

Interference of colliding Fock states and vacuum-generated photon pairs at time interfaces

Nikita Nefedkin^{1,*}, Emanuele Galiffi^{1,†,‡} and Andrea Alù^{1,2,§}¹Photonics Initiative, Advanced Science Research Center, City University of New York, New York, New York 10031 USA²Physics Program, Graduate Center, City University of New York, New York, New York 10016, USA

(Received 2 April 2025; accepted 11 February 2026; published 23 March 2026)

We investigate the scattering of counterpropagating photon Fock states at a time interface, induced by an abrupt variation of the dielectric permittivity of the host material. We unveil the role of vacuum-generated photons in determining the final photon number distribution and coherence properties. Such vacuum-generated photons interfere with the photons in the initial Fock states, producing distinctive interference patterns, which can be broadly controlled, together with the temperature of the final photon gas, by the time-interface parameters and the number of photons in the initial states.

DOI: [10.1103/z61w-gdvl](https://doi.org/10.1103/z61w-gdvl)

Time interfaces (TIs) are abrupt variations in the optical properties of a material, which can produce efficient, broadband frequency translation and time reflection [1,2]. Recently, experimental demonstrations of time interfaces for water waves [3], microwaves [4], and cold atoms [5] have been presented, alongside progress in near-single-cycle nonlinear nanophotonics [6–8]. These advances have stimulated significant interest in the engineering of wave scattering via temporal inhomogeneities, launching the paradigm for four-dimensional metamaterials [9].

Time-varying media hold the potential to unleash a host of opportunities beyond the current boundaries of wave manipulation, from surpassing the conventional bandwidth limitations of optical devices [10–12] to novel schemes for nonreciprocal propagation [13] and gain [14], radiative heat transfer [15–17], wave routing [18], amplification [19], coherent absorption [20], and wave-based computation [21] schemes, as well as photonic analogs of gravitational effects such as Hawking radiation [22,23], among others [24].

Time-varying systems have been extensively studied also in the condensed matter community, where the search for time crystals [25–27] has stimulated intense research in Floquet engineering [28] amid the search for new phases of matter, with recent experimental realizations in both quantum [29] and classical settings [30]. Quantum optics has been an active platform for dynamical photonic systems, e.g., in the context of the dynamical Casimir effect, i.e., the creation of photon pairs from the interaction of the vacuum field state with an oscillating mirror [31–36]. Similar, although nonequivalent

phenomena emerge in materials whose dielectric permittivity is temporally modulated [35,37–40]. In particular, photon-pair creation, which may be triggered by both moving media and time-modulated structures, was explored in detail for vacuum amplification in the continuous modulation regime [41], as well as for quantum state engineering, also in the presence of a single-input channel populated by Fock states [32,42].

Recently, however, the potential of interferometric control of light in time-varying media was identified [20,43]. In these schemes, the system is illuminated by multiple light beams, which undergo mutual interactions as a result of the temporal variation of the background material. This framework can be used to shape the output light in a linear fashion, coherently controlling the output via the relative phase of the incoming waves, leveraging the temporal analog of coherent perfect absorption [44], but benefiting from gain, loss, and broadband features inherent to time-varying media [20,43].

Here, we explore the application of coherent light control to quantum light by demonstrating how the action of a time interface on a pair of counterpropagating Fock states conspires with the photon-pair production resulting from the effect of the time interface on the vacuum state, i.e., the dynamical Casimir effect [45], synthesizing complex quantum states of light. The statistics of the final photon distribution emerging from the time interface, significantly departing from a thermal state, can be tailored in this process. In turn, these phenomena open new opportunities to exploit broadband dynamical Casimir physics emerging at time interfaces to engineer complex quantum states of light.

We start by studying the effect of an abrupt permittivity variation of a lossless, homogeneous, and nondispersive host medium $\epsilon_j \rightarrow \epsilon_{j+1}$ at instant t_j . We excite this TI with two counterpropagating Fock states with equal number of photons, $\mathbf{k}$ and $-\mathbf{k}$, and study the influence of the TI on their occupation number [Figs. 1(a) and 1(b)]. The state of the field can be expressed as $|\psi_k, \phi_{-k}\rangle_j \equiv |\psi_k\rangle_j \otimes |\phi_{-k}\rangle_j$, where j denotes the medium properties before the j th TI (in the case of multiple TIs). The instantaneous transition approximation is accurate if the switching time τ_0 is much shorter than a cycle of the input light, i.e., $\tau_0 \ll 2\pi/\omega_j$ [4]. Detailed derivations

*Contact author: nnefedkin@gc.cuny.edu†Contact author: egaliffi@gc.cuny.edu

‡Present address: Chandra Department of Electrical and Computer Engineering, University of Texas at Austin, Austin, Texas 78712, USA.

§Contact author: aalu@gc.cuny.edu

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI.

