Microgap thermophotovoltaic systems with low emission temperature and high electric output

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Abstract
We theoretically show that a thermophotovoltaic (TPV) system enhanced by a wire metamaterial opens the door to a prospective microgap TPVs which will combine high electric output with relatively low temperatures of the emitter. The suggested system comprises an array of parallel metal nanowires grown on top of a photovoltaic semiconductor and standing free in the vacuum gap between the host dielectric layer and the emitter, so that their ends are sufficiently close to the emitting surface. Due to the resonant near-field coupling between this wire medium and the emitter and due to the optimized layered structure of the whole system, the strongly super-Planckian radiative heat flux of resonant nature is engineered. In the suggested system, heavily doped silicon and indium antimonide are considered as the materials for the emitter and the photovoltaic cell, respectively. Also, the parallel nanowires are made of tungsten. Employing the minority-carrier transport model, it is shown that a power output equal to $26 \text{ kW m}^{-2}$ can be achieved when the temperature of the doped-silicon emitter is only $500 \degree C$.

Keywords: wire metamaterials, radiative heat transfer, microgap thermophotovoltaics, hyperbolic metamaterials

1. Introduction
In thermophotovoltaic (TPV) systems, a hot emitter of thermal radiation whose temperature is determined by a heat source (e.g. flame) transmits infrared photons to the photovoltaic (PV) cell as it is shown in figure 1 [1]. A free-space gap separates the emitter from the PV cell, preventing the conductive thermal flux [1, 2]. Most popular TPV systems are so-called far-field ones, in which the gap thickness $d$ is larger than the free-space wavelengths $\lambda_\text{L}$ (the largest one in the model spectrum of the emitted radiation). In these systems the process of radiative heat transfer (RHT) splits onto two independent stages: thermal radiation governed by the Stefan–Boltzmann and Wien laws, and its absorption by the PV material. If the emitter’s temperature $T_\text{E}$ is low, the semiconductor of the PV cell must possess a very narrow bandgap $E_\text{g} = h\omega_\text{g}$ ($\omega_\text{g}$ is called the bandgap frequency and $h$ is the Planck constant). Really, if the frequency $\omega_\text{M}$ at which the radiation spectrum is maximal becomes smaller than $\omega_\text{g}$, the most part of radiation cannot be converted to the photocurrent. Also, the dissipation in the PV material is not admissible. Practically, one needs $\omega_\text{M} \approx \omega_\text{g}$ [1]. For temperatures $T_\text{E} < 400 \text{ K}$, $\omega_\text{M} = 2\pi T_\text{E} \cdot (5.87 \times 10^{10}) < 1.47 \times 10^{14} \text{ rad s}^{-1}$, that requires $\omega_\text{g} < 1.5 \times 10^{14}$. PV materials with so narrow bandgap have been recently synthesized, but they cannot operate at room temperatures [1] and require the energy-consuming cooling. As a result, far-field TPV systems operating with such $T_\text{E}$ cannot serve electric generators. A far-field TPV system whose PV cell operates at room temperatures requires...
$T_E > 600$ K. However, even the temperature $T_E \approx 600$ K is not practically sufficient, because the thermal emission governed by the Stefan–Boltzmann law is low at this temperature, and the electric output is poor. Existing far-field TPV systems operate as electric generators at $T_E > 800$ K [1]. To extract the electric energy from bodies with lower temperatures one has to apply thermoelectric converters whose heat-to-electricity conversion efficiency, in practice, cannot exceed 3% [3]. Notice that the PV conversion efficiency in realistic TPV generators operating at $T_E \approx 2000$ K attains 20% [4].

