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**NOVEL WAVE MANIPULATION OF DESIGNER SURFACE
PLASMONS**

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NOVEL WAVE MANIPULATION OF DESIGNER SURFACE PLASMONS

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For my parents.

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List of Publications

First or Co-first author:

1. F. Gao*, **Z. Gao*** (*contribute equally), X. Shi, Z. Yang, X. Lin, H. Xu, J. D. Joannopoulos, M. Soljacic, H. Chen, L. Lu, Y. Chong, and B. Zhang, Probing topological protection using a designer surface plasmon structure, **Nat. Communication**. **7**, 11619 (2016).
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3. **Z. Gao**, Z. Yang, F. Gao, H. Xue, Y. Yang, J. Dong, B. Zhang, Valley surface-wave photonic crystal and its bulk/edge transport, **Phys. Rev B** 96, 201402(R) (2017).
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5. **Z. Gao**, F. Gao, Y. Zhang, B. Zhang, Deep-subwavelength magnetic-coupling-dominant interaction among magnetic localized surface plasmons, **Phys. Rev B**. 93, 195410 (2016).
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Abstract

Due to its subwavelength field confinement and low propagation loss, designer surface plasmon (DSP) holds considerable promise in microwave- to infrared- frequencies device applications. This thesis studies the novel wave manipulation of DSP and its various applications in the microwave frequency range. We first explore the band gap nature of surface-wave band-gap crystal implemented on a single metal surface. By merging the forbidden bandgap of photonic crystal and the deep subwavelength nature of surface plasmons, the concept of surface-wave band-gap crystal with subwavelength-scale unit cell was developed with superior properties including high-efficiency pulse routing around sharp corners and multi-directional splitting of surface waves with full bandwidth isolation. Then we extend the concept of surface-wave band-gap crystal to the well-known Sievenpiper “mushroom” structure and demonstrate that this thin high-impedance surface is a new type of ultrathin surface-wave band-gap crystal with thickness much smaller than the operating wavelength. After the study of propagating DSP, we transfer to localized DSP and demonstrate the forward/backward switching of plasmonic waves propagation using sign-reversal coupling between localized DSP resonators. Since the properties of DSP can be fine-tuned by simply altering the underlying structural parameters and the mode profiles of designer surface plasmons can be imaged directly, this highly tunable platform allows us to explore the recently emerged topological photonics. In last part of this thesis, we extend the concept of topological photonic into DSP and demonstrate a flexible photonic topological insulators that can be bent, folded, or even twisted to mold the flow of topological edge states.

Chapter 1 introduces the basic concept of surface plasmon polaritons (SPPs) that exist on the metal/dielectric interface and localized surface plasmons (LSPs) that exist on a close metal surface at optical frequencies, and how to mimic SPPs and LSPs at low frequencies based on the concept of designer (or spoof) surface plasmons and designer localized surface plasmons.

Chapter 2 studies a class of surface-wave band-gap crystal implemented on a single metal surface which exhibits a complete band gap for surface waves. Surface waves can be tightly guided along the line defect and near-perfect transmission around sharp corners can be achieved over a broad frequency band. Moreover, functional components including a T-shaped splitter and a square open resonator are also demonstrated.

Chapter 3 studies a multi-directional plasmonic surface-wave splitter with full bandwidth isolation based on a new waveguiding mechanism of surface modes through weak couplings between tightly localized defect cavities in an otherwise gapped surface-wave band-gap crystal. The dispersion of coupled defect surface modes is tight-binding-like dispersion, rather than the polaritonic dispersion of conventional designer surface plasmons which starts from light line and approaches to the asymptotic frequency at the first Brillouin zone boundary. Therefore, the waveguiding bands of different branches of the multidirectional plasmonic splitter can be completely isolated which cannot be achieved using conventional designer surface plasmon waveguides.

Chapter 4 studies an ultrathin surface-wave band-gap crystal with thickness much smaller than the operating wavelength implemented on a well-known Sievenpiper “mushroom” structure. Functional devices such as a sharp bend and a T-shaped splitter are demonstrated with good performance.

Chapter 5 studies the forward/backward switching of plasmonic waves propagation using sign-reversal coupling. By directly measuring the tight-binding Bloch waves on a periodic array of coupled designer surface plasmon resonators in the microwave regime, we demonstrate multi-band forward/backward switching of plasmonic wave propagation. Our approach makes use of sign-reversal coupling that occurs when switching the coupling configuration between two photonic resonators. Direct experimental measurements of plasmon dispersion curves confirm the forward/backward switching of plasmonic waves propagation in the same frequency range.