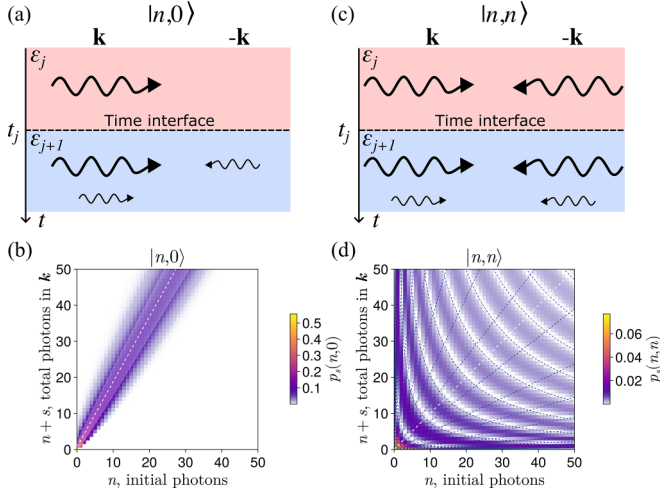


FIG. 1. (a) Sketch of two-mode field prepared in the Fock state $|n_k, n_{-k}\rangle_j = |n, 0\rangle_j$ before and after the TI; (b) photon number distribution (color axis) for asymmetric Fock state $|n, 0\rangle_j$ after a TI with $r_j = 0.2$ as a function of the initial number of photons in the modes and the total number of photons after the TI, where the white dashed line marks the average number of photons in k mode and the band denotes dispersion; (c) sketch for initial $|n, n\rangle_j$ state; and (d) photon number distribution (color axis) for a symmetric Fock state $|n, n\rangle_j$ and $r_j = 0.02$, where the white dashed line denotes the line where no new photons are created after the TI.

accounting for the effects of a noninstantaneous switching profile are available in the Supplemental Material [46].

At the TI, we enforce the continuity of the magnetic field operator and the displacement field operator: $\hat{\mathbf{B}}(\mathbf{r}, t = t_j^-) = \hat{\mathbf{B}}(\mathbf{r}, t = t_j^+)$ and $\hat{\mathbf{D}}(\mathbf{r}, t = t_j^-) = \hat{\mathbf{D}}(\mathbf{r}, t = t_j^+)$ [2], leading to canonical Bogoliubov transformations:

$$\begin{aligned}\hat{a}_j(\mathbf{k}) &= u_j \hat{a}_{j+1}(\mathbf{k}) - v_j \hat{a}_{j+1}^\dagger(-\mathbf{k}), \\ \hat{a}_j^\dagger(-\mathbf{k}) &= -v_j^* \hat{a}_{j+1}(\mathbf{k}) + u_j^* \hat{a}_{j+1}^\dagger(-\mathbf{k})\end{aligned}\quad (1)$$

for the bosonic creation and annihilation operators in the $\pm\mathbf{k}$ modes before (j) and after ($j+1$) the TI, which fulfill the symplectic relation $|u_j|^2 - |v_j|^2 = 1$ to preserve the commutation relations between operators. The coefficients read $u_j = (1 + r_j)/(2\sqrt{r_j})$, $v_j = (r_j - 1)/(2\sqrt{r_j})$ and the TI contrast r_j is defined as the ratio $\sqrt{\epsilon_j/\epsilon_{j+1}}$ between initial (j) and final ($j+1$) permittivity, which corresponds to the ratio ω_{j+1}/ω_j between final and initial frequencies. For one TI, the coefficients of the Bogoliubov transformation only depend on the permittivity before and after the TI, while, upon multiple TIs, they also depend on the interval between them [46].

The abrupt nature of the TI implies that the quantum state of the system does not change [47]. Thus, we can express the initial vacuum state $|0, 0\rangle_j$ in the Fock basis of the system after the TI [2]:

$$|0, 0\rangle_j = \frac{1}{|u_j|} \sum_{s=0}^{\infty} \left(\frac{v_j}{u_j}\right)^s |s, s\rangle_{j+1}. \quad (2)$$

Here, the initial vacuum state is written as a superposition of Fock states $|s, s\rangle_{j+1}$ after the TI, implying that the TI can create entangled photons from the vacuum. The coefficients

u_j and v_j act as squeezing parameters [48], implying the emergence of correlations between the photons in the two modes. In fact, the production of photons in pairs results from momentum conservation due to translational invariance, as discussed in Ref. [49], confirmed by the symplectic relation $|u|^2 - |v|^2 = 1$ between the Bogoliubov coefficients. In contrast, energy is not conserved in the process because of broken time-translation symmetry. This response is analogous to parametric amplification [50,51], which is described by the same Bogoliubov transformation as Eq. (1). The resulting photon number increase can be calculated from the probability distribution in the k and $-k$ modes.

To estimate the probability $p_s(0, 0)$ of creating s pairs of photons from the vacuum state, we project the state ${}_{j+1}\langle s, s|$ onto the initial vacuum state represented in the Fock basis of the $j+1$ medium [Eq. (2)] and then multiply it by its conjugate, yielding $p_s(0, 0) = {}_{j+1}\langle s, s|0, 0\rangle_j = |v_j|^{2s}/(1 + |v_j|^2)^{s+1}$. Using Eq. (2), we can also calculate the average number of photons produced by the TI in each mode from the initial vacuum state ${}_j\langle 0, 0|\hat{n}_{k/-k}|0, 0\rangle_j \equiv n_{\text{vac}} = |v_j|^2$ [52].

For a general initial Fock state $|n_k, n_{-k}\rangle_j$, applying the Bogoliubov transformations (1) yields

$$\begin{aligned}|n_k, n_{-k}\rangle_j &= \sum_{s=0}^{\infty} c_{s-\min(n_k, n_{-k})}(n_k, n_{-k}) |n_k + s \\ &\quad - \min(n_k, n_{-k}), n_{-k} + s - \min(n_k, n_{-k})\rangle_{j+1},\end{aligned}\quad (3)$$

where the probability amplitude $c_s(n_k, n_{-k})$ was derived in the Supplemental Material [46]. The photon number distribution for a general Fock initial state is thus given by

$$\begin{aligned}p_s(n_k, n_{-k}) &= {}_{j+1}\langle n_k + s, n_{-k} + s | n_k, n_{-k} \rangle_j^2 \\ &= |c_s(n_k, n_{-k})|^2,\end{aligned}\quad (4)$$

where $s \in [-\min(n_k, n_{-k}), +\infty)$ is the number of additional photon pairs created [46].