Planckian spectrum of thermal radiation is very broadband, and an optical filter with $\omega_{\text{min}} = \omega_q$ shown in figure 1 is obviously needed in the TPV system. The filter suppresses the reflection in the operation band of the PV cell (operating as an antireflector) and fully reflects the emitted harmonics beyond it. In fact, this filter is a bandpass filter and it has a relatively narrow pass-band. This is so because the so-called ultimate efficiency of the PV cell decreases versus the band of incident radiation that is called the Shokley–Queisser limit. Low ultimate efficiency implies high dissipation, strong heating, and finally, negligible electric output [1]. Therefore, the most part of the emitted power should be reflected that is known as Carnot loss (although the reflected radiation may maintain the operation temperature of the emitter). For maximal electric output the filtering is not optimal, and one needs to engineer a narrow-band thermal radiator. This is possible with a resonant emitter—nano-textured, nano-patterned or performed as an array of IR nanoantennas [1, 2, 5, 6]. Its emissivity attains unity at a resonant frequency and it is negligibly small beyond a narrow frequency band which is equal to the operation band of an optimized infrared PV cell. The conversion efficiency of the PV in such systems may theoretically achieve 80%–90% (see e.g. [1]). However, the electric output is determined not only by the efficiency of the conversion, but also by the emitted power. Far-field emission of a resonant emitter keeps fundamentally restricted by the black-body radiation. Then to maximize the emitted power in the operation band one has to engineer the resonance frequency $\omega_{\text{res}}$ nearly equal to $\omega_M$ of the black body having the same temperature. Therefore, far-field TPV systems with resonant emitters are justified only with $T_E > 1500$–2000 K [1, 2, 5].

Then, the question arises: is it possible to achieve a high electric output if the emitter has relatively low temperature. This can be an intriguing goal in the field of waste heat harvesting. As a reference value we admit the temperature $T_E = 500$ °C = 773 K, which is typical for the shell of a car exhausting pipe in the peak regime [7]. Specially designed thermoelectric generators may convert this waste heat into electricity [7]. At this temperature, the output power of the best commercial thermoelectric generator is equal to 2.2 W harvested from the area 0.4 m$^2$. This is the value we want to exceed with an advanced TPV system. In other words, we want to design a TPV system with $T_E = 500$ °C with electric power output noticeably higher than 5.5 W m$^{-2}$. Since at this temperature $\omega_M = 2.75 \times 10^{14}$ rad s$^{-1}$, the most suitable PV material is indium antimonide (InSb) with $\omega_q = 2.58 \times 10^{14}$ rad s$^{-1}$. It has sufficiently high PV spectral response $R$ at room temperatures, which at frequencies $\omega > \omega_q$ can be modeled as $R = e_0/\hbar \omega$, where $e_0$ is the electron charge [1]. The operation band of such a PV cell can be nearly $(2.6-4) \times 10^{14}$ rad s$^{-1}$ that allows the ultimate efficiency higher than 70%. However, a far-field TPV system with such a PV cell will not allow the power output higher than 5.5 W m$^{-2}$ due to two factors: $T_E = 500$ °C corresponds to low power radiated by the black body in the spectral range $(2.6-4) \times 10^{14}$. Black-body power flux integrated over this range is nearly equal to 20 W m$^{-2}$, and this is the maximal power for a resonant emitter. The realistic overall efficiency of the PV cell is below 20%, that means the output power below 4 W m$^{-2}$.

The way on which we theoretically attain our goal is the involvement of the near-field RHT. The near-field coupling, also called phonon tunneling, is illustrated in the inset of figure 1. If the thickness of the free-space gap (medium 2) between hot medium 1 and cold medium 3 (both these media are semitransparent) is subwavelength, the total internal reflection is frustrated. A part of the incident wave energy is transferred from the hot medium to the cold one by a couple of mutually coherent evanescent waves. In the black-body limit RHT is restricted by the waves incident at the interface between media 1 and 2 with the angles $\theta < \theta_{\text{IR}}$. Beyond the black-body limit (if $d \ll \lambda$) this spatial spectrum broadens up to $\theta < \pi/2$ and even more, because waves which are evanescent in medium 1 are also involved. Such RHT becomes super-Planckian [8, 9]. The super-Planckian RHT at some frequencies means that the spectrum of RHT between two parallel surfaces having temperatures $T_E$ and $T_C < T_E$ exceeds the spectrum of RHT between two parallel surfaces of two black bodies with same temperatures. In the last case RHT does not depend on the distance $d$ and reproduces the Planckian radiation spectrum. The efficient involvement of evanescent waves into RHT across a subwavelength gap is a resonant process. A pair of mutually coupled surface-phonon polaritons excited in the IR range at the interfaces of media 1 and 3 with medium 2 can dramatically enhance the RHT making it strongly super-Planckian [10–13]. This effect resulted in a scientific direction called near-field TPV systems.
(see [14–24]). Near-field TPV systems cannot be physically split onto an emitter and an absorber of emission, and the RHT process is not anymore the two-stage one. Obviously, the conventional Stefan–Boltzmann and Wien laws restricting the TPV systems and resulting in the requirement of high temperatures do not become applicable. In a near-field TPV system one may try to obtain a high electric output with rather low $T_E$ [24, 25]. Of course, the efficient PV conversion still requires the strong frequency selectivity of the RHT. A conventional filter is not helpful, because we need photon tunneling from the emitter to the PV cell which will be blocked by the filter. However, both emitter and PV cell of a near-field TPV system can be structured so that to engineer the resonant RHT in the predefined frequency range. A theoretically successful attempt in this direction is based on the application of graphene sheets [20]. However, near-field TPV electric generators requiring the gap as tiny as $d = 10–20$ nm are challenging because the so tiny gap can not be maintained between two parallel surfaces at high temperatures. If the gap is much larger, the RHT can not be strongly super-Planckian [1, 25, 26]. Therefore, in our group we have developed, starting from our initial work [27] the theory of strongly super-Planckian RHT in TPV systems with a micron or slightly submicron gap. Such systems called microgap TPV systems imply the vacuum gap $d \geq 0.5$ $\mu$m. Starting from the invention [28], this technology (based on tubular quartz spacers and adjusting springs) has developed to practical TPV devices with the working area of square decimeter [29, 30]. In this paper, we theoretically demonstrate that a micron gap TPV system enhanced by a metamaterial allows the frequency selective strongly super-Planckian RHT from the emitter to the PV material for $T_E = 500 ^\circ$C. The main novelty of the present paper is namely this relatively low emission temperature. We start from our previous design solution and explain how to modify this structure in order to achieve the desired goal. In next section, we present our numerical model and discuss the obtained results.