Chapter 6 studies a mechanically flexible photonic topological insulator that supports robust topological photonic states on a curved surface. Spatial topologies achieved by folding the flexible photonic topological insulator serve as a new freedom to manipulate the propagation of topological photonic states. This work bridges the gap between the emerging field of topological photonics and the technologically promising field of flexible photonics.

Chapter 1

Introduction

1.1 Surface plasmon polaritons (SPPs) and localized surface plasmons (LSPs)

Surface plasmon polaritons (SPPs), generally described as electromagnetic (EM) surface waves on a metal/dielectric interface [1], show promising potentials to merge photonics and electronics on the same metal circuitry [2,3] because of their inherent subwavelength field confinement. As shown in Fig. 1.1, surface plasmons are essentially light waves that are trapped on the metal surface because of their interaction with the free electrons of the conductor. Due to the resonant interaction between light and electrons that constitutes the SPPs, it can be tightly confined within the subwavelength scale below the diffraction limit of light. Due to its excellent electromagnetic property, SPPs have given rise to the large amount of investigations on applications for surface enhanced Raman spectroscopy (SERS) [4], nonlinear light generation [5] and subwavelength optical devices enabling the miniaturization of optical components [6-9].

Figure 1.1. Schematic of propagating surface plasmon polaritons. Bottom yellow region represents metal, and its upper region is dielectric (air).

Beside the propagating SPPs on flat metal surfaces, localized surface plasmons (LSPs) [10-11] in structures with closed surfaces, i.e., plasmonic particles, also attract much attention due to its wide and potential applications in bio-sensing [12], optical nanoantennas [13-14], and subwavelength optical devices [15-17]. The electric field distribution of LSPs on a plasmonic particle is shown in Fig. 1.2, exhibiting a dipolar characteristic.

Figure 1.2. Schematic of localized surface plasmon polaritons. Gold spheres represent plasmonic particles.

1.2 Designer (or spoof) surface plasmon polaritons

Although metals are remarkably lossy at high frequencies such as in the visible spectrum and beyond, they can be treated as perfect electric conductors (PECs) with negligible loss in the far-infrared, terahertz, and microwave spectra and support surface electromagnetic waves, known as Sommerfeld or Zenneck waves [18-19]. However, the poor confinement of Sommerfeld or Zenneck waves makes them not suitable for routing applications. To solve this problem, some artificial textured metallic structures were proposed [20-21] to increase the confinement of surface electromagnetic waves at microwave frequencies, which can be traced back to the 1950s and 1960s. Recently, in 2004, Prof. Pendry developed the concept of designer (or spoof) surface plasmon polaritons from the viewpoint of metamaterials [22]. Using the effective media theory, he verified that the surface modes supported on a holey structured metal surface at low frequencies exhibit plasma forms, as shown in Fig. 1.3. Later in 2005, A. P. Hibbins experimentally verified the surface plasmon-like nature of the surface EM modes supported on a square brass tubes [23], as shown in Fig. 1.4.

Figure 1.3. (a) Schematic of the holey structured metal surface. (b) Dispersion relation of the designer surface plasmon polaritons.

Figure 1.4. (a) Photograph of the experimental sample showing the square brass. (b) Measured dispersion relation of the designer surface plasmon polaritons.

Due to its low loss and subwavelength confinement, designer surface plasmons show great potential in building high-efficiency subwavelength photonic circuits and other unprecedented functional devices from microwave to infrared spectra [24-26]. In particular, a surface waveguide consists of an easy-to-manufacture periodic chain of metallic box-shaped elements protruding out of a metallic surface was theoretically proposed [27] and experimentally verified later at microwave frequencies [28-30].

Figure 1.5. (a) Dispersion relation of domino plasmons for various lateral widths L . Inset show the schematics of domino plasmon waveguide.

As shown in Fig. 1.5, this waveguide structure was termed as “domino array”, and the supported surface modes as “domino plasmons (DPs)”. From the dispersion relations of domino plasmon waveguides with different lateral widths, we can see that domino plasmon modes are rather insensitive to the lateral waveguide width, preserving tight confinement even when the waveguide transverse dimensions are well in the deep subwavelength regime. In 2013, Shen et al. then pushed the idea to the extreme and made the lateral width of domino plasmon waveguide to massive sub-wavelength [31], forming an ultrathin and flexible film which can be bended, folded, even twisted to mold the flow of propagating designer surface plasmons. As shown in Fig. 1.6. They termed the surface modes supported by this flexible films as “conformal surface plasmons (CSP)”, and demonstrate the flexibility of this structures by coiling the waveguide into complex geometric.

Figure 1.6. (a) Dispersion relation of conformal surface plasmons for various copper thickness t . (b) Fabricated samples of conformal surface plasmon waveguides.