We start by considering the known case of only one mode initially populated $|n_k, n_{-k}\rangle_j = |n, 0\rangle_j$ [2], which yields the probability distribution $|c_s(n, 0)|^2$ shown in Fig. 1(b) as a function of the initial (x axis, “ n ”) and final (y axis, “ $n + s$ ”) photon numbers in the forward k mode. The white dashed line shows the average number of photons, $\langle \hat{n}_k \rangle = (n + 1)|v_j|^2 + n$. Note how the variance $\text{Var}(\hat{n}_k) = |v_j|^2(|v_j|^2 + 1)(n + 1)$ also scales linearly with the initial photon number n . Expectedly, increasing the TI contrast flattens the distribution against the y axis, implying the creation of a larger number of photon pairs.

Populating the initial $-k$ state dramatically changes the character of the photon number distribution, as shown in Fig. 1(c), where we consider a symmetric initial Fock state $|n_k, n_{-k}\rangle_j = |n, n\rangle_j$, yielding a probability amplitude [46]

$$c_s(n, n) = \frac{1}{|u_j|} \left(\frac{u_j^*}{u_j}\right)^n \left(-\frac{v_j}{u_j}\right)^s P_n^{(s, 0)}\left(1 - 2\frac{|v_j|^2}{|u_j|^2}\right), \quad (5)$$

where $P_n^{(a, b)}$ are Jacobi polynomials [53]. By initially populating both channels, we induce interference effects in the photon number distribution, leading to peaks and troughs that form hyperbolas in Fock space, as captured by Eq. (5) [Fig. 1(d)].

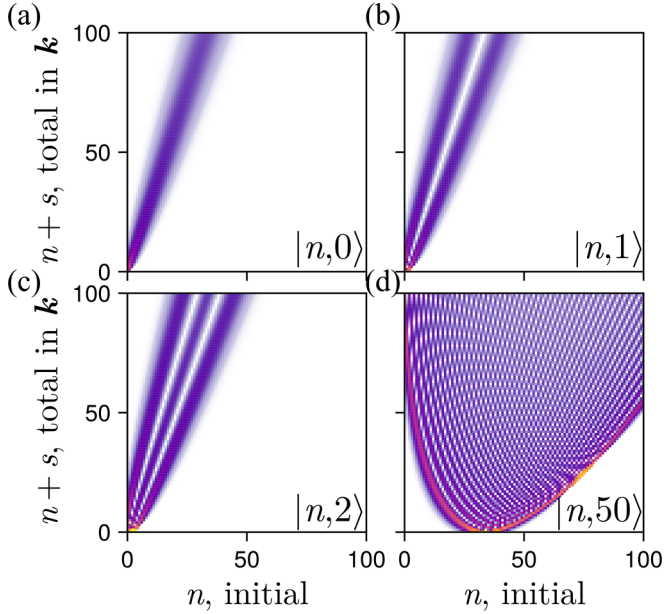


FIG. 2. Photon number distribution for initial states $|n, m\rangle_j$, where m varies from 0 [panel (a)] to 1 [panel (b)], 2 [panel (c)] and 50 [panel (d)], demonstrating the buildup of interference between initial photons and vacuum-generated pairs.

In the Supplemental Material [46], we derive the probability distribution for the symmetric case, which is characterized by a series of hyperbolas in the n - $n+s$ coordinate plane. The distance between the hyperbolic probability ridges in Fock space is fixed and can be determined from a cut of the distribution along the $s = 0$ line that intersects its vertices (white dashed line [46]).

The distribution can be simplified using Legendre polynomials and their asymptotes as [46]

$$p_0(n, n) \approx_{n \rightarrow \infty} \frac{1 + \sin((2n+1) \arccos(1 - 2\frac{|v_j|^2}{|u_j|^2}))}{2\pi n |v_j|}, \quad (6)$$

which reveals the frequency of the probability oscillations in Fock space $\omega_{\text{osc}} = 2 \arccos(1 - 2\frac{|v_j|^2}{|u_j|^2})$ or, if $n_{\text{vac}} \gg 1$, $\omega_{\text{osc}} \approx 4|v_j|^{-1} \equiv 4/\sqrt{n_{\text{vac}}}$. Concurrently, the probability decays as

$$g_k^{(2)}(0) = \frac{[n_{\text{vac}}^2(m^2 + 4mn + 3m + n^2 + 3n + 2) + n_{\text{vac}}(4mn + 2n^2 + 2n) + (n-1)n]}{(n_{\text{vac}}(m+n+1) + n)^2}. \quad (7)$$

The coherence function for the $-k$ mode can be obtained by swapping n and m in Eq. (7). The $g_k^{(2)}(0)$ function is shown in Fig. 3(a) as a function of the index contrast r_j , for initial states with (blue) one and (red) two (symmetrically) populated modes, and increasing values of n (arrows). In the limiting case $n \rightarrow 0$ and $m \rightarrow 0$, which refers to the initial vacuum state, we obtain $g_k^{(2)}(0) = 2$, independently of the TI contrast, as seen in Fig. 3(a), indicating thermal-like emission statistics. For the symmetric case $n = m$ and $n \rightarrow \infty$, the coherence $g_k^{(2)}(0) \rightarrow 3/2 - 1/[2(2n_{\text{vac}} + 1)^2] \rightarrow 3/2$ in the limit of large TI contrasts.