2. New design of a micron gap TPV system enhanced by nanowires

The idea of our new design becomes clear if we start from the previously suggested structure which was studied in [34] and considered there as the most practical one. Initially, in order to achieve the strongly super-Planckian RHT, we suggested to grow nanowires at the sides of the micron gap such that they formed an interdigital structure [27]. In our further works [31–33], we have developed this approach and shown that the narrow-band strongly super-Planckian RHT is achievable in such systems with feasible design parameters. Finally, in [34], we have studied the power output and optimized the whole structure imposing some practical restrictions (no energy-consuming cooling and no challenging interdigital geometry of the wire medium). In [34], we obtained an unprecedent $3.3 \times 10^4$ $\text{W m}^{-2}$ power output on the condition that the temperature of the PV cell is maintained near the room temperature by using a water cooling system. This structure (the cooling frame is not shown) is schematically depicted in figure 2(a). The photon tunneling occurred in our model across the tiny gap ($d = 10$ nm) at the ends of hot tungsten nanowires connected to the hot silicon carbide plate. The wire medium supports propagating the huge spatial spectrum of thermal radiation. Spatial harmonics with transversal (in-plane) wave numbers $q > k_0 \equiv \omega/c$ (which in the vacuum gap correspond to evanescent waves) are propagating in the wire medium [35–37]. This feature of the wire medium allows a huge resonant RHT across the structure. It occurs in a predefined narrow band of frequencies. To design the structure we have used the method [31], where the radiative heat is expanded onto both frequency $\omega$ and spatial $q$ spectrum, and for each spatial harmonics the impedance matrices of all effective layers of the structure are simulated. The method takes into account the generation of radiative heat in all effective layers and fits the alternative method based on the Green function of the whole structure [31, 33]. Though the method of the Green function [38] is more popular, our method seems to be more suitable for the design of frequency-selective TPV systems. Our method delivers the surface impedances at the internal interfaces of the structure as intermediate results. This is very important for the optimization since our design strategy is to match the impedances at both sides of every interface in the predefined frequency

![Figure 2. (a) Previous structure with hot nanowires. (b) New structure with cold nanowires.](image-url)
range. If at one frequency the perfect matching is achieved, then far from this frequency (i.e. beyond the operation band of the PV cell) the spectrum of RHT turns out to be suppressed due to the strong impedance mismatch [31].

The structure shown in figure 2(a) satisfies to all these requirements and it was considered in [34] as most suitable for experimental realization. However, the target of [34] was different from the present one. We aimed to theoretically beat the record for the power output previously claimed in the literature for the same frequency range. If at one frequency the perfect matching is achieved, the power output will be far from this frequency range. If at one frequency the perfect matching is achieved, the power output will be.