1.3 Designer localized surface plasmons

Besides the propagating designer surface plasmons (designer-SPPs) mimic real surface plasmons on a flat metal surface, the concept of designer localized surface plasmons (designer-LSPs) was also proposed to mimic the properties of localized surface plasmons (LSPs) in structures with closed surface, i.e., plasmonic particles [32]. As shown in Fig. 1.7, the two-dimensional corrugated closed metal surface can be modeled using effective medium theory as an inhomogeneous and anisotropic layers of thickness $(R-r)$ wrapped around a PEC cylinder of radius r . The scattering cross section (SCS) of the corrugated closed metal surface is very similar with that of a Drude metal cylinder with radius R , revealing the almost perfect analogy between designer-LSPs and conventional LSPs. In 2013, the designer-LSPs was first experimentally verified on an ultrathin plasmonic metamaterial [33], as shown in Fig. 1.8. The ultrathin plasmonic metamaterial supported dipole, quadrupole, hexapole, octopole, decapole, dodecapole, quattuordecapole modes, respectively. To further decrease the size of the designer

Figure 1.7. (a) Schematic of closed PEC surface textured with grooves. (b) Equivalent core-shell metamaterial model.

plasmonic resonator to deep subwavelength, a spiral metallic structure which supports both electric dipole mode and magnetic dipole mode is proposed and experimentally demonstrated [34], as shown in Fig. 1.9.

Figure 1.8. Experimental demonstration of designer-LSPs on an ultrathin plasmonic metamaterial and its resonance modes profiles.

Figure 1.9. (a) Experimental demonstration of an ultrathin textured metallic spiral disk with electric and magnetic dipole modes. (b) Mode profile of electric dipole mode. (c) Mode profile of magnetic dipole mode.

1.4 Motivation and Outlook

First, surface plasmon polaritons (SPPs), or localized electromagnetic surface waves propagating on a metal/dielectric interface, are deemed promising information carriers for future subwavelength terahertz and optical photonic circuitry. However, surface waves fundamentally suffer from scattering loss when encountering sharp corners in routing and interconnection of photonic signals. Previous approaches enabling scattering-free guidance of surface waves around sharp corners are limited to either a bulky waveguide environment or extremely narrow bandwidth, being unable to guide a surface-wave pulse (SPP wavepacket) on a single metal surface. Here, in a surface-wave band-gap crystal implemented on a single metal surface, we demonstrate in time domain routing a surface-wave pulse around multiple sharp corners without perceptible deterioration. Moreover, a new waveguiding mechanism for coupled surface-wave defect modes and a multi-directional plasmonic splitter with full bandwidth isolation are also proposed and experimentally demonstrated. Our Part I work not only offers a solution to surface-wave pulse routing, but also provides spatio-temporal information of surface-wave pulse propagation in presence of sharp corners, both of which are desirable in developing high-performance, large-scale integrated photonic circuits.

Second, backward waves with wave-front propagation opposite in direction to that of energy flow have attracted considerable interest in the context of photonic metamaterials. However, switching between forward and backward waves in the same frequency range has remained a challenge. Here, by directly measuring the tight-binding Bloch waves on a periodic array of coupled designer surface plasmon resonators in the microwave regime, we demonstrate

multi-band forward/backward switching of plasmonic wave propagation. Our approach makes use of sign-reversal coupling that occurs when switching the coupling configuration between two photonic resonators. Direct experimental measurements of plasmon dispersion curves confirm the forward/backward propagation of plasmonic waves in the same frequency range. Our work provides a solution to forward/backward switching of plasmonic wave propagation, and may find potential applications in photonic integrated systems.

Third, photonic topological insulators have received extensive attentions recently due to its unique topological-protected backscattering-immune propagation against defects. To date, however, all the proposed photonic topological insulators have a major limitation associated with their inherent 3D bulky and rigid structures. A flexible and ultrathin photonic topological insulator that can be bent, folded, or even twisted to manipulate the topologically protected edge states in non-planar 3D physical spaces still remain unrealized. In chapter 6 we propose and experimentally demonstrate a flexible photonic topological insulator supporting robust topological edge states on curved surfaces. We show that the topological edge states can exist on curved surfaces and exhibit robustness against intrinsic and exotic disorders. We implement and investigate two exotic global topologies to manipulate the propagation of topological edge states.

