

Article

High-Sensitivity Bistable Biosensor Based on the Photonic Crystal Fabry-Pérot Cavity with Weyl Semimetal

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Abstract

Optical bistability (OB) with low threshold and high tunability is crucial for advanced photonic devices. This work theoretically investigates low-threshold and tunable OB in a one-dimensional photonic crystal Fabry-Pérot (FP) cavity embedded with a Weyl semimetal (WSM) layer. By combining the local field enhancement of the cavity and the large third-order nonlinear refractive index of the WSM, we achieve OB in the terahertz regime with a low threshold of $\sim 10^5$ V/m. Furthermore, the OB threshold and hysteresis width can be flexibly manipulated via the Fermi energy of WSM, incident angle, photonic crystal period, the refractive index of the dielectric inside the cavity, and the position of the WSM inside the FP cavity. Moreover, based on the high sensitivity of switching thresholds to refractive index and displacement variations, we also propose a dual-parameter sensing scheme. This exhibits excellent sensing performance for gas sensing (sensitivity up to 96.48×10^5 V/m·RIU) and high sensitivity for displacement sensing (0.91×10^6 V/m· μm). We believe that the above results can provide reference schemes for nonlinear photonic devices.

Keywords: optical bistability; weyl semimetal; biosensor; Fabry-Pérot cavity

1. Introduction

Optical bistability (OB) is the hysteresis switching phenomenon where a stable input light intensity could produce two different output light intensities [1]. This distinctive nonlinear response not only enriches the theoretical underpinnings of light-matter interactions but also provides the physical basis for the design and realization of photonic devices, including optical memories [2], optical logic gates [3,4], optical switches [5,6], and high-sensitivity optical sensors [7]. While the physical mechanisms and fundamental theoretical framework of OB have been largely established, the generally high threshold for OB realization remains a major bottleneck, limiting the practical deployment of related devices. Concurrently, the limited flexibility of active control over OB hinders its adaptation to diverse real-world application scenarios. Consequently, achieving low-threshold and tunable OB has emerged as a critical research frontier and development direction in contemporary nonlinear optics. To date, numerous physical mechanisms and device architectures have been demonstrated to effectively reduce the OB threshold, with the



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Fabry-Pérot (FP) cavity being the most prominent example [8,9]. The FP cavity has a simple structure, versatile design, and exceptional local field enhancement capability, which can drastically lower the OB threshold via cavity resonance effects. Beyond the optical resonator cavity, physical mechanisms such as surface plasmon polaritons [10], Fano resonances [11], topological edge states [12], and optical Tamm states [13] can also induce strong local field enhancement, thereby reducing the OB onset threshold.

In addition to the aforementioned local field enhancement achieved through structural and mechanism engineering, exploring novel materials with high nonlinear coefficients represents another important avenue for reducing the OB threshold. In recent years, researchers have successfully realized low-threshold and tunable OB using various materials such as graphene [14,15] and Dirac semimetals [16], providing entirely new insights for the design and optimization of related optical bistability schemes. In recent years, a new class of topological materials, Weyl semimetals (WSMs) [17,18], has attracted extensive attention from researchers. We know that linear dispersion, chiral anomaly, and Fermi arc are the three most important characteristics of WSMs, which are the physical sources of their peculiar optical properties. This also means that WSMs may have optoelectronic properties comparable to graphene and other Dirac materials. It is worth mentioning that Zhang et al. studied in detail the connection between the solutions to the Dirac and Weyl equations and the associated electromagnetic four-potentials [19]. This provides a solid theoretical foundation for us to utilize the unique characteristics of WSMs and apply them to nanophotonics. In the field of nonlinear optics, WSMs also exhibit exceptional properties. Almutairi et al. investigated the four-wave mixing effect in WSMs and discovered that doped bulk WSMs exhibit a third-order nonlinear Hall effect. Even as a lossy material, its nonlinear conversion efficiency is significantly higher than that of conventional nonlinear optical materials [20]. Zhong et al. found that the third-order nonlinear conductivity of WSMs is much larger than that of traditional materials in long-wavelength bands such as the terahertz region [21]. Furthermore, Bednik et al. demonstrated that the presence of a finite magnetic field can substantially enhance the nonlinear optical response of WSMs [22]. These studies collectively confirm that WSMs possess excellent and flexibly tunable nonlinear optical properties, holding tremendous application potential in the design of low-threshold tunable OB devices. Specifically, WSMs have a high nonlinear refractive index, which is similar to that of graphene and Dirac semimetals. However, compared with graphene, WSMs have easier preparation and stable performance. Meanwhile, WSMs can also adjust the dielectric constant by changing the Fermi energy. Therefore, the high nonlinear refractive index of WSMs is expected to play a crucial part in providing strong optical nonlinearity to realize low-threshold OB.

In this work, we introduce a Weyl semimetal (WSM) sheet as the nonlinear dielectric into a one-dimensional (1D) photonic crystal (PhC) structure and construct an FP resonant configuration. By combining the excellent third-order nonlinear optical properties of the WSM with the strong local field enhancement effects of FP resonance, we achieve low-threshold and tunable OB in the terahertz frequency range, with a threshold as low as the order of 10^5 V/m. The OB curve can be conveniently and flexibly tuned by adjusting parameters such as the Fermi energy of the WSM and the incident angle. In addition, we can also realize dynamic modulation of the OB threshold by varying structural parameters, including the number of photonic crystal periods and the position of the WSM within the cavity. The results of this study not only validate the great application potential of the WSM in low-threshold OB devices but also provide important theoretical references for the design and realization of miniaturized, high-density, integrated photonic devices based on OB.

2. Theoretical Model and Method

As shown in Figure 1, a single layer of WSM is placed at the center of the FP cavity, which is symmetrically flanked by 1D PhCs with identical period numbers on both sides. It should be noted that this structure is mainly for illustrative purposes. In the actual case, the WSM sheet can be sandwiched between two dielectric slabs. For example, flexible polymer material, owing to its high transparency in the THz frequency, was chosen to act as the dielectric spacer layer required for the FP cavity. Furthermore, by introducing an air cavity, we ensure that the cavity length L satisfies the FP resonance condition $2n_a L = m\lambda_0$, where the cavity is filled with gas, n_a is the refractive index of the gas inside the cavity, λ_0 is the wavelength in vacuum, and m is the resonance order, which is typically an integer. Unlike the complex preparation process of Dirac materials such as graphene, many WSMs that have been theoretically calculated can be naturally synthesized without the need for detailed and complex control such as doping, which also brings convenience to material preparation. In addition, for the preparation of layered micro/nano structures, there are currently very mature processes and technologies that can support it. Therefore, the physical implementation of the above structure does not have unsolvable problems. Here, we adopt the resonance order $m = 1$. For computational convenience, we assume that the cavity is filled with air, and the refractive index of air n_a is approximated as 1.