Figure 2 schematically shows this structure. Instead of SiC enhanced by W nanowires as in [34], we suggest a flat emitter of heavily doped Si (doping level \(2 \times 10^{20}\) cm\(^{-3}\)). This geometry allows the coupled surface plasmon polariton (SPP) to be excited at the interfaces silicon-vacuum and wire medium-vacuum. Calculations based on the effective-medium model for the wire medium of W nanowires [35] have shown that the band where the SPP is excited is adjustable via the design parameters of the wire medium. The involvement of SPP combines the advantages of our reference design figure 2(a) and those of the near-field TPV generator suggested in [20]. Really, in our reference design the role of the wire medium was only to effectively approach the emitter to the PV cell. Therefore, the RHT though super-Planckian was not very strong, and the gain compared to the black-body RHT in the operation band was nearly equal 6. Now, the wire medium layer is strongly coupled to the emitter via the SPP, and the expected gain is much higher.

The optical parameters of heavily-doped silicon at \(T_E = 500\) °C were taken from [39]. The array of nanowires is assumed to be grown in the dielectric medium with thickness \(W_g\) and permittivity \(\varepsilon_h\) to be found in the numerical optimization. Parameters of W were taken from [40]. An array of tungsten nanowires partially hosted by the dielectric and partially free-standing is replaced in our calculations by two homogeneous layers of effective uniaxial media which are presented through an effective permittivity tensor as following

\[
\varepsilon = \begin{pmatrix}
\varepsilon_\parallel & 0 & 0 \\
0 & \varepsilon_\parallel & 0 \\
0 & 0 & \varepsilon
\end{pmatrix}
\]  

(1)

Here, \(\varepsilon_\parallel\) and \(\varepsilon\) are called the transverse and the axial components, respectively, and they are given by [33]

\[
\varepsilon_\parallel = \varepsilon_h (1 + f_c) \varepsilon_m + (1 - f_c) \varepsilon_h,
\]

\[
(1 - f_c) \varepsilon_m + (1 + f_c) \varepsilon_h
\]

(2)

In the above equation, \(f_c\) is the volume fraction of nanowires, \(\varepsilon_m\) is the tungsten permittivity and \(\varepsilon_h\) represents the relative dielectric constant of the corresponding host medium in each layer. One layer corresponds to the free-standing parts of nanowires (\(\varepsilon_h = 1\)), another layer—to the hosted parts. Both these effective media have hyperbolic dispersion (this kind of dispersion surface is useful for many applications, see e.g. [36, 41, 42]). Here, we use the effective-medium model (equation (2)) whose application for the analysis of RHT was validated in our previous work [33]. In principle, nanowires can be made of another materials such as gold. However, for gold nanowires the simple effective-medium model turns out to be in accurate [33]. So, the choice of tungsten allowed us to avoid full-wave simulations evaluating the impedance matrices of wire-medium layers. Generally speaking, the applicability of the effective-medium model for wire media is determined by two main conditions: (1) sufficiently low spatial frequencies \(q a \ll \pi\) (a is the array period) and (2) long wavelengths \(k_0 a \ll 2\pi\). If the radius of the wires is much smaller than the wavelength \(k_0 r_0 \ll 1\) (\(r_0\) is the radius of each nanowire) and the fraction of wires is rather small, the effective-medium model is quite accurate.

The tiny gap \(W_g\) prevents the harmful contact of nanowires and the emitter. The realistic fabrication tolerances of wire media [36] impose the restriction \(W_g \geq 10\) nm. The gap between the emitter and the dielectric interface \(d = W_g + W_s\) is taken minimal possible in the microgap technology [29, 30] i.e. equals 500 nm. As it was already mentioned, to ensure the parallelism of two stacked surfaces with \(d = 10\) nm gap (and allowed deviation plus minus 10 nm) is presently possible only on the area of 1 μm\(^2\). To ensure the parallelism of two stacked surfaces with \(d = 500\) nm gap is possible on the area of up to 1 square decimeter whereas the deviation also is plus minus 10 nm. Therefore, the gap \(W_g = 10\) nm between the ends of free-standing nanowires and the emitter means that the PV cell of the area 1 square decimeter will be thermally disconnected from the emitter. Our wires will not touch the emitter.

The volume fraction of nanowires (\(f_c = \pi r_0^2/a^2\)) was varied in our calculations from 0 to 0.3 in order to adjust the band of SPP combining in with the band of the best matching. As it was explained above, we chose InSb as a material for the PV cell. Complex permittivity of the InSb was taken from [17, 20].