Part I

Surface-wave band-gap crystal

Chapter 2

Surface-wave pulse routing around sharp corners

2.1 Surface-wave scattering loss at the sharp corners

Surface plasmon polaritons (SPPs), generally described as electromagnetic (EM) surface waves on a metal/dielectric interface [1], show a promising potential to merge photonics and electronics on the same metal circuitry [2,3] because of their inherent subwavelength confinement. Although metals are remarkably lossy in the visible spectrum and beyond, they can be treated as perfect electric conductors (PECs) with negligible loss in the far-infrared, terahertz, and microwave spectra. Designer SPPs are EM surface modes at these frequency bands, supported on a periodically textured metal surface [22,23]. They show great potential in building high-efficiency subwavelength photonic circuits and other unprecedented functional devices from microwave to infrared spectra [24-31]. However, while being routine in electronic circuits, near-perfect transmission through sharp corners is fundamentally difficult for SPP waveguides in low frequency range (Terahertz or Gigahertz) because of the intrinsic scattering loss of guiding EM surface waves around sharp corners. The scattered energy will not only deteriorate transmission, but also introduce unwanted crosstalk of signals [42], being especially detrimental in compactly integrated photonic circuits. In this Chapter, we will overcome this challenge and realize near-perfect transmission through sharp corners in microwave frequencies.

To date, scattering-free guidance of EM surface waves around sharp corners has been realized only in photonic topological insulators [43-47], which support topological EM surface states as an analogue of electronic chiral edge states in topological insulators, and later in surface-wave transformation cloaks [48], which apply transformation-optics to effectively warp the EM space around sharp corners analogously to the way in which gravity curves space in general relativity. However, these approaches generally require a bulky waveguide environment, usually between two parallel metal plates, or only apply to a very narrow bandwidth, being unable to manipulate surface-wave pulses (SPP wave packets) on a single metal surface.

On the other hand, the capability of manipulating surface-wave pulses is extremely desirable in ultrafast optical data processing and recording. It is therefore necessary to have a clearer and more complete understanding in the spatio-temporal interplay between surface-wave pulses and the underlying photonic structures carrying them [49-52]. The surface-wave pulse propagation at optical frequencies has been directly imaged along a straight path [53-57] using the femtosecond laser pulse tracking technique. However, how a surface-wave pulse interacts with a sharp corner remains unobserved in reality, not to mention the imaging of the process of routing a surface-wave pulse around sharp corners with near-perfect transmission.

It is the purpose of this article to develop an approach of routing surface-wave pulses around multiple sharp corners, especially with the spatio-temporal information in the routing process. We first show that by introducing a line defect in a surface-wave band-gap crystal, the surface waves can be tightly guided along the line defect in the surface-wave band gap. (Recent preprints [58-59] have shown the feasibility of guiding and slowing surface waves in a surface-

wave band-gap crystal.) Then by tracking and mapping a surface-wave pulse as it propagates around multiple sharp corners in a surface-wave band-gap crystal, as schematically shown in

Figure 2.1. Schematic of a surface-wave pulse routing through multiple sharp corners along an arbitrary path in a surface-wave band-gap crystal.

Fig. 2.1, we schematically demonstrate that near-perfect transmission around multiple sharp corners can be achieved for surface-wave pulses.

The remarkable high-efficiency time-domain surface-wave pulse routing around multiple sharp corners in the surface-wave photonic crystals is mainly because of the following two reasons. First, the wide forbidden band gap of the surface-wave photonic crystal suppresses scattering of surface waves into the ambient space in a broad bandwidth. This feature is similar to previous wavelength-scale photonic circuits in a photonic crystal [60-62]. Second, because the eigenmode supported by the line-defect waveguide of the surface-wave photonic crystal is essentially coupled toroidal resonance mode with four-fold rotational symmetry [63-65], the momentum of the surface EM waves before and after passing the sharp corners in an extremely compact space is inherently matched [66]. Thus there is no need to purposely construct extra local resonances at the sharp corners in order to facilitate tunneling the EM energy through, being in sharp contrast to previous photonic crystal-based proposals [60-62]. Similar high-

efficiency transmission through sharp corners can also be achieved by previous wave manipulation along a line defect in a 3D bulk wire medium metamaterial [67,68]. However, the waveguiding frequency of the line defect in the 3D wire medium metamaterial is almost fixed by the length of a local resonant defect wire, resulting in an extremely narrow waveguiding bandwidth [67] and cannot support a real broadband pulse signal. On the contrary, the (spoof) surface plasmons supported by the surface-wave photonic crystal are broadband ensemble resonances, which can be utilized to bend surface waves in a relatively broad bandwidth, more importantly, to bend a real broadband time-domain SPP wave packet on a single metal surface.