$\sim (n|v_j|)^{-1}$, i.e., linearly with the initial number of photons and with the square root of the number n_{vac} of photon pairs generated from the vacuum by the TI. We see from Eq. (6) that both the form of the probability distribution and its modulation frequency are inversely proportional to $\sqrt{n_{\text{vac}}}$, pointing to a crucial role of the interaction between the TI and the vacuum state in shaping the photon number distribution. Importantly, however, this interference only occurs if both forward and backward modes are initially populated.

To reveal the physical origin for the probability density modulation, we recall that the initial state consists of two counterpropagating equal-amplitude plane waves, which, classically, form a standing wave. After the TI, the pairs of photons are created everywhere from vacuum. Depending on the number of initial photons, the vacuum-generated photons can then interfere constructively or destructively with the scattered photons in the system. However, this interference occurs at the wave function level and results in a modulation of the probability distribution. In reality, we cannot observe interference patterns directly in the fields due to the phase uncertainty of Fock states.

We can control the probability distribution patterns by tuning the TI parameters and by applying multiple TIs. To qualitatively verify this possibility, we consider initial states in the form $|n, m\rangle_j$, where $m = 0, \dots, 50$. When $m = 0$, i.e., we only have photons in the $+k$ mode, a standing wave is not formed, and no interference pattern can be observed [Fig. 2(a)]. This is analogous to what has been investigated previously in Refs. [2,32,42]. As the number of photons m in the $-k$ channel increases from 0 [panel (a)] to 1 [panel (b)], the interference manifests as a minimum that cuts along the probability distribution [Fig. 2(b)], and each increase in m traces a new minimum in the probability distribution, e.g., 2 in panel (c) and 50 in panel (d), such that the number of interference minima corresponds to the minimum initial photon number in either mode, i.e., $\min(n, m)$.

The TI also drastically alters the photon statistics: We analytically calculate the second-order coherence function $g_k^{(2)}(0) = \langle \hat{a}_k^\dagger \hat{a}_k^\dagger \hat{a}_k \hat{a}_k \rangle / \langle \hat{a}_k^\dagger \hat{a}_k \rangle^2$ after the TI for a general initial Fock state, $|n, m\rangle_j$, which takes the form [46]

The asymmetry of the initial state impacts the emission statistics. For a fully asymmetric initial state, $m = 0$, when $n \rightarrow \infty$, the function $g_k^{(2)}(0) \rightarrow 1$ for any TI contrast, implying that in this limit, the photons detected after the TI are coherent, as seen by the blue lines in Fig. 3(a). Similarly, fixing the number of photons in one mode and letting the number in the other mode approach infinity results in $g_k^{(2)}(0) \rightarrow 1$.

In Fig. 3(b), we illustrate the more general case of a finite initial number of photons $n = 8$ in the $+k$ mode and vary the number m of initial photons in the $-k$ mode. The $g_k^{(2)}(0)$ function first increases to a maximum value of $3/2$ for the

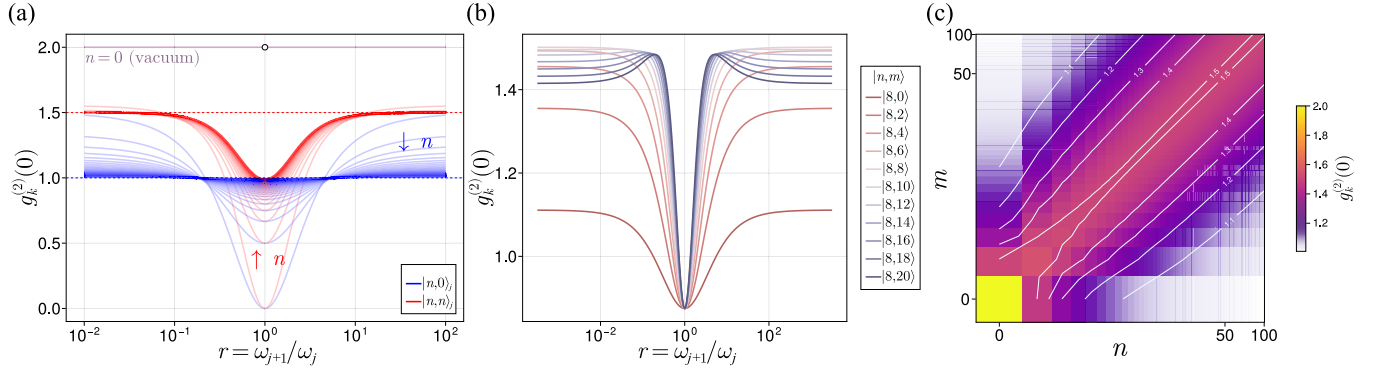


FIG. 3. Second-order coherence function for different initial states: (a) $g_k^{(2)}(0)$ as a function of the TI contrast r for symmetric $|n, n\rangle_j$ (red lines) and fully asymmetric $|n, 0\rangle_j$ (blue lines) initial states; (b) $g_k^{(2)}(0)$ function for the initial $|n, m\rangle_j$ state, where $n = 8$ is fixed and m changes; and (c) $g_k^{(2)}(0)$ in the parameter space of $n_k = n$ and $n_{-k} = m$ for the initial state $|n, m\rangle_j$ at fixed TI contrast $r = 0.02$.

symmetric initial state and then approaches 1 as the initial photon number in the second mode grows. The dependence of the $g_k^{(2)}(0)$ function on the TI contrast r_j is less apparent. From Fig. 3(b), we observe that when m exceeds n , the coherence function has a maximum $|g_{k,\max}^{(2)}(0)| = (3mn + m + 2n + 2)/(2mn + m + n + 1) \rightarrow_{m \gg 1} (3n + 1)/(2n + 1)$, when the average number of photon pairs created $n_{\text{vac}} = n_{\text{vac}}^{\max} = (n + 1)/(m - n - 1)$.