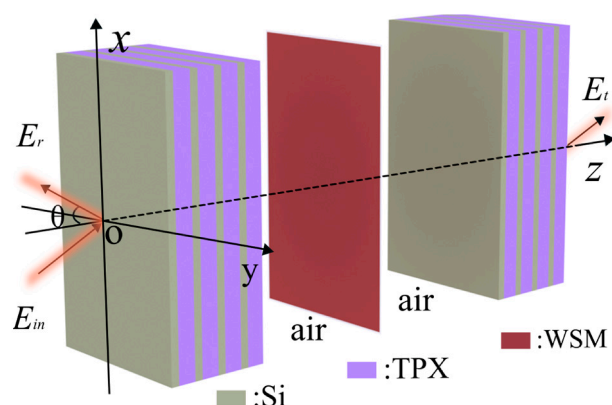


Figure 1. Schematic diagram of a 1D PhCs-FP cavity structure incorporating a WSM layer. The incident light has an amplitude of E_{in} , and the amplitudes of the generated reflected and transmitted light are E_r and E_t , respectively.

The left and right 1D PhCs are composed of alternating layers of two dielectrics, silicon (Si) and polymethylpentene (TPX). Consequently, the entire structure can be denoted as $(P_1 P_2)^N (A) (WSM) (A) (P_1 P_2)^N$, where P_1 and P_2 represent the Si layer and TPX layer, respectively, A denotes the cavity section, and N is the number of photonic crystal periods. The refractive indices of Si and TPX are $n_{P1} = 3.42$ and $n_{P2} = 1.45$, respectively. In addition, the thicknesses of the P_1 and P_2 layers, d_{P1} and d_{P2} , satisfy Bragg's law with a Bragg wavelength $\lambda_0 = 50 \mu\text{m}$.

We know that the massless properties of WSM quasiparticles significantly enhance and reshape the electrical conductivity characteristics of the material through ultra-high mobility, negative magnetic resistance caused by chiral anomalies, and giant anomalous Hall effect driven by intrinsic Berry curvature. These provide various manipulation methods for the conductivity characteristics of WSMs. In order to simplify the calculation model, we will incorporate the material properties of WSMs into the conductivity model in the following calculations. Under certain conditions, the linear conductivity of WSMs can be expressed as [23]:

$$\sigma^{(1)}(\omega) = \frac{e^2 v_F P_F^2 g_s g_\omega}{6\pi^2 \hbar^3 (\gamma - i\omega)} \quad (1)$$

where e denotes the elementary charge, the Fermi velocity is $v_F = 10^6$ m/s, $P_F = E_F/v_F$ represents the magnitude of the Fermi momentum, and E_F is the Fermi energy of WSM. g_s is the spin degeneracy factor, and g_ω is the Weyl node degeneracy factor. Here, we adopt $g_s = 2$ and $g_\omega = 4$. $\hbar$ is the reduced Planck constant, γ is the electron relaxation rate with $\hbar\gamma = 5$ meV, and ω is the angular frequency of the incident light. In addition, the WSM is considered an isotropic material in this work, and its linear permittivity is given by:

$$\varepsilon_1 = \varepsilon_b + \frac{i4\pi\sigma^{(1)}}{\varepsilon_0\omega} \quad (2)$$

where the background permittivity $\varepsilon_b = 10$, and ε_0 is the vacuum permittivity. The third-order nonlinear conductivity of WSM can be expressed as [23]:

$$\sigma^{(3)}(\omega) = \frac{e^4 v_F g_s g_\omega}{15\pi^2 \hbar^3 (-i\omega + \gamma)} \left[\frac{2}{(i\omega + \gamma)(-i\omega + \gamma)} + \frac{1}{(-i2\omega + \gamma)(-i\omega + \gamma)} \right] \quad (3)$$

The third-order nonlinear susceptibility of WSM is given by:

$$\chi^{(3)} = \frac{i\sigma^{(3)}}{\varepsilon_0\omega}. \quad (4)$$

The refractive index of WSM includes both linear and nonlinear components, which can be written as:

$$n_W = \sqrt{\varepsilon_1 + \chi^{(3)}|E_W|^2} \quad (5)$$

where E_W denotes the electric field intensity in the WSM layer, which varies with the position of the WSM layer. There are many options for WSMs that meet the above characteristics, such as TaAs, TaP, NbAs, and NbP, etc. Taking TaAs as an example, its preparation method is simple, and its nonlinear optical properties have also received sufficient attention.

For layered structures, we can use the transfer matrix method to determine the relationships between the incident, reflected, and transmitted electric fields. In this work, we assume that light propagates along the z -axis, and the 1D PhCs and WSM layer are parallel to the x - y plane. Using the transfer matrix method, we can also calculate the transmittance and reflectance of the composite structure. Here, we only consider the case of TE polarization. Although the calculation methods for the two polarizations are similar, the OB calculation for TE polarization only needs to consider the electric field in the y direction, while the OB calculation for TM polarization needs to consider both the x and z directions simultaneously. Under TE polarization, the transfer matrix for a single layer can be expressed as:

$$M_j = \begin{pmatrix} \cos(k_{jz}d_j) & -\frac{i}{\eta_j}\sin(k_{jz}d_j) \\ -i\eta_j\sin(k_{jz}d_j) & \cos(k_{jz}d_j) \end{pmatrix} \quad (6)$$

where M_j represents the transfer matrix for the P_1 , P_2 , or air layer. The z -component of the wave vector in the j -th layer is $k_{jz} = n_j \frac{\omega}{c} \cos(\theta_j)$, and $\eta_j = \sqrt{\varepsilon_0/\mu_0} n_j \cos(\theta_j)$ denotes the optically corrected admittance under TE polarization, where n_j is the refractive index of the j -th layer, θ_j is the refraction angle, and d_j is the thickness of that layer.

As can be seen from Equation (5), the refractive index of WSM is highly dependent on the local electric field intensity, making the standard transfer matrix method inapplicable. Hence, we adopt a layer-by-layer approximation approach, dividing the WSM layer into n sufficiently thin sublayers along the propagation direction. In this case, the electric field intensity can be approximated as the electric field value at the right boundary of each sublayer. Therefore, the transfer matrix for each sublayer is identical to the ordinary transfer matrix, and the transfer matrix corresponding to the n -th sublayer of the WSM is M_n . The

total transfer matrix for the entire WSM layer is obtained by multiplying the transfer matrices of the n sublayers: $M_W = M_1 M_2 M_3 \cdots M_{n-2} M_{n-1} M_n$. Finally, the total transfer matrix of the entire structure can be expressed as $M = (M_{P_1} M_{P_2})^N M_A M_W M_A (M_{P_1} M_{P_2})^N$:

After obtaining the total transfer matrix, we can calculate the transmission coefficient and the reflection coefficient of the entire structure using the following equation:

$$t = \frac{2\eta_0}{M_{11}\eta_0 + M_{12}\eta_0\eta_N + M_{21} + M_{22}\eta_N}, r = \frac{M_{11}\eta_0 + M_{12}\eta_0\eta_N - M_{21} - M_{22}\eta_N}{M_{11}\eta_0 + M_{12}\eta_0\eta_N + M_{21} + M_{22}\eta_N} \quad (7)$$

where η_0 and η_N are the corrected admittances of the external environments in contact with the left and right sides of the structure, respectively. Finally, we can further obtain the transmittance and reflectance of the entire structure as: $R = |r|^2$, $T = |t|^2$, $A = 1 - R - T$. We can also express the relationship between the incident electric field and the transmitted electric field using:

$$|E_t| = |E_{in}| \cdot |t|. \quad (8)$$

3. Results and Discussion

In this section, we investigate the OB in the terahertz band. First, we analyze the relevant characteristics of OB using an analytical approach. For a standard FP cavity, the light intensity at the position of the WSM layer can be approximately calculated using the following equation:

$$I_W \approx \frac{1 + R_{Phcs}}{1 - R_{Phcs}} I_t \quad (9)$$

where R_{PhC} is the reflectance of the 1D PhCs, I_t is the transmitted light intensity, and I_W is the light intensity near the WSM. Traditional OB phenomena employ optical Kerr media with nonlinear refractive indices. Therefore, we express the nonlinearity of WSM in terms of refractive index $n = n_0 + n_2 I_W$. Where n_0 is the linear refractive index of WSM, and n_2 is the nonlinear refractive index of WSM. n_0 is the real part of the linear complex refractive index N , which can be calculated by $N = n_0 + ik = \sqrt{\epsilon_1}$ [24]. Where k is the extinction coefficient. Thus, $n_0 = \text{Re}(N)$ and $k = \text{Im}(N)$. For the nonlinear refractive index n_2 , we can calculate it using [24]:

$$n_2 = \frac{3}{4\epsilon_0 c (n_0^2 + k^2)} \left[\text{Re}(\chi^{(3)}) + \frac{k}{n_0} \text{Im}(\chi^{(3)}) \right] \quad (10)$$

When the absorption is weak ($k \ll n_0$), the above equation can be simplified to:

$$n_2 \approx \frac{3\text{Re}(\chi^{(3)})}{4\epsilon_0 c n_0^2} \quad (11)$$

Obviously, under the conditions of this work, we can approximately neglect the influence of the extinction coefficient k on the nonlinear refractive index. Substituting Equation (9) into $n = n_0 + n_2 I_W$, we can obtain:

$$n = n_0 + n_2 \frac{1 + R_{Phc}}{1 - R_{Phc}} I_t \quad (12)$$

In Equation (12), we denote $(1 + R_{Phc})/(1 - R_{Phc})$ as the field enhancement factor α . The cavity length of the FP cavity is L , so the phase difference ϕ for a round-trip propagation of light inside the cavity is:

$$\phi = \frac{2\pi}{\lambda} n(2L) = \frac{4\pi}{\lambda} nL. \quad (13)$$

where λ is the wavelength of the incident light. Substituting Equation (12) into Equation (13), we can obtain:

$$\phi = \frac{4\pi}{\lambda}(n_0 + n_2\alpha I_t)L = \frac{4\pi}{\lambda}Ln_0 + \frac{4\pi}{\lambda}L\alpha n_2 = \phi_0 + \phi_2 I_t. \quad (14)$$

where ϕ_0 is the initial phase shift with $\phi_0 = \frac{4\pi}{\lambda}Ln_0$, and ϕ_2 is a constant related to the nonlinear refractive index with $\phi_2 = \frac{4\pi}{\lambda}L\alpha n_2$. From the FP resonance condition $2n_aL = m\lambda_0$, ϕ_0 can be simplified to $2\pi n_0$ and ϕ_2 can be simplified to $2\pi\alpha n_2$. From Equation (14), we can derive the relationship between the phase shift ϕ and the transmittance T : $T = I_t/I_{in} = (\phi - \phi_0)/\phi_2 I_t$. Obviously, the relationship between T and ϕ is a straight line with a slope dependent on the reciprocal of the incident light intensity. This equation is also known as the feedback relation of OB. According to the characteristics of FP interferometers, we can also express the periodic relationship between the phase shift ϕ and the transmittance T using:

$$T = \frac{I_t}{I_{in}} = \frac{1}{1 + \frac{4R_{Phc}}{T_{Phc}^2} \sin^2 \frac{\phi}{2}} \quad (15)$$

where T_{Phc} is the transmittance of the 1D PhCs. Clearly, T reaches its peak value of 1 when $\phi = 2m\pi$ ($m = 1, 2, 3, \dots$). This equation is also known as the modulation relation of OB.

As shown in Figure 2a, both T_{Phc} and R_{Phc} can be calculated using the transfer matrix method. Clearly, we observe a distinct photonic band gap in the frequency range of 5–7 THz, where the reflectance approaches unity, and the absorptance A is extremely low and can be approximately neglected. Here, for the convenience of observing the modulation curve, we take $T_{PhCs} = 0.1$ and $R_{PhCs} = 0.9$. As depicted in Figure 2b, we plot the modulation curve of OB. The transmission peaks have a period of 2π and a full width at half maximum of $\delta\phi$.

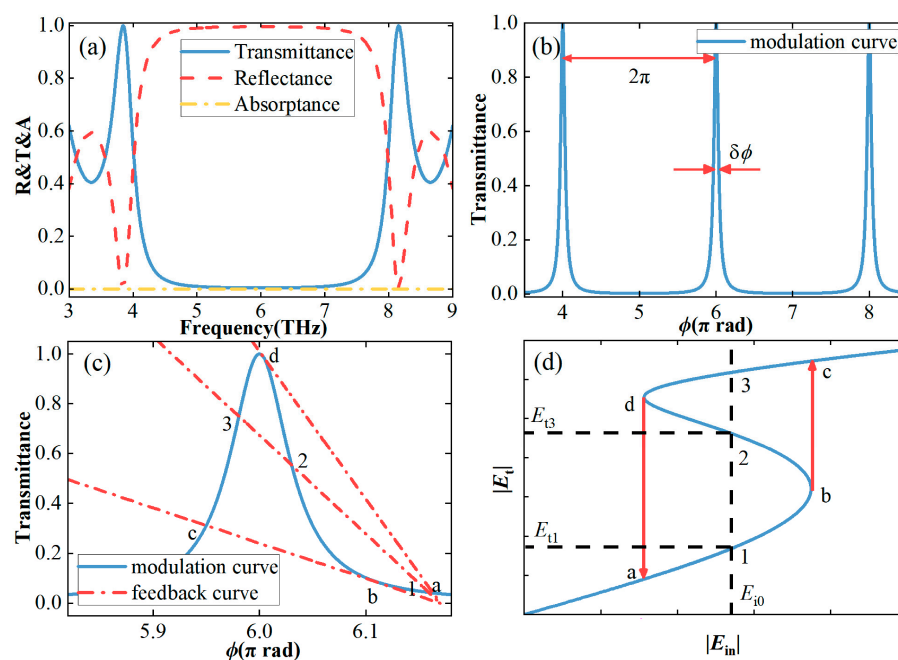


Figure 2. (a) Transmittance, reflectance, and absorptance spectra of the 1D PhCs; (b) Modulation curve of OB; (c) Graphical determination of the OB operating point; (d) OB characteristic curve.

