of cosmic neutrino background (CνB) could have an impact on long-range force search experiments which are soaked in this ubiquitous background. We find that the medium effect of CνB can significantly shorten the interaction ranges of long-range forces, due to coherent forwarding scattering of the force mediator with neutrinos. Consequently, the experimental bounds on long-range forces can be altered by the CνB with sizable couplings of the mediator to neutrinos.

Specifically, for experiments sensitive to new forces with interaction ranges longer than $\sim 10^8$ km (the Sun-Earth distance, corresponding to a mediator mass of $\sim 10^{-18}$ eV) or $\sim 10^5$ km (the Earth-Moon distance), the CνB could effectively screen the forces from detection with the neutrino coupling greater than 5×10^{-13} or 3×10^{-10}, respectively, causing such experimental bounds to be entirely eliminated. For experiments utilizing the Earth as an attractor, the bounds can be substantially alleviated by about one order of magnitude with cosmologically allowed values of the neutrino coupling.

Our results can also be used in the studies of neutrino oscillations with long-range forces — see e.g. [39, 40, 65–70]. Besides, future gravitational wave experiments such as LIGO, VIRGO, Einstein Telescope, LISA, NANOGrav, etc., will be able to probe a multitude of long-range force effects in both astrophysical and terrestrial environments. For instance, neutron star mergers may provide a novel avenue to long-range force searches, since the observed gravitational waves could be altered by extra forces between two neutron stars as well as extra radiations [36, 37]. The high density of a neutron star could modify such forces due to the medium effect discussed in this work. On the other hand, it is also possible that the extra radiations are enhanced or suppressed by the mass correction caused by the surrounding cosmic neutrino background. We leave dedicated studies on these possibilities for future work.

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A Some minus signs

In this appendix, we shall clarify a few minus signs important to our analysis.

A.1 Minus signs responsible for repulsive and attractive forces

If the scalar mediator ϕ is replaced by a vector mediator ϕ^μ, the long-range force between two ψ particles becomes repulsive instead of attractive. This change arises essentially from the differences in their kinetic and mass terms. The mass term of a vector boson is positive, implying a sign flip in the mass term:

$$-\frac{1}{2}m_\phi^2\phi^2 \;\to\; +\frac{1}{2}m_\phi^2\phi^\mu\phi_\mu\,. \tag{A.1}$$

On the other hand, the kinetic terms also differ by a minus sign:

$$\frac{1}{2}(\partial\phi)^2 = -\frac{1}{2}\phi\partial^2\phi \;\to\; -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} = \frac{1}{2}\phi^\mu\partial^2\phi_\mu\,, \tag{A.2}$$

where $F^{\mu\nu} \equiv \partial^\mu\phi^\nu - \partial^\nu\phi^\mu$. Consequently, the EOM in eq. (2.3) changes to

$$-(\partial^2 + m_\phi^2)\phi^\mu + y_\psi\overline{\psi}\gamma^\mu\psi = 0\,. \tag{A.3}$$

Due to the minus sign in eq. (A.3), the force becomes repulsive between two ψ particles. One may wonder if the sign of y_ψ in the Lagrangian (2.1) is relevant. This sign is unimportant because the effective potential and the mass correction due to the medium effect are all proportional to y_ψ^2. So flipping the sign of y_ψ has no physical consequences.

A.2 When do antiparticle and particle contributions cancel?

Antiparticles may introduce some extra minus signs to the calculations. In eqs. (2.5), (2.8) and (2.11), only particles (no antiparticles) are included. In the presence of antiparticles, one may be concerned with possible cancellations between antiparticles and particles.