3. Numerical study

In this section we present the results on RHT across the structure as well as the results for the electric current passing the load in our PV cell.
In the above equation we have [18, 23]
\[ j_{h,p}(\omega) = e_0 \left[ D_e \frac{d n_{h,p}(\omega, z = z_0)}{dz} + D_h \frac{d n_{h,p}(\omega, z = z_0)}{dz} \right] \]

where \( e_0 \) is the electron charge, \( D_{e,h} \) is the diffusion constant, \( z_{a,p} \) are coordinates of the edges of the depletion region, and \( \text{Im}(\beta) \) is the imaginary part of the vertical component of the wave vector in InSb: \( \beta = \sqrt{k_{\text{InSb}}^2 - q^2} \), where \( q \) is the transverse component of the wave vector of spatial harmonic (see above), and \( k_{\text{InSb}} \) is the complex wave number of InSb at frequency \( \omega \). In equation (4), \( Q(\omega, q) \) is the frequency-and-space spectrum of RHT i.e. power absorbed by the unit area of the PV cell per unit interval of frequencies and unit interval of \( q \). Finally, \( n_{e,h} \) denote the minority carrier concentrations.

These concentrations are solutions of the equation [23]
\[ \frac{D_{e,h}}{\text{d}z^2} \frac{d n_{e,h}}{\text{d}z} = \frac{n_{e,h}}{\tau_{e,h}} + \frac{1}{\pi} \int_0^{\infty} \text{Im}(\beta) Q(\omega, q) e^{-2\text{Im}(\beta)qz} \text{d}q = 0. \]

Here, \( \tau_{e,h} \) is the relaxation time. To calculate \( Q(\omega, q) \), we used the method introduced in [31]. The calculation is shown in appendix A.1.

The solutions of equation (5) were obtained imposing four boundary conditions: two for \( n_e \) (at \( z = 0 \) and at \( z = W_e \)) and two for \( n_h \) (at \( z = z_0 \) and at \( z = W_e + W_h \)). The first-type condition is \( D_{e,h}(d n_{e,h}/dz) = S_{e,h} n_{e,h} \) (at \( z = 0 \) for electrons and at \( z = W_e + W_h \) for holes), where \( S_{e,h} \) is the surface recombination rate. The second-type condition is \( n_{e,h} = 0 \) at \( z = W_e \) for electrons and at \( z = W_e + W_h \) for holes. Solving equation (5) we find both components of the photocurrent in equation (3). In accordance to the equivalent circuit in figure 1, we find the total current collected from unit area \( J \) (current in the load divided by the illuminated area of the PV cell [1]) as the difference \( J = J_{\text{diode}} - J_{\text{ph}} \), in other words
\[ J = \left[ e_0 n_i^2 \left( \frac{D_e}{N_p L_e} + \frac{D_h}{N_n L_h} \right) \exp \left( \frac{e_0 V}{K_B T} \right) - 1 \right] - J_{\text{ph}}, \]

in which \( K_B \) is the Boltzmann constant, \( n_i \) denotes the intrinsic carrier concentration, \( N_{p,n} \) represents the carrier concentration, \( L_{e,h} = \sqrt{D_{e,h} \tau_{e,h}} \) is the diffusion length and \( V \) is the load voltage which varies from zero (short-circuit case) to its maximum value as the load is considered to be infinity (open-circuit case). Table 1 contains the parameters that we use to calculate \( J \). These parameters were taken from [17, 22, 43] (see also appendix A.2).

Table 1. Input parameters for calculating the photocurrent density and the current density referred to the load.

| Parameter          | p-region    | n-region    |
|--------------------|-------------|-------------|
| \( D_e \) cm² s⁻¹  | 186         | 5.21        |
| \( \tau_e \) ns    | 1.45        | 1.81        |
| \( S_e \) cm⁻¹ ns⁻¹| 10^4        | 0           |
| \( W_e \) nm       | 100         | 10          |
| \( W_h \) nm       | 500         | 730         |
| \( N_p \) cm⁻³     | 10¹⁹        | 10¹⁹        |
| \( n_i \) cm⁻³     | 2 × 10¹⁶    |             |

For the optimal load, the electric power density \( P_{\text{elec}} = J \cdot V \) gets the maximum value. Then we can obtain the photovoltaic conversion efficiency as the ratio between the maximal (over all possible loads) electric power output per unit area of the cell to the integral value of RHT \( P_{\text{radiative}} \)—power flux at the interface \( z = 0 \) between the dielectric layer and InSb. In other words, we have
\[ \eta = \frac{\text{Max}[P_{\text{elec}}]}{\text{Max}[J \cdot V]} = \frac{1}{2\pi} \int_0^{\infty} \text{d}\omega \int_0^{\infty} Q(\omega, q) \text{d}q. \]