By combining Equations (14) and (15), we employ the graphical method to identify the OB operating points where the two curves intersect, as shown in Figure 2c. Due to

the negative nonlinear refractive index of the WSM, we observe a descending feedback curve. Here, we use electric field instead of light intensity. As the incident electric field gradually increases, the slope of the straight line decreases, and the intersection points of the two curves are a-1-b-c in sequence; as the incident electric field gradually decreases, the intersection points become c-3-d-a. In Figure 2d, we change the coordinates to the incident electric field amplitude $|E_{in}|$ and the transmitted electric field amplitude $|E_t|$. Connecting the corresponding operating points in order, we obtain an S-shaped OB curve: a-1-b-2-d-3-c. Obviously, the OB curve exhibits upward and downward jumps at points b and d, respectively. Within the ranges of lines bc and da, a single incident electric field E_{i0} corresponds to two stable transmitted electric fields E_{t1} and E_{t3} , which constitute the bistable operating range of the OB device.

Subsequently, we analyze the formation conditions of OB using an analytical approach. For Equation (15), since OB occurs near the peaks where $\phi = 2m\pi$ ($m = 1, 2, 3, \dots$), the sine function can be approximated as $\sin(\phi/2) \approx \phi/2$, and the equation can be rewritten as:

$$I_{in} \approx \left(1 + \frac{R_{Phc}\phi^2}{T_{Phc}^2} \right) I_t \quad (16)$$

Let $\frac{R_{Phc}}{T_{Phc}^2} = k$. Substituting Equation (14) into Equation (16) yields:

$$I_{in} = k\phi_2^2 I_t^3 + 2k\phi_0\phi_2 I_t^2 + (1 + k\phi_0^2) I_t. \quad (17)$$

In this work, we assume that both the incident and exit media are vacuum. Therefore, we can relate the electric field and light intensity using:

$$I = \frac{1}{2} c \epsilon_0 |E|^2 \quad (18)$$

Among them, c represents the speed of light, ϵ_0 represents the dielectric constant in vacuum. Let $X = |E_t|^2$, $Y = |E_{in}|^2$, and $a = \frac{1}{2} c \epsilon_0$. We rewrite the relationship between light intensities in terms of electric fields:

$$Y = k\phi_2^2 a^2 X^3 + 2k\phi_0\phi_2 a X^2 + (1 + k\phi_0^2) X. \quad (19)$$

Clearly, the relationship between X and Y is a cubic curve. Its derivative is given by $\frac{dY}{dX} = 3k\phi_2^2 a^2 X^2 + 4k\phi_0\phi_2 a X + (1 + k\phi_0^2)$. The occurrence of OB requires the derivative to have two distinct roots. We can verify the root distribution by calculating the discriminant Δ of the derivative $\Delta = 4k\phi_2^2 a^2 (k\phi_0^2 - 3)$. Within the frequency range investigated in this work, the condition $k\phi_0^2 - 3 > 0$ holds, thus $\Delta > 0$, which confirms the existence of OB in the WSM-based 1D PhCs-FP cavity. The parameter k depends on the reflectance and transmittance of the 1D PhCs, and the photonic band gap generated by the PhC results in an extremely large value of k , making the realization condition of OB much easier to satisfy.

Subsequently, we employ an accurate numerical algorithm to further investigate the OB phenomenon in the proposed structure. Based on the transfer matrix method, we obtain the total transmittance, reflectance, and absorptance of the composite structure, as shown in Figure 3. To simplify the calculations, we neglect the contribution of the third-order nonlinear conductivity of the WSM to the results. In Figure 3a,b, we first investigate the transmission and reflection spectra of the 1D PhC FP cavity without WSM in the frequency range of 3–8 THz. Here, the number of photonic crystal layers N is set to 4, and the incident angle is set to 0° . We observe distinct photonic band gaps in the ranges of 4.3–5.1 THz and 5.3–7.5 THz for the 1D PhC FP cavity. At a frequency of 5.2 THz, we observe clear

transmission and reflection peaks. After introducing the WSM layer, we observe a relatively weak transmission peak and a more pronounced reflection peak at 7.01 THz, and the peak intensities decrease with increasing number of photonic crystal layers. In Figure 3c, we observe a significant absorption peak at 7.01 THz. In the terahertz band, the optical absorption of the entire structure is mainly attributed to the WSM layer. When the number of 1D PhC layers $N = 4$, the absorption peak reaches 0.61 and decreases significantly with increasing N . This is because the increase in the thickness of the 1D PhC leads to an increase in the effective cavity length of the FP cavity, causing a redshift of the resonance frequency and ultimately weakening the optical field inside the cavity. Based on the above discussion, we select $f = 7.01$ THz as the operating frequency.

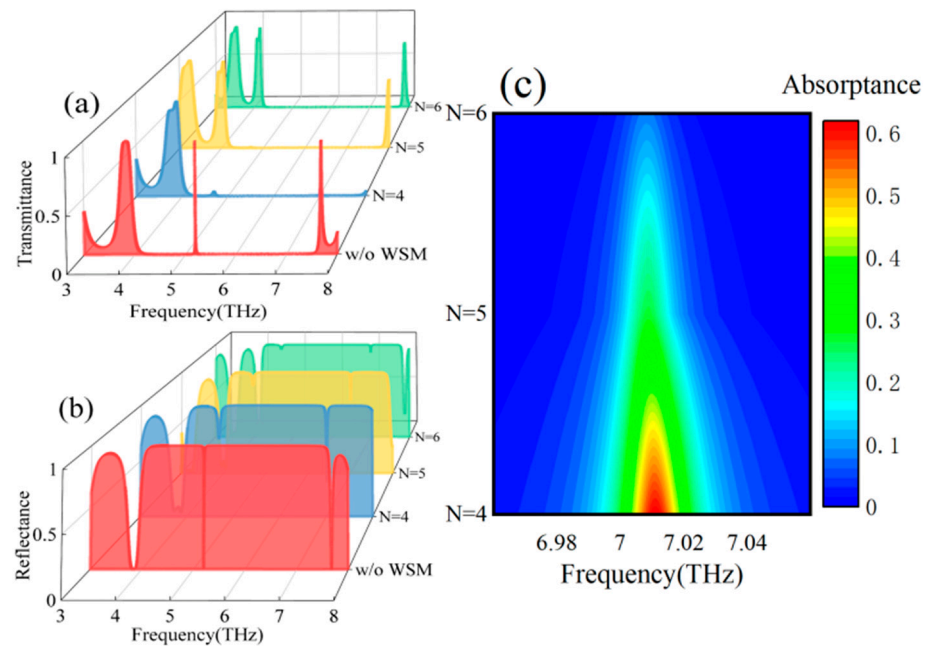


Figure 3. Optical responses of the structures at $E_F = 3$ meV. (a) Reflectance; (b) Transmittance; (c) Absorbance contour maps for $N = 4, 5, 6$ (bottom to top).