Figure 3(c) extends our analysis of the coherence function $g_k^{(2)}(0)$ across varying initial photon numbers at fixed large TI contrast. For large initial occupation numbers, $g_k^{(2)}(0)$ decreases monotonically from $3/2$ in the fully symmetric state to 1 in the fully asymmetric state, with linear isolines parallel to the diagonal in logarithmic scale. The coherence function has its maximum value of 2 for the vacuum state, where both occupation numbers are zero, highlighting the transition from photon bunching to coherent emission as asymmetry increases. In summary, the emission statistics and coherence of photons generated after the TI are tightly linked to the initial state and the TI contrast. Bunched emission dominates for large TI contrasts and finite initial photon numbers, while in the limit of large photon numbers, the $g_k^{(2)}(0)$ function tends to unity, revealing a coherent photon distribution.

The fact that the second-order coherence function $g_k^{(2)}(0) = 2$ for photons produced from the vacuum independently of the TI contrast r , mirroring the emission characteristics of a thermal source, prompts further investigation into the features of the resulting photon gas. Noting that the average number of “vacuum” photons in each mode is simply the squared modulus of the Bogoliubov coefficient v_j [46], we can rewrite the probability distribution as

$$p_s(0, 0) = \frac{1}{n_{\text{vac}} + 1} \left(\frac{n_{\text{vac}}}{n_{\text{vac}} + 1} \right)^s, \quad (8)$$

which resembles the Gibbs distribution [54], considering that the photons obey Bose-Einstein statistics. Furthermore, calculating the entanglement entropy [55] (see the Supplemental Material [46] for details), which monotonically increases with n_{vac} , reveals that the pairs of photons produced by the TI from the vacuum are entangled. The entropy growth with the TI contrast reflects the information lost upon tracing over the unobserved mode and reveals the strong correlation between photons in the k and $-k$ modes.

We also observe that the expectation values for the number of particles, coherence functions, and photon number distributions are symmetric with respect to an increase or decrease in permittivity, i.e., $r_j < 1$ and $r_j > 1$, depending only on the magnitude of the logarithm of the permittivity ratio $|\ln r_j|$.

In order to distinguish between increasing and decreasing the medium permittivity, we can employ the temperature of the photon gas created from the vacuum state after the TI, using the fact that the photons produced from a vacuum are thermally distributed Bose particles and expressing $k_B T$ through n_{vac} :

$$T = \frac{\hbar \zeta}{2\pi k_B}, \quad \zeta = \frac{2\pi r_j \omega_j}{\ln \left(\frac{n_{\text{vac}} + 1}{n_{\text{vac}}} \right)}. \quad (9)$$

Equation (9) connects the temperature of the photon gas to the number of photons created from the vacuum and, therefore, with the TI contrast. Note that the temperature (9) has the same form as the Unruh temperature [56] or the Hawking temperature [57,58]. The only difference is the particular form of the squeezing parameter $|v_j/u_j|^2$ [52].

Figure 4(a) shows the dependence of the temperature $k_B T$ on the TI contrast $r_j = \omega_{j+1}/\omega_j$. Although for $r_j = \sqrt{\epsilon_j/\epsilon_{j+1}} < 1$, the average number of photons $n_{\text{vac}} = (1 - r_j)^2/4r_j$ produced from the vacuum grows to infinity as $r_j \rightarrow 0$, the photon frequency $\omega_{j+1} = r_j \omega_j$ tends to zero, resulting in a finite average total energy in the system, $E_{\text{tot}} = \hbar r_j \omega_j [(1 - r_j)^2/4r_j + 1/2]$, which approaches $\hbar \omega_0/4$ in this limit. Thus, the temperature of the photon gas also has a finite limit $k_B T|_{r \rightarrow 0} = \hbar \omega_j/4$, equal to the average total energy.

In contrast, when the TI reduces the permittivity, $r_j = \sqrt{\epsilon_j/\epsilon_{j+1}} > 1$, the temperature of the photon gas grows to infinity as r increases, $k_B T|_{r \rightarrow \infty} = \hbar \omega_j r^2/4$. This behavior arises because n_{vac} is symmetric with respect to $r_j = 1$ and grows to infinity in both limits $r_j \rightarrow 0$ and the $r_j \rightarrow \infty$, while the frequency of the produced photons grows monotonically as $r_j \rightarrow \infty$, leading to an increase in the average total energy of the system and, consequently, the temperature of the photon gas. Interestingly, the temperature of the photon gas produced from the vacuum distinguishes the cases where the TI reduces or increases the permittivity, as opposed to other average values, such as the photon number and coherence functions, which are symmetric with respect to $\ln r_j = 0$.