Table 2 illustrates the optimized values for the structure parameters. The spectrum of RHT (power absorbed per unit area of the PV cell per unit interval of angular frequencies) is shown in figure 3(a). This spectrum is shown in comparison with that of the black-body RHT. We see that it is much more narrow-band than the black-body RHT and exceeds it by an order of magnitude in the band between \( \omega_{\text{min}} = 2.6 \times 10^{14} \text{ rad s}^{-1} \) and \( \omega_{\text{max}} = 5 \times 10^{14} \) (the band-gap frequency of the PV material is marked by red line). The spectral maximum at which the RHT spectrum exceeds the black-body one by two orders of magnitude holds at \( \omega_M \approx 3.3 \times 10^{14} \text{ rad s}^{-1} \), and the crucial part of the resonant spectrum is above \( \omega_e \). As figure 3(a) indicates, increasing \( W_e \) from 10 to 20 and 30 nm causes the decreasing in the maximum value at \( \omega_M \).

For comparison we have calculated RHT in our structure removing the nanowires. The result is visually very close to the black-body RHT. Therefore, the difference between two curves in figure 3(a) is practically equal to the gain granted by nanowires. Figure 3(b) shows the corresponding volt–ampere
the PV conversion efficiency 24% for a near-field TPV system also operating at 500 °C. In that structure, the gap between the emitter (hexagonal boron nitride) and the semiconductor (indium antimonide) was equal 16 nm (also an atomic layer of graphene with the chemical potential of 0.5 eV was located on the surface of InSb). Clearly, this structure essentially refers to near-field TPV systems, whereas we claim the results of the same order for a microgap TPV system.

Moreover, our results can be improved if we soften the restriction for the thickness of the host dielectric. We have taken above $W_h = 730$ nm that is approximately 60% of the total length $W_h + W_b$ of a nanowire because such free-standing wire media were reported in the literature (see e.g. [36]). However, there are no physical reasons prohibiting smaller values for the hosted length $W_h$ than 60% of the total length for a so robust metal as tungsten. As table 3 indicates, decreasing the thickness $W_b$ from 730 to 300 nm results in the increasing of the conversion efficiency from 13% to 16.30%. Nanowires with a free-standing portion longer than the hosted portion are not yet reported in the literature, however, we believe that they are feasible. In the case of $W_b = 300$ nm, the maximum electric power is about 40.64 kW m⁻². By fixing $W_h = 300$ nm and trying to change other parameters of the structure, we can even match better the emitter to the semiconductor near its band gap. The best case is when the parameters in table 2 remain the same except the relative dielectric constant of the host medium. If this constant changes from 2.25 to unity (a transparent dielectric at infrared), the conversion efficiency and the maximum output power increase to 18.7% and 55.2 kW m⁻². These values are closer to those claimed in [20].

Finally, in comparison with the practical macroscopic system, where the heat-to-electricity conversion resulted in the power output 5.5 W m⁻² for the temperature 500 °C [7], our 55.2 kW m⁻² mean the gain 10⁴. This amazing gain, to our opinion, will justify the fabrication costs of a nanostructured micro-gap TPV system.

4. Conclusions

In the present work we have introduced a new type of efficient and feasible low-temperature TPV system with very high electric output. The system is based on the existing microgap technology for TPV systems, and, therefore, potentially allows one to collect the photocurrent from macroscopic areas covering the emitter with temperatures 400 °C–600 °C. The use of the realistic tungsten wire medium grown on top of an IR PV cell together with the heavily-doped silicon emitter results in our calculations in the excitation of a surface-phonon-polariton at the effective internal interfaces of the system. Then the wire medium not only effectively collects the electrons from the surface of the PV cell but also couples them in a resonant way. The frequency selectivity is achieved not only due to the effectively multi-layer structure of the TPV system and high contrast between layers as in our precedent works. It is enhanced by the resonant

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Figure 3. (a) Spectrum of radiative heat transfer at $z = 0$. The blue and black curves correspond to our optimized structure and to the black-body RHT, respectively. The dashed and dotted curves show the effect of the thickness of the tiny gap $W_g$. (b) The current density over the load voltage ($J$-$V$ curve) (here, $W_g = 10$ nm). Notice that the sign of the current density ($J$) is negative due to its direction (see figure 1).