On this basis, we further investigate the nonlinear relationships between the incident electric field amplitude and the transmitted electric field amplitude, structural reflectance, and absorbance, as well as the influence of relevant parameters on these curves. First, since the nonlinear response of WSM is dominated by Weyl fermions near the Fermi energy E_F , we investigate the effect of E_F on the OB curves, as shown in Figure 4. In Figure 4a, we plot the relationship curves between the incident electric field amplitude $|E_{in}|$ and the transmitted electric field amplitude $|E_t|$ for different E_F values. Taking the red curve in the figure as an example, in the initial stage, $|E_t|$ increases linearly with increasing $|E_{in}|$. When $|E_{in}|$ increases to the upward threshold $|E_{in}|_{up}$, $|E_t|$ abruptly jumps to a larger value and then decreases slowly with further increase of $|E_{in}|$. When $|E_{in}|$ gradually decreases, $|E_t|$ first increases until it reaches the downward threshold $|E_{in}|_{down}$, where $|E_t|$ jumps to a lower value. We calculate that the hysteresis loop width $\Delta |E_{in}| = |E_{in}|_{up} - |E_{in}|_{down} = 0.38 \times 10^5$ V/m when $E_F = 3$ meV. Clearly, a decrease in E_F leads to a significant rightward shift of the OB curves. To more intuitively demonstrate the effect of E_F on the bistability thresholds, we plot the variations of $|E_{in}|_{up}$ and $|E_{in}|_{down}$ with E_F , as shown in Figure 4b, and the variation of $\Delta |E_{in}|$ with E_F , as shown in the inset of Figure 4b. As E_F decreases from 3 meV, $|E_{in}|_{up}$ increases significantly while $|E_{in}|_{down}$ remains nearly unchanged, which well reflects the trend observed in Figure 4a.

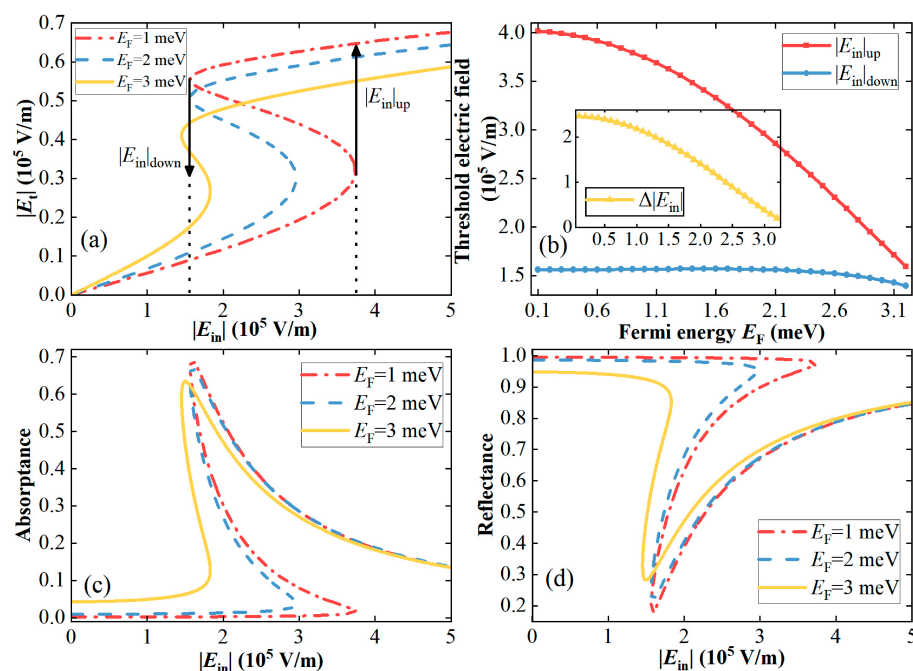


Figure 4. Effect of WSM Fermi energy E_F on OB. (a) Transmitted electric field OB curves; (b) Thresholds vs. E_F (inset: threshold difference $\Delta|E_{in}|$ vs. E_F); (c) Absorbance vs. $|E_{in}|$; (d) Reflectance vs. $|E_{in}|$.

Furthermore, we also investigate the relationship curves between the incident electric field amplitude $|E_{in}|$ and the absorbance A and reflectance R for different E_F values, as shown in Figure 4c,d. In Figure 4c, we observe a similar trend to that in Figure 4a. In comparison, the jumps in absorbance A and reflectance R are more pronounced, with the absorption peak reaching approximately 0.68 and the reflection valley as low as 0.2. We also find that when E_F is greater than 3.5 meV, the real part of the linear conductivity of the WSMs increases significantly, which enhances the absorption of incident light by WSMs. At this point, it becomes difficult to observe a clear OB phenomenon. The above discussion demonstrates that manipulating the Fermi level E_F of WSMs is an important means for tuning OB in practical applications.

In subsequent studies, we fix $E_F = 3$ meV. On this basis, we further investigate the tuning effect of the incident angle θ on OB, as shown in Figure 5. As mentioned earlier, all discussions in this paper are based on incident light under TE polarization. In Figure 5a, we observe that with increasing incident angle θ , both the upward threshold $|E_{in}|_{up}$ and downward threshold $|E_{in}|_{down}$ of OB shift toward lower electric fields. As shown in Figure 5b, when θ increases to 5° , the switching threshold can even reach the order of 10^4 V/m, approximately 0.93×10^4 V/m. Similarly, the hysteresis loop width $\Delta|E_{in}|$ also decreases with increasing θ , and the values of $|E_{in}|_{up}$ and $|E_{in}|_{down}$ tend to converge. The same conclusion applies to the absorbance OB and reflectance OB in Figure 5c,d. It should be noted that when the incident angle θ continues to increase beyond 5° , the S-shaped characteristic of the OB curve will gradually weaken until the hysteresis phenomenon disappears. Therefore, we should control the incident angle θ to not exceed 5° .

In practical device fabrication, the structural parameters of the cavity are also an important means for tuning the OB curves. We investigate the effect of the period number N of the 1D PhC on OB, as shown in Figure 6. Clearly, the OB curves are extremely sensitive to the period number N . In Figure 6a, we observe that with increasing period number N , both the upward threshold $|E_{in}|_{up}$ and downward threshold $|E_{in}|_{down}$ of OB shift toward higher electric fields, and the transmitted electric field $|E_t|$ corresponding

to the thresholds also decreases significantly. Clearly, the hysteresis loop width $\Delta|E_{in}|$ also increases significantly. The above trends can be intuitively observed in Figure 6b. Furthermore, in Figure 6c,d, we also find that with increasing period number N of the 1D PhCs, the maximum value of absorptance A decreases significantly, while the minimum value of reflectance R increases significantly. This is not only because the increase in N leads to a change in the effective cavity length of the FP cavity, but also because more incident light is reflected at the surface of the 1D PhCs, reducing the coupling efficiency of the incident light and preventing further enhancement of the field enhancement effect inside the cavity. Therefore, in practical applications, we should set the period number of the 1D PhCs to 4. Similarly, this also indicates that tuning the period number N of the 1D PhCs can effectively manipulate the OB curves.