2.2 Experimental demonstration of the surface-wave band gap

We start the experimental demonstration with a surface-wave band-gap crystal that consists of a periodic (period $d = 5.0$ mm) 25×25 square lattice of circular metal rods with height $H = 5.0$ mm and radius $r = 1.25$ mm, all standing on a flat metallic surface, as shown in Fig. 2.2a. The sample was fabricated with mechanical machining with tolerance of 0.1 mm and the whole size of 125×125 mm². We first fabricated an aluminum plate with square array of holes and the circular metal rods, then we inserted the metal rods in to the holes to assemble them together. Thus the metal rods can be replaced as will to change the rod height. The commercial software CST Microwave Studio is used for numerical simulation of the band structure of the surface-wave band-gap crystal. In band structure simulation, a single metallic rod of the surface-wave band-gap crystal is adopted as the unit cell. The periodic boundary conditions are imposed with

phase shifts in both x and y directions. By scanning the phase shifts in the first Brillouin zone of a square lattice as shown in the inset of Fig. 2.2b, the complete band structure of the perfect surface-wave photonic crystal can be obtained. The calculated band structure, as shown in Fig. 2.2b, reveals a complete forbidden band gap (purple region) for surface waves from 12.6 GHz to 27 GHz for the current surface-wave band-gap crystal.

Figure 2.2. Experimental demonstration of a surface-wave band-gap crystal. (a) Photograph of the surface-wave band-gap crystal. A square array of aluminum rods (with height $H = 5$ mm and periodicity $d = 5$ mm) stand on a metallic surface. (b) Calculated band structure of the surface-wave band-gap crystal. Purple region represents the forbidden band gap. (c) Measured transmission spectrum through the surface-wave band-gap crystal.

A vector network analyser (VNA; Agilent Technologies N5225A) and two small near-field homemade antennas were adopted in measurement. The near-field antennas, serving as the point source and probe, were constructed by cutting a SMA coaxial cable in such a way that a short (5 mm) piece of its core shows up outside of the shielding. Both the point source and probe were oriented along the vertical direction and thus measured only the E_z component of the electric field. In the VNA setup, 1001 points were sampled in the frequency spectrum and the power to 10 dB was set in order to gain a sufficient signal to noise ratio.

At two opposite sides of the perfect surface-wave band-gap crystal (without shortening any metallic rod), we placed the point source and probe to measure the S-parameter (S_{21}) as the relative transmission between the point source and probe. The absolute value of the S_{21} parameter has no practical meaning because of the large impedance mismatch between the point source/probe and the crystal. The measured transmission spectrum supported by the surface-wave band-gap crystal is shown in Fig. 2.2b. The measured transmission exhibits a substantial drop in the frequency range from 12.6 GHz to 27 GHz, corresponding to the forbidden band gap of the surface-wave band-gap crystal. When the sample was removed, the signal level reduced below -40 dB, which means negligible energy is received by the probe. Hereafter, all data of transmission and near-field measurements will be normalized.

2.3 Spectral measurement and near-field imaging of the waveguiding along a straight line-defect waveguide

A straight line defect can be constructed by shortening the height of a row of rods in the middle from $H = 5.0$ mm to $h = 4.3$ mm, as indicated in Fig. 2.3a. In the experiment, we remove one row of metal rods with height $H = 5.0$ mm and insert metal rods with height $h = 4.3$ mm to construct the straight line-defect waveguide. Figure 2.3b shows the dispersion (red line) of the line-defect waveguide that is calculated with the eigenmode solver of the commercial software CST Microwave Studio. In the simulation of the dispersion of line-defect waveguide, a supercell that consists of seven metallic rods in y direction and one metallic rod in x direction is adopted. The center rod in the super cell is shortened from height $H = 5$ mm to $h = 4.3$ mm while the other rods keep unchanged. Periodic boundary condition has been applied in the x direction, and the electric boundary conditions have been used in both the y and z directions. By scanning the phase shift between the periodic boundaries from 0 to π (first Brillouin zone), we can get one eigenfrequency for the line-defect eigenmode with 0.05π step for each phase shift; thus, the dispersion of the line-defect waveguide in the first Brillouin zone can be obtained. Note that only the part below the light line and within the band gap (purple region) can confine and guide surface waves along the line-defect waveguide.