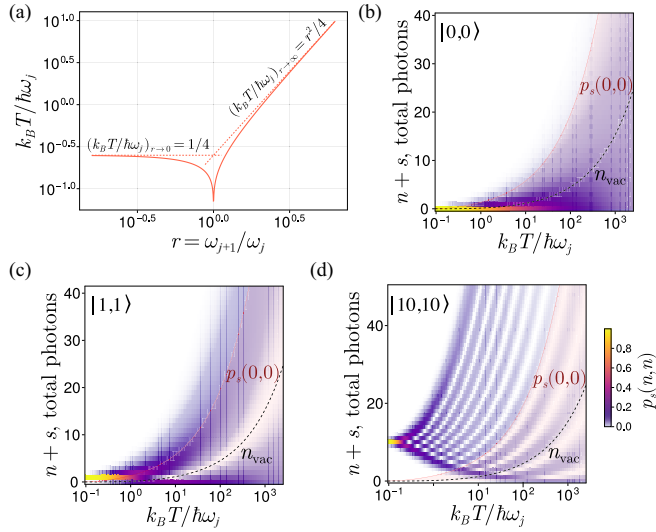


FIG. 4. (a) Dependence of the temperature of photon gas generated by the TI from the vacuum state for permittivity contrast r . Probability of finding $n + s$ number of photons in the k mode for different temperatures of vacuum photons when the initial state is (b) $|0, 0\rangle_j$, (c) $|1, 1\rangle_j$, and (d) $|10, 10\rangle_j$, where n is the initial photon number and s is the number of generated pairs; the red area represents the distribution for initial vacuum state, and the black dashed line shows the average number of vacuum photons.

Next, we examine how the photon number distribution depends on the temperature of the photon gas (note that at zero temperature, the TI contrast $r_j = 1$). Adding initial photons to the modes before the TI deviates the photon statistics from a Gibbs distribution, as shown in Fig. 4(b). In Fig. 4(c), the photon number distribution for an initial state with one photon in each mode $|1, 1\rangle_j$ reveals interference with a single minimum, deviating from the Gibbs distribution that arises from the vacuum state (red area). Finally, in Fig. 4(d), ten photons in each mode $|10, 10\rangle_j$ are considered, showing how the interference structure intensifies, showcasing ten minima and a distribution shape significantly different from the Gibbs distribution (red area). These results illustrate how increasing the initial photon number in each mode enhances the interference pattern, modulating the final photon gas away from a Gibbs distribution. This interference can be read off our analytical solution by comparing the probability distribution (8) under one or two populated channels:

$$p_s(n, 0) = p_s(0, 0) \frac{(n + s)!}{n! s! (n_{\text{vac}} + 1)^n}, \quad (10)$$

$$p_s(n, n) = p_s(0, 0) \left| P_n^{(s,0)} \left(\frac{1 - n_{\text{vac}}}{1 + n_{\text{vac}}} \right) \right|^2. \quad (11)$$

For the fully asymmetric initial state, $|n, 0\rangle_j$, the modulation term changes the position and width of the distribution maximum without introducing oscillations, supporting the results

in Fig. 2. On the other hand, for a symmetric initial state, the modulation term is a Jacobi polynomial of order n , whose number of oscillations (and zeros) defines the polynomial order. These oscillations reflect the interference of the wave function of the initial photons with those produced from the vacuum by the TI.

Finally, we discuss the robustness of our results to the presence of material dispersion and losses. For a general Lorentz-dispersive medium, one can show (via a Hopfield-Lagrangian treatment [59,60]) that when operating sufficiently far above (i.e., far-detuned from) a material resonance, dispersion primarily enters through a wave vector dependence of the effective mode normalization and hence of the Bogoliubov coefficients, which smoothly approaches the nondispersive limit at high frequencies. Regarding loss, the subcycle nature of the proposed effect implies that it need not rely on resonant probing; with suitable engineering, the relevant wave-medium interaction can be made much shorter than the intrinsic decay time, so that dissipation only weakly perturbs the predicted interference signatures.

In conclusion, in this Letter we have analytically studied the unique photon dynamics enabled by photonic TIs in the presence of two frequency-degenerate counterpropagating modes initially populated by Fock states. Our study reveals the intricate interference patterns produced by the interplay between the photon pairs generated from the vacuum state by the TI and the input photons in the two channels, producing an equal number of zeros in the photon number distributions as the minimum number of initial photons in either channel. Crucially, both input channels need to be populated to observe any of the phenomenology demonstrated, which qualitatively differs from any single-input scenarios.

By analyzing the coherence properties of the final photon distribution and identifying its departure from a thermal state, we reveal a behavior of the photon gas produced by the time interface analogous to the one found in Hawking and Unruh radiation. We can also define a corresponding temperature, characterizing its dependence on the TI parameters and its departure from the perfect thermal state induced by an initial state where both the forward and the backward modes are populated.

Our results shed light on the fundamental quantum processes triggered by TIs in two-mode systems, highlighting the crucial role of the vacuum in time-varying photonic systems. These findings pave the way for future research into photonic quantum state manipulation and frequency conversion using TIs, which may be readily tested in microwave QED setups [35].

This work was supported by the Air Force Office of Scientific Research NOFO and SBIR programs, and the Simons Collaboration on Extreme Wave Phenomena. E.G. acknowledges funding from the Simons Foundation through a Junior Fellowship of the Simons Society of Fellows (855344).

Data availability. The code and scripts are available at [61].

[1] F. R. Morgenthaler, Velocity modulation of electromagnetic waves, *IRE Trans. Microwave Theory Tech.* **6**, 167 (1958).

[2] J. T. Mendonça, A. Guerreiro, and A. M. Martins, Quantum theory of time refraction, *Phys. Rev. A* **62**, 033805 (2000).