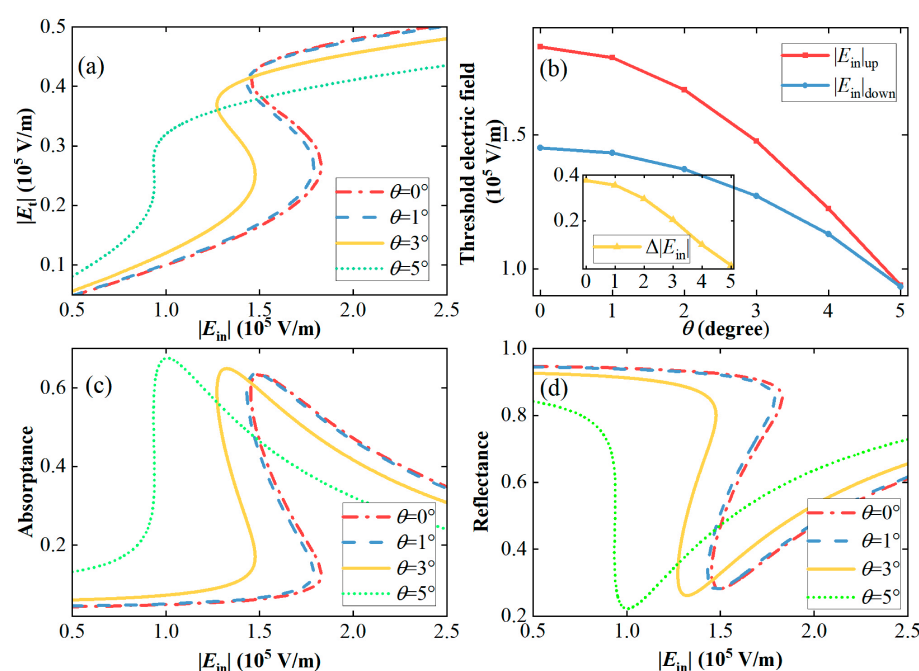


Figure 5. Effect of incident angles θ on OB. (a) Transmitted electric field OB curves; (b) Thresholds vs. N (inset: threshold difference $\Delta|E_{in}|$ vs. θ); (c) Absorptance vs. $|E_{in}|$; (d) Reflectance vs. $|E_{in}|$.

In addition to the period number N , the refractive index n_a of the gas filled in the cavity and the position of the WSM layer within the cavity are also important structural parameters. Therefore, we further investigate their tuning effects on the OB curves, as shown in Figure 7. In Figure 7a–c, we observe that with increasing n_a , the switching thresholds of OB also shift toward higher electric fields, and the corresponding hysteresis loop width increases. Furthermore, we find that if the cavity is filled with liquid, i.e., the refractive index continues to increase, a clear OB phenomenon can still be observed. However, correspondingly, the order of magnitude of the switching threshold will increase to 10^6 V/m, which provides a new perspective for our structural design. In Figure 7d–f, we investigate the effect of different positions of the WSM layer within the cavity on OB. We define the leftward displacement distance of the WSM layer as Δz , and $\Delta z = 0$ μm when the WSM layer is located at the center of the cavity. Clearly, with increasing Δz , both the upper and lower thresholds shift significantly toward higher electric fields, and the value of the dependent variable corresponding to the downward threshold increases considerably. When $\Delta z = 1$ μm , the order of magnitude of the upper and lower thresholds jumps to 10^6 V/m.

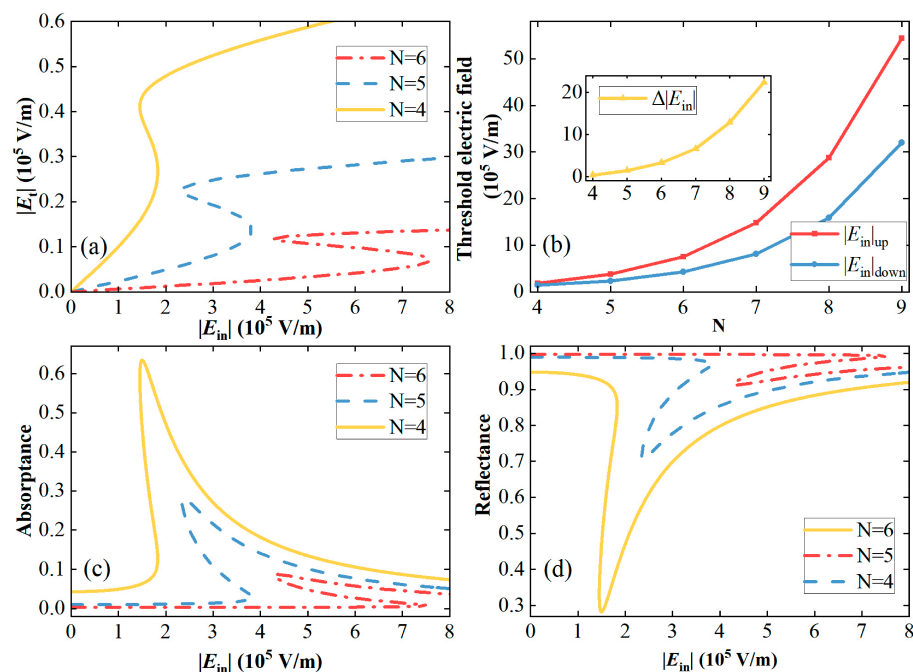


Figure 6. Effect of the number of 1D PhCs periods N on OB. (a) Transmitted electric field OB curves; (b) Thresholds vs. N (inset: threshold difference $\Delta|E_{in}|$ vs. N); (c) Absorbance vs. $|E_{in}|$; (d) Reflectance vs. $|E_{in}|$.

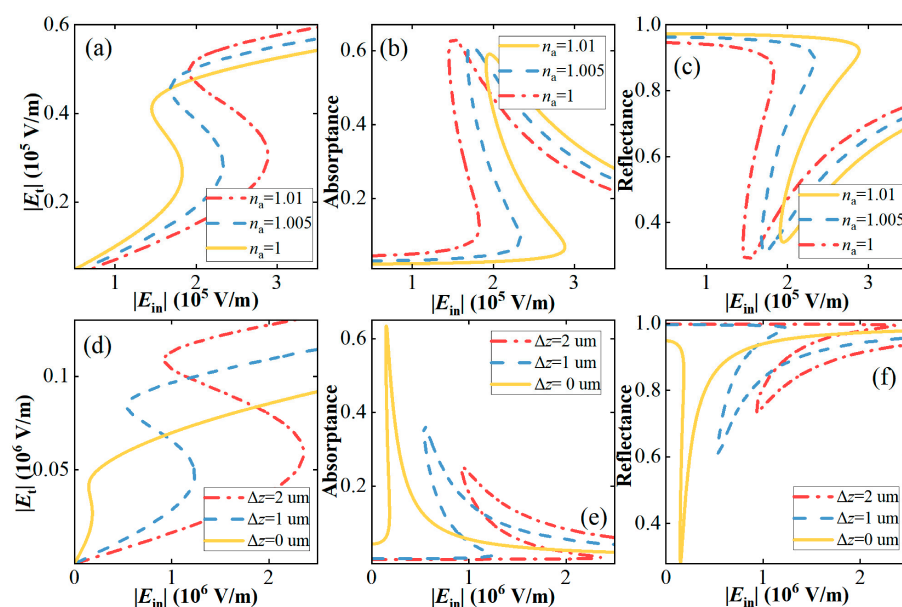


Figure 7. Effects of gas refractive index n_a in the cavity and leftward displacement Δz of the WSM layer on the OB characteristics. (a) Transmitted electric field OB curves for different n_a ; (b) Absorbance vs. $|E_{in}|$; (c) Reflectance vs. $|E_{in}|$; (d) Transmitted electric field OB curves for different Δz ; (e) Absorbance vs. $|E_{in}|$; (f) Reflectance vs. $|E_{in}|$.