In fact, even without the surrounding surface-wave band-gap crystal, this row of shortened rods themselves also can guide surface waves. This is a typical traditional surface-wave waveguide (widely referred to as “domino plasmons” [27]) that has been studied extensively with various device applications [28-31]. For comparison, we also simulated the dispersion (blue line in Fig. 2.3b) of surface waves guided along this domino-plasmon waveguide, when all the other rods of the surrounding surface-wave band-gap crystal were removed. In the simulation of the dispersion of domino-plasmon waveguide, a single metallic rods with height

$h = 4.3$ mm is adopted. Periodic boundary condition has been applied in the x direction, and the electric boundary conditions have been used in both the y and z directions. By scanning the phase shift between the periodic boundaries from 0 to π (first Brillouin zone in x direction), we can get one eigenfrequency for the domino-plasmon eigenmode with 0.05π step for each phase shift; thus, the dispersion of the domino-plasmon waveguide in the first Brillouin zone can be obtained. It can be seen that this domino-plasmon waveguide exhibits a polaritonic dispersion (blue line in Fig. 2.3b) that starts from the light line and then tends to zero group velocity at the Brillouin zone edge. Due to the much tighter confinement provided by the surrounding surface-wave photonic crystal, the tight-binding-like dispersion of the line-defect waveguide (red line in Fig. 2.2b) is rather different with the polaritonic dispersion of the domino-plasmon waveguide (blue line in Fig. 2.2b).

To test the line-defect waveguide, we used a vector network analyzer (Agilent Technologies N5225A) to measure the transmission in the frequency range from 10 to 15 GHz. Two monopole antennas, serving as the source and the probe, respectively, were placed next to the input and the output of the line-defect waveguide, as illustrated in the inset of Fig. 2.3c. We first measured the transmission spectrum (grey line in Fig. 2.3c) of the perfect surface-wave band-gap crystal without shortening any rod. The band gap starting from 12.6 GHz can be clearly seen (the upper edge of the band gap at 27 GHz is not shown here). We then measured the transmission spectrum of the straight line-defect waveguide (red line in Fig. 2.3c). Evidently, the line-defect waveguide transmits surface waves over a broad frequency range from 12.8 to 14.1 GHz within the band gap of the surface-wave band-gap crystal, being consistent with the dispersion in Fig. 2.3b. To map the near-field distributions of the surface

Figure 2.3. Spectral measurement and near-field imaging of the waveguiding along a line-defect waveguide in a surface-wave band-gap crystal. (a) Photograph of the straight line-defect waveguide formed by shortening a row of metallic rods from $H = 5$ mm to $h = 4.3$ mm in the middle of the surface-wave band-gap crystal. Inset shows the schematic representation of the

pulse propagation. (b) Dispersion of the line-defect waveguide compared to that of the traditional domino-plasmon waveguide. Purple region represents the forbidden band gap of the surface-wave band-gap crystal. (c) Measured transmission through the straight line-defect waveguide and through the perfect surface-wave band-gap crystal. Data are normalized to the maximum field amplitude on top of the third rod from the input. Purple region represents the band gap of the surrounding surface-wave band-gap crystal. (d-f) Measured field profiles in a transverse plane 1 mm above the line-defect waveguide at 13.1, 13.4, and 13.7 GHz, respectively.

waves, the source was a 5 mm-long monopole antenna that was placed near the surface-wave band-gap crystal. The probe, another 5 mm-long monopole antenna, was able to scan in the horizontal plane 1 mm above the sample, measuring both the amplitude and phase of the resonance wave fields. By stepping the probe with 0.5 mm increment in both x and y directions, a full 2D spatial field map can be obtained experimentally. Both the source and probe are connected to a vector network analyzer (Agilent Technologies N5225A). Figure 2.3d-f show the measured spatial field profiles in a transverse plane 1 mm above the line-defect waveguide at 13.1, 13.4 and 13.7 GHz, which was mapped with a microwave near-field scanning probe. It can be seen that tightly confined surface waves are guided along the line-defect waveguide in a broad frequency range.

2.4 Time-domain surface-wave pulse propagation along a straight line-defect waveguide

Since tightly confined surface waves can be supported by the line-defect waveguide in a broad frequency range, we used a surface-wave pulse to study the dynamic properties of the line-defect waveguide in such a surface-wave band-gap crystal. In the experiment, we used an arbitrary waveform generator (Tektronix AWG70001A) to generate an 800 MHz wide Gaussian pulse centered around 13.4 GHz within the line-defect waveguiding band and launch it into the line-defect waveguide through an input monopole antenna. Using the other output monopole antenna connected to an oscilloscope (Tektronix MSO 72004C), we measured the pulse signal at the third rod from the input as the input pulse (grey line in Fig. 2.4a; delay-time 0 ns) to minimize the interference from the source. Similarly, the output pulse (red line in Fig. 2.4a) was measured at the third rod from the output. To avoid reflections from the exit-end and source-end, we put absorption materials around the samples to eliminate the reflection from the source-end and exit end. As shown in Fig. 2.4a, it takes 6.2 ns for the pulse to propagate through the 100 mm-long straight line-defect waveguide.