Let us first inspect the role of antiparticles in the ensemble averages $\langle\overline{\psi}\psi\rangle$ and $\langle\overline{\psi}\gamma^\mu\psi\rangle$, which determine the scalar and vector field strengths of ϕ (ϕ^μ), respectively. Under charge conjugation (the C transformation), we have

$$\langle\overline{\psi}\gamma^\mu\psi\rangle \overset{C}{\to} -\langle\overline{\psi}\gamma^\mu\psi\rangle\,, \tag{A.4}$$

$$\langle\overline{\psi}\psi\rangle \overset{C}{\to} \langle\overline{\psi}\psi\rangle\,, \tag{A.5}$$

which implies that $\langle\overline{\psi}\gamma^\mu\psi\rangle$ and $\langle\overline{\psi}\psi\rangle$ should be antisymmetric and symmetric under the particle-antiparticle interchange:

$$\langle\overline{\psi}\gamma^\mu\psi\rangle = \int\frac{d^3\mathbf{p}}{(2\pi)^3}\frac{p^\mu}{E_\mathbf{p}}\left[f_+(\mathbf{p}) - f_-(\mathbf{p})\right], \tag{A.6}$$

$$\langle\overline{\psi}\psi\rangle = \int\frac{d^3\mathbf{p}}{(2\pi)^3}\frac{m_\psi}{E_\mathbf{p}}\left[f_+(\mathbf{p}) + f_-(\mathbf{p})\right], \tag{A.7}$$

where f_+ and f_- denote the phase space distribution functions of ψ particles and antiparticles, respectively. Note that here $\frac{d^3\mathbf{p}}{(2\pi)^3}\frac{1}{E_\mathbf{p}}$ is Lorentz invariant.

Eqs. (A.6) and (A.7) are obtained simply from the argument of the C symmetry. They can also be obtained by explicitly computing the ensemble averages (i.e. performing Wick contractions of creation and annihilation operators) using the following identities:

$$\overline{u}\gamma^\mu u = 2p^\mu\,, \qquad \overline{v}\gamma^\mu v = 2p^\mu\,, \tag{A.8}$$

$$\overline{u}u = 2m_\psi\,, \qquad \overline{v}v = -2m_\psi\,, \tag{A.9}$$

where u and v denote the particle and antiparticle solutions of the Dirac equation in momentum space, i.e. $(\not{p} - m_\psi)u = 0$ and $(\not{p} + m_\psi)v = 0$. Note that the $u \leftrightarrow v$ interchange causes

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a minus sign for the scalar case in eq. (A.9), as opposed to eq. (A.5) or (A.7) where the scalar product is symmetric under the particle-antiparticle interchange. This minus sign is canceled by interchanging anti-commuting fermionic creation and annihilation operators. For the vector case, the same reason accounts for the anti-symmetry in eq. (A.6) when deriving it from eq. (A.8).

Finally, let us inspect the role of antiparticles in the mass correction (2.8). The medium-induced mass correction for the scalar case arises from the evaluation of the following ensemble average:

$$\Delta m_\phi^2 \phi^2 \propto y_\psi^2 \langle \overline{\psi}\phi\psi\overline{\psi}\phi\psi \rangle \,, \tag{A.10}$$

which is obviously a C-even quantity according to eq. (A.5). Hence antiparticles contribute positively to eq. (2.8), which after including such a contribution should be

$$m_\phi^2 \to m_\phi^2 + y_\psi^2 \frac{\tilde{n}_\psi + \tilde{n}_{\overline{\psi}}}{m_\psi} \,. \tag{A.11}$$

As for the vector case, although $\overline{\psi}\gamma^\mu\psi$ is C-odd, the medium-induced mass correction $\langle \overline{\psi}\gamma\cdot\phi\psi\overline{\psi}\gamma\cdot\phi\psi \rangle$ is C-even. Hence the antiparticle contribution is also positive, which is a known conclusion for the plasmon mass in e^+e^- plasma.

B Derivation of the medium effect from coherent forward scattering

As mentioned in section 2.2, the medium effect given by eqs. (2.8) and (2.9), known from the finite-temperature field theory, is also valid for non-thermal distributions, because it fundamentally arises from coherent forward scattering of ϕ with medium particles. In this appendix, we rederive this medium effect using the theory of coherent forward scattering — similar calculations for photon-dark photon or photon-axion conversions can be found in refs. [42, 43].

In coherent forward scattering, the medium particles do not change their states after such scattering, implying that the amplitude of scattering with each medium particle can be added coherently, as is illustrated by figure 2. The combined effect corresponds to the mass correction in eq. (2.8).