- [3] V. Bacot, M. Labousse, A. Eddi, M. Fink, and E. Fort, Time reversal and holography with spacetime transformations, *Nat. Phys.* **12**, 972 (2016).
- [4] H. Moussa, G. Xu, S. Yin, E. Galiffi, Y. Ra'di, and A. Alù, Observation of temporal reflection and broadband frequency translation at photonic time interfaces, *Nat. Phys.* **19**, 863 (2023).
- [5] Z. Dong, H. Li, T. Wan, Q. Liang, Z. Yang, and B. Yan, Quantum time reflection and refraction of ultracold atoms, *Nat. Photon.* **18**, 68 (2024).
- [6] K. Komatsu, Z. Pápa, T. Jauk, F. Bernecker, L. Tóth, F. Lackner, W. E. Ernst, H. Ditlbacher, J. R. Krenn, M. Ossiander, *et al.*, Few-cycle surface plasmon polaritons, *Nano Lett.* **24**, 2637 (2024).
- [7] E. Lustig, O. Segal, S. Saha, E. Bordo, S. N. Chowdhury, Y. Sharabi, A. Fleischer, A. Boltasseva, O. Cohen, V. M. Shalae, *et al.*, Time-refraction optics with single cycle modulation, *Nanophotonics* **12**, 2221 (2023).
- [8] R. Tirole, S. Vezzoli, E. Galiffi, I. Robertson, D. Maurice, B. Tilmann, S. A. Maier, J. B. Pendry, and R. Sapienza, Double-slit time diffraction at optical frequencies, *Nat. Phys.* **19**, 999 (2023).
- [9] N. Engheta, Four-dimensional optics using time-varying metamaterials, *Science* **379**, 1190 (2023).
- [10] C. Firestein, A. Shlivinski, and Y. Hadad, Absorption and scattering by a temporally switched lossy layer: Going beyond the Rozanov bound, *Phys. Rev. Appl.* **17**, 014017 (2022).
- [11] A. Mekawy, H. Li, Y. Radi, and A. Alù, Parametric enhancement of radiation from electrically small antennas, *Phys. Rev. Appl.* **15**, 054063 (2021).
- [12] Z. Hayran and F. Monticone, Beyond the Rozanov bound on electromagnetic absorption via periodic temporal modulations, *Phys. Rev. Appl.* **21**, 044007 (2024).
- [13] N. A. Estep, D. L. Sounas, J. Soric, and A. Alu, Magnetic-free non-reciprocity and isolation based on parametrically modulated coupled-resonator loops, *Nat. Phys.* **10**, 923 (2014).
- [14] T. T. Koutserimpas and R. Fleury, Nonreciprocal gain in non-Hermitian time-Floquet systems, *Phys. Rev. Lett.* **120**, 087401 (2018).
- [15] S. Buddhiraju, W. Li, and S. Fan, Photonic refrigeration from time-modulated thermal emission, *Phys. Rev. Lett.* **124**, 077402 (2020).
- [16] Y. Hadad, J. C. Soric, and A. Alu, Breaking temporal symmetries for emission and absorption, *Proc. Natl. Acad. Sci. USA* **113**, 3471 (2016).
- [17] J. E. Vázquez-Lozano and I. Liberal, Incandescent temporal metamaterials, *Nat. Commun.* **14**, 4606 (2023).
- [18] V. Pacheco-Peña and N. Engheta, Temporal aiming, *Light Sci. Appl.* **9**, 129 (2020).
- [19] E. Galiffi, P. A. Huidobro, and J. B. Pendry, Broadband non-reciprocal amplification in luminal metamaterials, *Phys. Rev. Lett.* **123**, 206101 (2019).
- [20] E. Galiffi, A. C. Harwood, S. Vezzoli, R. Tirole, A. Alù, and R. Sapienza, Optical coherent perfect absorption and amplification in a time-varying medium, *Nat. Photon.* **20**, 163 (2026).
- [21] C. Rizza, G. Castaldi, and V. Galdi, Short-pulsed metamaterials, *Phys. Rev. Lett.* **128**, 257402 (2022).
- [22] R. Schützhold and W. G. Unruh, Hawking radiation in an electromagnetic waveguide? *Phys. Rev. Lett.* **95**, 031301 (2005).
- [23] S. A. Horsley and J. B. Pendry, Quantum electrodynamics of time-varying gratings, *Proc. Natl. Acad. Sci. USA* **120**, e2302652120 (2023).
- [24] E. Galiffi, R. Tirole, S. Yin, H. Li, S. Vezzoli, P. A. Huidobro, M. G. Silveirinha, R. Sapienza, A. Alù, and J. B. Pendry, Photonics of time-varying media, *Adv. Photonics* **4**, 014002 (2022).
- [25] F. Wilczek, Quantum time crystals, *Phys. Rev. Lett.* **109**, 160401 (2012).
- [26] H. Watanabe and M. Oshikawa, Absence of quantum time crystals, *Phys. Rev. Lett.* **114**, 251603 (2015).
- [27] K. Sacha, *Time Crystals* (Springer, Cham, Switzerland, 2020), Vol. 114.
- [28] M. S. Rudner and N. H. Lindner, Band structure engineering and non-equilibrium dynamics in Floquet topological insulators, *Nat. Rev. Phys.* **2**, 229 (2020).
- [29] X. Mi, M. Ippoliti, C. Quintana, A. Greene, Z. Chen, J. Gross, F. Arute, K. Arya, J. Atalaya, R. Babbush, *et al.*, Time-crystalline eigenstate order on a quantum processor, *Nature (London)* **601**, 531 (2022).
- [30] T. Liu, J.-Y. Ou, K. F. MacDonald, and N. I. Zheludev, Photonic metamaterial analogue of a continuous time crystal, *Nat. Phys.* **19**, 986 (2023).
- [31] G. T. Moore, Quantum theory of the electromagnetic field in a variable-length one-dimensional cavity, *J. Math. Phys.* **11**, 2679 (1970).
- [32] S. Sarkar, Photon statistics and moving mirrors, *Quantum Opt.* **4**, 345 (1992).
- [33] V. V. Dodonov, A. B. Klimov, and D. E. Nikonov, Quantum phenomena in nonstationary media, *Phys. Rev. A* **47**, 4422 (1993).
- [34] V. Dodonov, Current status of the dynamical Casimir effect, *Phys. Scr.* **82**, 038105 (2010).
- [35] C. M. Wilson, G. Johansson, A. Pourkabirian, M. Simoen, J. R. Johansson, T. Duty, F. Nori, and P. Delsing, Observation of the dynamical Casimir effect in a superconducting circuit, *Nature (London)* **479**, 376 (2011).
- [36] D. Dalvit, P. Milonni, D. Roberts, and F. Da Rosa, *Casimir Physics* (Springer, Berlin, Heidelberg, 2011), Vol. 834.
- [37] E. Yablonovitch, Accelerating reference frame for electromagnetic waves in a rapidly growing plasma: Unruh-Davies-Fulling-Dewitt radiation and the nonadiabatic Casimir effect, *Phys. Rev. Lett.* **62**, 1742 (1989).
- [38] C. Law, Effective Hamiltonian for the radiation in a cavity with a moving mirror and a time-varying dielectric medium, *Phys. Rev. A* **49**, 433 (1994).
- [39] A. V. Dodonov and V. V. Dodonov, Dynamical Casimir effect via modulated Kerr or higher-order nonlinearities, *Phys. Rev. A* **105**, 013709 (2022).
- [40] S. D. Liberato, C. Ciuti, and I. Carusotto, Quantum vacuum radiation spectra from a semiconductor microcavity with a time-modulated vacuum Rabi frequency, *Phys. Rev. Lett.* **98**, 103602 (2007).
- [41] A. Ganfornina-Andrades, J. E. Vázquez-Lozano, and I. Liberal, Quantum vacuum amplification in time-varying media with arbitrary temporal profiles, *Phys. Rev. Res.* **6**, 043320 (2024).
- [42] M. S. Mirmoosa, T. Setälä, and A. Norrman, Quantum state engineering and photon statistics at electromagnetic time interfaces, *Phys. Rev. Res.* **7**, 013120 (2025).