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characteristics of our PV cell ($J$-$V$ curve). The black and blue curves are diode current and the total current, respectively. At the optimal load, the electric power is equal $P_{elec} = 26$ kW m⁻² corresponding to the total (PV) conversion efficiency $\eta = 13\%$ (here, it is worth noting that we consider an isothermal assumption. In other words, we assume that the temperature of the PV cell is maintained near $T_C = 27$ °C. Practically, if thermal losses are taken into account, the temperature of the cell is different and it may increase which results in an increase of the diode current. Therefore, the power output reduces. We hope that this decrease compared to our predictions will not be drastic). Work [20] has reported...
excitation of the SPP, that results in further frequency squeeze of the spectrum of RHT and enhancement of its maximum. A huge gain—four orders of magnitude—for the electric power output compared to commercial analogs operating at same temperatures is obtained. The next stage of our studies will be the further decrease of $T_{ph}$ in a completely new structure which will allow the TPV electric generators to replace the thermoelectric ones in the electric power supplies, computer processors, etc.

Appendix

A.1. Frequency-and-spatial spectrum of radiative heat transfer

In equation (4), the spectrum $Q(\omega, q)$ can be expressed as

$$Q(\omega, q) = \frac{2}{\pi} \frac{\Theta(\omega, T_{ph}) R_{\text{eff}}}{|Z^{(5)}|^2} \text{Re}[Z^{(5)}].$$  

Here, $\text{Re}[Z^{(5)}]$ is the real part of the wave impedance in indium antimonide semi-infinite layer, and the function $\Theta(\omega, T)$ is the Planck’s mean energy of a harmonic oscillator, in other words

$$\Theta(\omega, T) = \frac{\hbar \omega}{\exp\left(\frac{\hbar \omega}{k_B T}\right) - 1},$$
$$k_B = 1.38 \times 10^{-23} \text{ J K}^{-1}, \quad \hbar = \frac{6.626 \times 10^{-34}}{2\pi} \text{ J s}.$$  

In equation (8), the effective resistance ($R_{\text{eff}}$) is given by

$$R_{\text{eff}} = F^{(4)} F^{(1)} F^{(2)} \text{Re}[Z^{(1)}],$$

where

$$F^{(i)} = \frac{|Z^{(i)}|^2}{|Z^{(i)} \cos(\beta^{(i)} d^{(i)}) + j Z^{(i-1)} \sin(\beta^{(i)} d^{(i)})|^2},$$
$$Z^{(i)}_{m} = \frac{Z^{(i)}_{m-1} + j Z^{(i-1)} \tan(\beta^{(i)} d^{(i)})}{Z^{(i)}_{m} + j Z^{(i-1)} \tan(\beta^{(i)} d^{(i)})}; \quad i = 2, 3, 4,$$
$$Z^{(1)}_{m} = Z^{(1)}.$$

In the above equations, $d^{(2)} = W_e$, $d^{(3)} = W_e$ and $d^{(4)} = W_b$. Also, $Z$ and $\beta$ represent the wave impedance and the propagation factor, respectively, described by

$$\beta^{(1)}_{(2)} = \sqrt{\kappa_{0}^{2} \frac{A^{(1)}_{(2)}}{\gamma^{2}}} = \varepsilon_{r}^{(1)} = \varepsilon_{r}^{(2)} = 1,$$
$$Z^{(1)}_{(2)} = \frac{\beta^{(1)}_{(2)}}{\omega \varepsilon_{0} \varepsilon_{r}^{(1,2)}},$$

for the isotropic layers and

$$\beta^{(3,4)} = \sqrt{\kappa_{0}^{2} \varepsilon_{\perp}^{(3,4)} - \varepsilon_{\parallel}^{(3,4)}} q^2,$$
$$Z^{(3,4)} = \frac{\beta^{(3,4)}}{\omega \varepsilon_{0} \varepsilon_{\perp}^{(3,4)}},$$

for each wire metamaterial layer (layer three is the layer of free-standing nanowires and layer four is the layer where the nanowires are grown in the dielectric material). According to the effective-medium model [33], the transverse ($\varepsilon_{\perp}$) and the axial ($\varepsilon_{\parallel}$) components of the effective permittivity of the wire metamaterial are expressed as

$$\varepsilon_{\parallel}^{(3,4)} = \frac{(1 + f_{e}) \varepsilon_{m} + (1 - f_{e}) \varepsilon_{h}^{(3,4)}}{(1 - f_{e}) \varepsilon_{m} + (1 + f_{e}) \varepsilon_{h}^{(3,4)}},$$
$$\varepsilon_{\perp}^{(3,4)} = f_{e} \varepsilon_{m} + (1 - f_{e}) \varepsilon_{h}^{(3,4)},$$

where $\varepsilon_h$ and $\varepsilon_m$ are the relative dielectric constants of the host medium and the metal (tungsten), respectively.