In addition, it should be noted that when the WSM layer moves toward the right side of the cavity, it is difficult to observe the OB phenomenon. Obviously, the movement of the WSM layer will alter the distribution of the standing wave field inside the cavity, leading to the shift of the OB curves.

By examining the curves in Figure 7, we can readily find that a small change in structural parameters can induce a relatively large variation in the upward threshold $|E_{in}|_{up}$ and downward threshold $|E_{in}|_{down}$. As shown in Figure 8a, when the refractive index n_a of the gas in the cavity is 1 and 1.005, respectively, the refractive index difference

$\Delta n_a = 0.005$, corresponding to an upward threshold difference $\Delta |E_{in}|_{up} = 0.51 \times 10^5$ V/m and a downward threshold difference $\Delta |E_{in}|_{down} = 0.23 \times 10^5$ V/m between the two OB curves. As shown in Figure 8f, when the leftward displacement Δz of the WSM layer is 1 μm and 2 μm , respectively, the displacement difference $\Delta z = 1 \mu\text{m}$, corresponding to an upward threshold difference $\Delta |E_{in}|_{up} = 1.12 \times 10^6$ V/m and a downward threshold difference $\Delta |E_{in}|_{down} = 0.38 \times 10^6$ V/m between the two OB curves. This phenomenon provides a novel approach for the design of high-sensitivity sensors in the terahertz band.

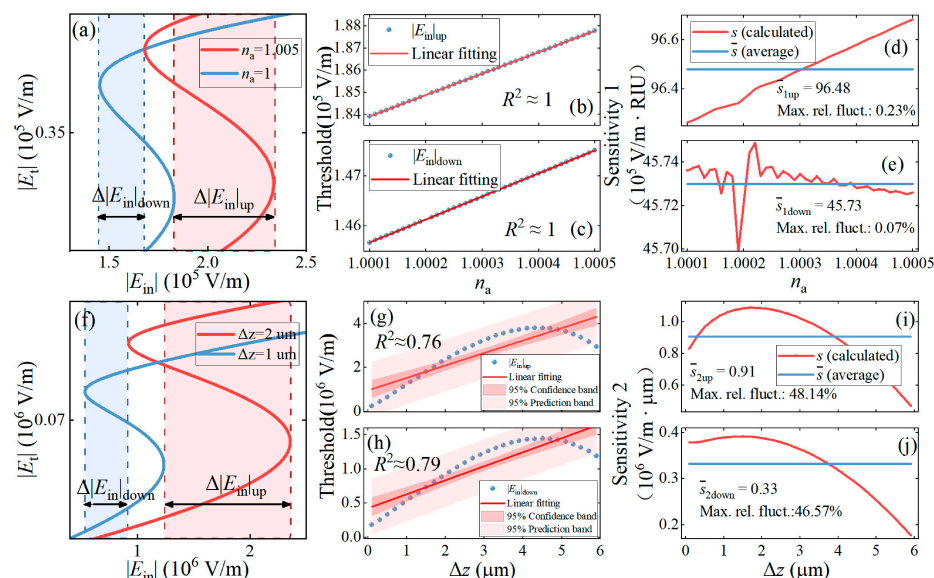


Figure 8. Sensing performance of the optical bistable sensor. (a) Transmission electric field bistability curves at $n_a = 1$ and $n_a = 1.005$, showing the variations of the upper and lower thresholds; (b) $|E_{in}|_{up}$ and (c) $|E_{in}|_{down}$ vs. n_a with their linear fittings; (d) S_{1up} and (e) S_{1down} vs. n_a and average sensitivity; (f) Transmission electric field bistability curves at $\Delta z = 1 \mu\text{m}$ and $\Delta z = 2 \mu\text{m}$, showing the variations of the upper and lower thresholds; (g) $|E_{in}|_{up}$ and (h) $|E_{in}|_{down}$ vs. Δz with their linear fittings, 95% confidence band, and 95% prediction band; (i) S_{2up} and (j) S_{2down} vs. Δz and average sensitivity.

Based on the above phenomena, we can propose two types of refractive index sensors and displacement sensors based on the modulation of switching threshold electric fields. First, we define the sensitivities as follows:

$$S_{1up} = \frac{\Delta |E_{in}|_{up}}{\Delta n_a}, S_{1down} = \frac{\Delta |E_{in}|_{down}}{\Delta n_a}, S_{2up} = \frac{\Delta |E_{in}|_{up}}{\Delta z}, S_{2down} = \frac{\Delta |E_{in}|_{down}}{\Delta z} \quad (20)$$

In Equation (20), $\Delta |E_{in}|_{up}$ is the upward threshold difference of the OB curves under different refractive indices, and $\Delta |E_{in}|_{down}$ is the downward threshold difference of the OB curves under different refractive indices. Δn_a is the change in the refractive index of the gas filled in the cavity. Here, we set the reference refractive index of the gas in the cavity $n_{a0} = 1.0002$.

Furthermore, $\Delta |E_{in}|_{up}$ is the upward threshold difference of the OB curves under different displacements, and $\Delta |E_{in}|_{down}$ is the downward threshold difference of the OB curves under different displacements. Δz is the leftward displacement of the WSM layer relative to the central position.

We first discuss the refractive index sensor. Considering common gases under standard conditions, we limit the research range of the refractive index n_a to 1.0001–1.0005. As shown in Figure 8b,c, we extract the values of $|E_{in}|_{up}$ and $|E_{in}|_{down}$ in this range and plot them as scatter plots as a function of n_a . It can be observed that all data points of the upper and lower thresholds exhibit a nearly perfect linear distribution over the entire refractive

index variation range. We further perform linear regression analysis on them using the least squares method, and the results show that the coefficient of determination R^2 of the fits is approximately 1. This result indicates that there exists a nearly ideal linear response relationship between the thresholds and the gas refractive index within the limited range. The nearly perfect linearity means that the sensor has a nearly constant sensitivity over the entire dynamic range, eliminating the need for complex nonlinear calibration algorithms. This not only reduces the complexity and cost of subsequent signal processing circuits but also mitigates the systematic errors introduced by nonlinear calibration, significantly improving the measurement accuracy and reliability.

In Figure 8d,e, we further plot the variations of sensitivities S_{1up} and S_{1down} with different n_a , as shown by the red curves, and the average sensitivities $\bar{S}_{1up}$ and $\bar{S}_{1down}$ within the limited range, as shown by the blue curves. The calculation results show that the average sensitivity of the upward threshold $\bar{S}_{1up}$ reaches 96.48×10^5 V/m·RIU, with a maximum relative fluctuation of only 0.23% over the entire measurement interval; the average sensitivity of the downward threshold $\bar{S}_{1down}$ is 45.73×10^5 V/m·RIU, with a maximum relative fluctuation as low as 0.07%. This result confirms the linear fitting analysis we conducted earlier. Although the sensitivities are approximately constant, as shown in Figure 8d,e, S_{1up} increases slightly with increasing refractive index within the limited range, while S_{1down} is more sensitive to gases, with n_a to the right of 1.0002 and less sensitive to those to the left.