Figure 2.4b shows the spectral signature of measured pulse signal. The spectral signatures of the pulse signals are obtained by Fourier transforming the measured Gaussian pulse. For comparison, the ideal pulse profile as the input to the arbitrary waveform generator is also shown. As the surface-wave pulse propagates along the straight line-defect waveguide, it maintains its amplitude and spectral signature without perceptible distortion.

In the following, we reconstruct the dynamic images of pulse propagation based on the microwave near-field imaging measurement. Two steps are adopted in the experiment. First, we mapped the surface-wave near-field distributions (electric field component E_z) on a

transvers plane 1 mm above the line-defect waveguide, covering the frequency range from 13 GHz to 13.8 GHz with a step of 0.01 GHz (80 different frequencies) within the waveguiding band. The spectral measurement and near-field imaging results are shown in Fig. 2.3d-f.

Figure 2.4. Surface-wave pulse propagation along a straight line-defect waveguide. (a) Measured surface-wave pulse signals at the input and output of the waveguide. (b) Spectral signatures of the measured input pulse and output pulse compared with the ideal case. (c) Reconstructed surface-wave pulse propagation in the straight line-defect waveguide at different time frames.

Second, we did an inverse Fourier transform on measured near-field profiles to reconstruct the Gaussian pulse (band width 0.8 GHz; central frequency w_0 : 13.4 GHz) $\exp(iw_0t) \cdot \exp(-((t-t_0)/\sigma)^2)$, where w_0 is the central frequency, $t_0 = 5$ ns is the central time, and $\sigma = 1.5$ is the pulse width.

Figure 2.4c shows the reconstructed images of surface-wave pulse propagation through the line-defect waveguide. The pulse signal reaches the last rod at 6.4 ns, being consistent with the measured time-delay of 6.2 ns in Fig. 2.4a.

2.5 Spectral measurement and near-field imaging of waveguiding through a U-shaped line-defect waveguide

We then examine the performance of a U-shaped line-defect waveguide with two right-angle sharp corners, as shown in Fig. 2.5a. The source and the probe were placed next to the input and output in measuring the transmission. To facilitate comparison, all transmission data hereafter are normalized to the same value as the maximum field amplitude on top of the third rod from the input (this finite distance from the input is to minimize the interference from the source) along the straight line-defect waveguide as demonstrated in Fig. 2.3a. It can be seen that the transmission through two sharp corners is almost unity in the whole waveguiding band from 12.8 to 14.1 GHz. We present in Fig. 2.5c-e the measured field profile in a transverse plane 1 mm above the U-shaped line-defect waveguide at 13.1, 13.4, and 13.7 GHz, respectively. Surface waves propagate smoothly around two sharp corners without perceptible decay. For comparison, we also measured the transmission (blue line in Fig. 2.5b) through the domino-plasmon U-shaped waveguide after all the other rods in the surrounding surface-wave

Figure 2.5. Spectral measurements and near-field imaging of waveguiding around two sharp corners. (a) Photograph of the U-shaped line-defect waveguide. Inset shows the schematic of the pulse propagation. (b) Measured transmission spectrum through the U-shaped line-defect waveguide with two sharp corners compared to that through the U-shaped domino plasmon waveguide with two sharp corners. Inset shows the schematic of the U-shaped line-defect waveguide. (c-e) Measured field profiles of the U-shaped line-defect waveguide at 13.1, 13.4, and 13.7 GHz. (f-h) Measured field profiles of the U-shaped domino-plasmon waveguide at 13.1, 13.4, and 13.7 GHz.

band-gap crystal were removed. The transmission is dramatically reduced by about 8 dB compared to the line-defect waveguide in presence of the surface-wave band-gap crystal. The peak of transmission around 13.7 GHz is close to the asymptotic frequency of domino plasmons at which the field confinement reaches maximum. The transmission decreases when the frequency gets lower because of the weaker field confinement. Figure 2.5f-h show the measured field profiles of the domino-plasmon U-shaped waveguide at 13.1, 13.4, and 13.7 GHz, in which the scattering at the sharp corners is evident. Note that though the domino-plasmon waveguide generally induces serious radiation loss at the sharp corner due to the dramatic momentum mismatch before and after passing the sharp corner, its waveguide structure is much simpler and the fabrication is easier than the line-defect waveguide with surrounding surface-wave photonic crystal. Moreover, the compactness of the integrated plasmonic circuits with domino-plasmon waveguide will be higher than that of the line-defect waveguide.