A.2. Indium antimonide (InSb) properties

At the room temperature ($T = 300$ K), based on Einstein relation, the diffusion coefficient is given by

$$D = 25.875 (\text{mV}) \cdot \mu,$$

in which $\mu$ is the electron/hole mobilities for InSb, described by [43]

$$\mu = \mu_{\text{min}} + \frac{\mu_{\text{max}} - \mu_{\text{min}}}{1 + \frac{N}{N_{\text{ref}}}}.$$

Here, $\mu_{\text{max}}(e) = 7.8 \times 10^{3} \text{ cm}^2 \text{V}^{-1} \text{s}^{-1}$, $\mu_{\text{min}}(e) = 5000 \text{ cm}^2 \text{V}^{-1} \text{s}^{-1}$, $\mu_{\text{max}}(h) = 100 \text{ cm}^2 \text{V}^{-1} \text{s}^{-1}$, $N_{\text{eff}}(e) = 7 \times 10^{16} \text{ cm}^{-3}$, $N_{\text{ref}}(h) = 6 \times 10^{17} \text{ cm}^{-3}$, $\Phi(e) = 0.7$ and $\Phi(h) = 0.6$. In the calculations, we assume that $N_0 = N_p = 10^{19} \text{ cm}^{-3}$. From equations (16) and (17), we achieve $D_e \approx 186 \text{ cm}^2 \text{s}^{-1}$ and $D_h \approx 5.2 \text{ cm}^2 \text{s}^{-1}$.

The thickness of the depletion region is calculated as follows

$$W_0 = \left(\frac{2 \varepsilon_{0} \varepsilon_{\text{InSb}}}{\varepsilon_{r}} \left(\frac{1}{N_n} + \frac{1}{N_p}\right) V_0\right),$$

Table 3. Overall conversion efficiency versus $W_b$.

| $W_b$ (nm) | 300     | 400     | 500     | 600     | 700     | 800     | 900     | 1000    |
|-----------|---------|---------|---------|---------|---------|---------|---------|---------|
| $\eta$%   | 16.3    | 15.7    | 14.8    | 13.9    | 13.2    | 12.7    | 12.4    | 12.1    |
where $V_0 = 25.875$(mV)$\ln[N_d \cdot N_p/n_i^2]$ is the junction built-in potential, $\varepsilon_0$ is the free-space permittivity and $\varepsilon_{\text{InSb}} = 16$ represents the relative dielectric constant of InSb. Therefore, the thickness of the depletion region is estimated to be $W_0 \approx 10$ nm.

The total minority carrier lifetime $\tau$ depends on the nonradiative Shockley–Read–Hall (SRH) and Auger recombination lifetimes and as well the radiative lifetime as

$$\frac{1}{\tau} = \frac{1}{\tau_{\text{SRH}}} + \frac{1}{\tau_{\text{Auger}}} + \frac{1}{\tau_{\text{radiative}}}. \tag{19}$$

Here, the SRH and radiative lifetimes are expressed as [22]

$$\tau_{\text{SRH}} = \frac{1}{\sigma N_i} \sqrt{\frac{m^*}{3k_B T}}, \tag{20}$$

$$\tau_{\text{radiative}} = \frac{1}{B \cdot N},$$

in which $\sigma = 1.5 \times 10^{-19}$ m$^2$, $N_i = 1.17 \times 10^{21}$ m$^{-3}$, $B = 5 \times 10^{-11}$ cm$^3$ s$^{-1}$ and $m^*$ represents the effective mass for InSb ($m^*(e) = 0.13 \times 10^{-31}$ kg, $m^*(h) = 3.91 \times 10^{-31}$ kg). The Auger recombination lifetime is calculated based on the method introduced in [43]. Finally, we achieve $\tau(e) \approx 1.45$ ns and $\tau(h) \approx 1.81$ ns.

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