Next, we investigate the displacement sensor. As discussed earlier, we found that the upper and lower thresholds of OB are extremely sensitive to small displacements of the nonlinear layer. As shown in Figure 8g,h, we first extract the values of $|E_{in}|_{up}$ and $|E_{in}|_{down}$ for Δz in the range of 1–6 μm and plot them as scatter plots as a function of Δz . Clearly, the switching thresholds first increase and then slightly decrease with increasing leftward displacement. Linear regression analysis using the least squares method shows that the coefficient of determination R^2 for the upward threshold is approximately 0.76, and that for the downward threshold is approximately 0.79. Clearly, the linearity of the displacement sensor is lower than that of the refractive index sensor. The difference in linearity arises from the fundamental differences between the two sensing mechanisms: changes in gas refractive index uniformly alter the effective optical length of the entire FP cavity, resulting in a strictly linear threshold drift; in contrast, the displacement of the WSM layer not only changes the effective cavity length but, more importantly, modifies the optical field intensity at the position of the WSM layer, leading to a different introduction mode of the nonlinear component. Nevertheless, R^2 values above 0.76 still meet the linearity requirements for most displacement sensing applications.

We further calculate the sensitivity of the displacement sensor, as shown in Figure 8i,j. The average sensitivity of the upward threshold $\bar{S}_{2up}$ reaches 0.91×10^6 V/m· μm , and the average sensitivity of the downward threshold $\bar{S}_{2down}$ is 0.33×10^6 V/m· μm . Unlike gas refractive index sensing, the sensitivity of displacement sensing exhibits significant fluctuations, with a maximum relative fluctuation of approximately 48.14% for the upward threshold and 46.57% for the downward threshold. Furthermore, the variation trends of sensitivity with displacement Δz are also slightly different. As Δz increases, both the upper and lower thresholds first increase and then decrease. The sensitivity reaches its maximum value at $\Delta z = 2 \mu\text{m}$. The two sensing modes possess different response characteristics and complement each other, greatly expanding the application scope of the device.

Finally, we investigate the effects of relevant parameters on the sensor sensitivity, as shown in Figure 9. Based on the results shown in Figure 8, we specify the refractive index variation $\Delta n_a = 0.0003$ and the displacement $\Delta z = 2 \mu\text{m}$. In Figure 9a,b,e,f, we investigate

the effect of the Fermi level E_F of WSM on the sensitivity. In Figure 9c,d,g,h, we investigate the effect of the period number N of the 1D PhC on the sensitivity.

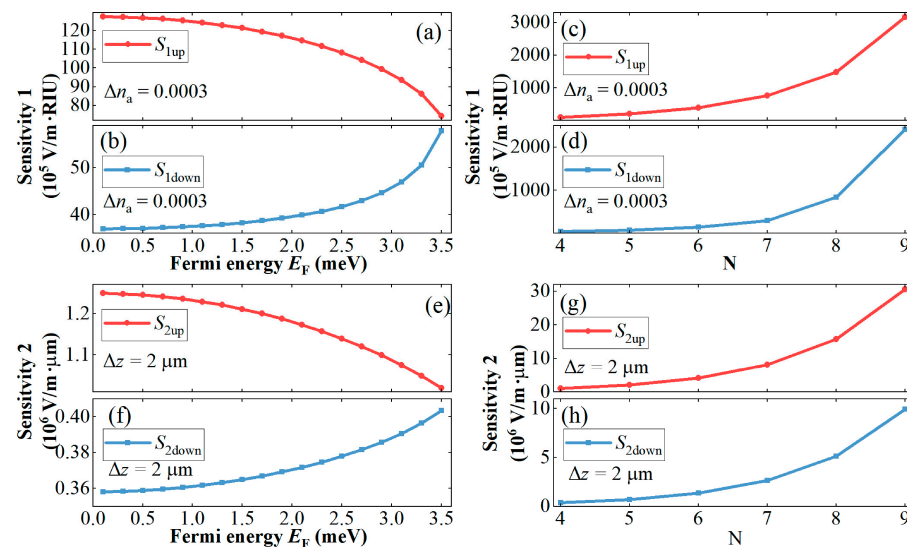


Figure 9. Influence of relevant parameters on sensitivity. (a) S_{1up} (b) S_{1down} (e) S_{2up} and (f) S_{2down} vs. WSM Fermi energy E_F . (c) S_{1up} (d) S_{1down} (g) S_{2up} and (h) S_{2down} vs. 1D PhC periods N .

Clearly, the variation trends of sensitivity with these parameters are consistent. As E_F increases, S_{1up} and S_{2up} decrease, while S_{1down} and S_{2down} increase. As N increases, S_{1up} , S_{1down} , S_{2up} , and S_{2down} all increase. This result well reflects the tuning effects of E_F in Figure 4 and N in Figure 6 on the OB curves. The above studies confirm the feasibility of the dual-parameter sensing scheme based on OB threshold modulation, which exhibits excellent linearity and stability in gas detection and high sensitivity in displacement detection. Moreover, the sensor performance can be flexibly tuned via structural parameters. We have also compared our proposed structure with other reported papers and found that our structure has better sensitivity and a simpler structure, as shown in Table 1.

Table 1. Comparison between different refractive index sensing schemes.

Ref.	Mechanism	Structure	Sensitivity	Frequency Range
[25]	Mode coupling sensor	Graphene-based Bragg reflector structure	1085°/RIU	THz
[26]	Bloch surface wave sensor	Prism-photonic crystal structure	117°/RIU	THz
[27]	SPR sensor	Grating structure	237°/RIU	Near Infrared
[28]	SPR sensor	Otto structure	34.11°/RIU	THz
[29]	Mode coupling sensor	Otto structure	3260 RIU ⁻¹	THz
This work	FP cavity nonlinearity	WSM- FP cavity	96.48 × 10 ⁵ V/m·RIU	THz

4. Conclusions

In this work, we theoretically investigate the low-threshold and tunable OB in a 1D PhCs-FP cavity structure incorporating WSM and its application potential as gas sensors and displacement sensors in the terahertz band. By combining the local field enhancement effect of the FP cavity and the large third-order nonlinear conductivity of WSM, we successfully achieve OB with a low threshold. By adjusting the Fermi energy of WSM, the incident angle, the period number of the 1D PhCs, the refractive index of the gas in the cavity, and the position of the WSM layer within the cavity, we can conveniently realize the flexible tuning of the upper and lower thresholds of OB and the width of the hysteresis

loop. Furthermore, the dual-parameter sensing scheme based on OB threshold modulation proposed in this study successfully achieves ultra-high linearity and excellent stability for gas refractive index sensing, as well as high sensitivity for displacement sensing.

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Data Availability Statement: The data that support the findings of this study are available from the corresponding author upon reasonable request.

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