Let us first compute the scattering of ϕ with a single ψ particle, for which the amplitude reads

$$i\mathcal{M} = y_\psi^2 \overline{u}(p) \left[\frac{i}{(p+k)\cdot\gamma - m_\psi} + \frac{i}{(p-k)\cdot\gamma - m_\psi} \right] u(p) \,, \tag{B.1}$$

where u and $\overline{u}$ denote the initial and final fermion states; p and k denote the momenta of the fermion and the scalar, respectively. The second term in square brackets accounts for the interchange of the two Yukawa vertices. Applying on-shell conditions for the fermion, $p^2 = m_\psi^2$ and $\not{p}u = m_\psi u$, we obtain

$$i\mathcal{M} = i y_\psi^2 \overline{u} \left[\frac{4k^2 m_\psi - 4(p\cdot k)(k\cdot\gamma)}{k^4 - 4(k\cdot p)^2} \right] u \,. \tag{B.2}$$

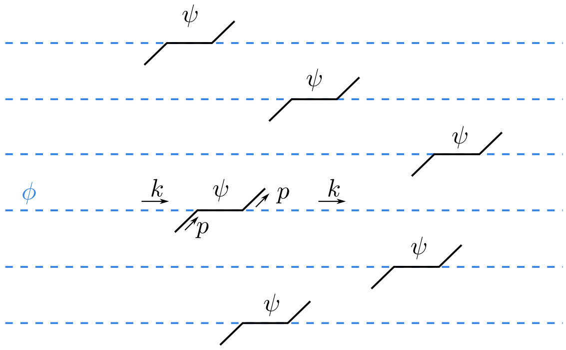

Figure 2. A schematic picture illustrating coherent forward scattering of ϕ with medium particles. The momenta of both ϕ and ψ remain unchanged after the scattering. As a consequence, all scattering amplitudes can be added coherently. The combined medium effect slows down the propagation of the ϕ field.

Next, we consider such scattering in a medium containing a large number of ψ particles, which form a background state denoted by $|\text{bkg}\rangle$. In this case, we need to evaluate the following matrix element:

$$\langle \text{bkg}|\overline{\psi}\phi\psi\overline{\psi}\phi\psi|\text{bkg}\rangle = \langle \text{bkg}|\overline{\psi}\phi\psi\overline{\psi}\phi\psi|\text{bkg}\rangle + \text{vertex interchange}\,, \tag{B.3}$$

where the first and third Wick contractions correspond to $\overline{u}$ and u in eq. (B.1), and the second contraction corresponds to the fermion propagator in eq. (B.1), respectively. The "vertex interchange" term corresponds to the second term in eq. (B.1). Therefore, eq. (B.2) implies

$$\begin{aligned}\langle \text{bkg}|\overline{\psi}\phi\psi\overline{\psi}\phi\psi|\text{bkg}\rangle &= \phi^2 \langle \text{bkg}|\overline{\psi}\, \frac{4k^2 m_\psi - 4(p\cdot k)(k\cdot\gamma)}{k^4 - 4(k\cdot p)^2}\, \psi|\text{bkg}\rangle \\ &= \phi^2 \int \frac{d^3\mathbf{p}}{(2\pi)^3}\frac{f_+(\mathbf{p})}{E_\mathbf{p}}\frac{4k^2 m_\psi^2 - 4(p\cdot k)(k_\mu p^\mu)}{k^4 - 4(k\cdot p)^2}\,,\end{aligned} \tag{B.4}$$

where in the second step we have used eqs. (A.6) and (A.7). For simplicity, we assume that the background does not contain antiparticles. The antiparticle contribution can be readily included by performing charge conjugation on the final result.

For applications in generating effective potentials, one can take the limit $k^2 \to 0$. Then eq. (B.4) implies

$$y_\psi^2 \langle \text{bkg}|\overline{\psi}\phi\psi\overline{\psi}\phi\psi|\text{bkg}\rangle = y_\psi^2 \phi^2 \int \frac{d^3\mathbf{p}}{(2\pi)^3}\frac{f_+(\mathbf{p})}{E_\mathbf{p}} = y_\psi^2\phi^2 \frac{\tilde{n}_\psi}{m_\psi}\,, \tag{B.5}$$

which is exactly the mass correction in eq. (2.8).

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