- [43] E. Galiffi, G. Xu, S. Yin, H. Moussa, Y. Ra'di, and A. Alù, Broadband coherent wave control through photonic collisions at time interfaces, *Nat. Phys.* **19**, 1703 (2023).
- [44] D. G. Baranov, A. Krasnok, T. Shegai, A. Alù, and Y. Chong, Coherent perfect absorbers: Linear control of light with light, *Nat. Rev. Mater.* **2**, 17064 (2017).
- [45] J. T. Mendonça and A. Guerreiro, Time refraction and the quantum properties of vacuum, *Phys. Rev. A* **72**, 063805 (2005).
- [46] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/z61w-gdvl> for the derivations and the additional information.
- [47] L. D. Landau and E. M. Lifshitz, *Quantum Mechanics: Non-relativistic Theory* (Elsevier, Oxford, UK, 2013), Vol. 3.
- [48] D. Walls and G. J. Milburn, *Quantum Optics* (Springer, Berlin, Heidelberg, 2008), pp. 307–346.
- [49] E. Galiffi, S. Yin, and A. Alù, Tapered photonic switching, *Nanophotonics* **11**, 3575 (2022).
- [50] N. D. Birrell and P. C. W. Davies, *Quantum Fields in Curved Space* (Cambridge University Press, Cambridge, UK, 1984).
- [51] U. Leonhardt, *Essential Quantum Optics: From Quantum Measurements to Black Holes* (Cambridge University Press, Cambridge, UK 2010).
- [52] P. D. Nation, J. R. Johansson, M. P. Blencowe, and F. Nori, *Colloquium: Stimulating uncertainty: Amplifying the quantum vacuum with superconducting circuits*, *Rev. Mod. Phys.* **84**, 1 (2012).
- [53] T. S. Chihara, *An Introduction to Orthogonal Polynomials* (Courier Corporation, Mineola, NY, 2011).
- [54] L. D. Landau and E. M. Lifshitz, *Statistical Physics* (Elsevier, Oxford, UK, 2013), Vol. 5.
- [55] M. Srednicki, Entropy and area, *Phys. Rev. Lett.* **71**, 666 (1993).
- [56] W. G. Unruh, Notes on black-hole evaporation, *Phys. Rev. D* **14**, 870 (1976).
- [57] S. W. Hawking, Black hole explosions? *Nature (London)* **248**, 30 (1974).
- [58] S. W. Hawking, Particle creation by black holes, *Commun. Math. Phys.* **43**, 199 (1975).
- [59] B. Huttner and S. M. Barnett, Quantization of the electromagnetic field in dielectrics, *Phys. Rev. A* **46**, 4306 (1992).
- [60] F. Belgiorno, S. L. Cacciatori, and F. Dalla Piazza, The Hopfield model revisited: Covariance and quantization, *Phys. Scr.* **91**, 015001 (2016).
- [61] https://github.com/venichka/ti_fock