2.6 Time-domain surface-wave pulse propagation along a U-shaped line defect waveguide

We then examine the time-domain performance of the U-shaped line-defect waveguide with two right-angle sharp corners. Using the same experimental method as in Fig. 2.4, we measured the input pulse and the output pulse, as shown in Fig. 2.6a. The measured two pulse signals are almost the same after the pulse has propagated through two sharp corners. We reconstruct the images of dynamic pulse propagation from the measured near-field distributions (The spectral

measurement and near-field imaging results are shown in Fig. 2.5). The reconstructed images are shown in Fig. 2.6b. It can be seen that the pulse reaches the first and second sharp corners at 4 ns and 6 ns, respectively. The pulse is guided around both corners, leaving the U-shaped line-defect waveguide at 8 ns with imperceptible scattering.

Figure 2.6. Surface-wave pulse propagation around a U-shaped line-defect waveguide with two sharp corners. (a) Measured surface-wave pulse at the input and output of the U-shaped line-defect waveguide. (b) Reconstructed surface-wave pulse propagation in the U-shaped line-

defect waveguide at different time frames. (c) Reconstructed surface-wave pulse propagation in the U-shaped domino-plasmon waveguide at different time frames.

For comparison, we also reconstruct the surface-wave pulse propagation along a U-shaped domino-plasmon waveguide, when all the other rods in the surrounding surface-wave band-gap crystal were removed. As presented in Fig. 2.6c, the pulse reaches the first sharp corner at 1.4 ns and the second sharp corner at 2 ns. Both corners cause strong scattering and reflection, leading to almost zero transmission at the output.

2.7 Mechanism of surface-wave pulse guidance around sharp corners

To understand the underlying mechanism of the high-efficiency surface-wave pulse guidance around a sharp corner, we first simulate the magnetic and electric field distributions of an isolated point defect that can be realized by shortening only one single rod (black circle) in the surface-wave band-gap crystal, as shown in Fig. 2.7a and Fig. 2.7d, in which a transverse plane at the half height of the 5-mm-high metallic rods is adopted. The closed loop of magnetic field and the quadrupolar profile of electric field reveal that the eigenmode of a point defect in the surface-wave band-gap crystal is a localized magnetic toroidal mode and an electric quadrupole mode [63-64] with four-fold rotational symmetry. In the numerical simulation of the electric and magnetic field distributions, we using the transient solver of CST Microwave Studio. All boundaries are set as open boundaries (perfect matched layers), and the metals are modeled as perfect electric conductors. We use a dipole current (discrete port) as the point source to excite

the waveguide structure and using another discrete port as the probe and monitor the electric and magnetic field distributions. When the energy decayed to -60 dB, the simulation will be stopped and the electric and magnetic field distributions at different frequencies are obtained to analyze the mechanism of surface-wave pulse guidance around sharp corners.

Then we simulate the magnetic and electric field distributions of a line-defect waveguide with a sharp corner, as shown in Fig. 2.7b and Fig. 2.7e. We choose the frequency of 13.09 GHz for illustration in order to make the phase difference between fields of neighboring shortened rods exactly equal to $\pi/2$, such that when the magnetic and electric fields are maximum around one shortened rod, they are zero around the neighbouring shortened rod. Thus the overlap of fields between neighboring shortened rods will not obscure the mode profile of each shortened rod to be observed. It can be seen that, even in an array of closely positioned shortened rods, the magnetic toroidal mode and the electric quadrupole mode of each shortened rod still maintain their mode profiles with four-fold rotational symmetry. This is in sharp contrast to typical plasmonic waveguides formed by an array of plasmonic resonators [46]. For typical coupled plasmonic resonators, only when their coupling strength is weak can each resonator maintain its mode profile. Yet in a surface-wave band-gap crystal, the four-fold rotational symmetry of mode profile around a shortened rod is maintained by the crystal symmetry, rather than the coupling strength between neighboring shortened rods.

For comparison, we also plot the magnetic and electric field distributions of a domino-plasmon sharp bend, when all other rods in the surface-wave band-gap crystal are removed, as shown in Fig. 2.7c and Fig. 2.7f. The frequency of 12.00 GHz is used to set the phase difference between neighboring rods equal to $\pi/2$. It can be observed that, without the surrounding

Figure 2.7. Field distributions of a point defect, a line-defect sharp bend in the surface-wave band-gap crystal and a domino-plasmon sharp bend without the surrounding surface-wave

band-gap crystal. Magnetic (a-c) and electric (d-f) field distribution of (a,d) an isolated point defect at 13.20 GHz, (b,e) a line-defect sharp bend at 13.09 GHz and (c,f) a domino-plasmon sharp bend at 12.00 GHz in a transverse xy plane at half height of the 5-mm-high metallic rods. Black circles represent the line defect with shortened rods. Grey circles represent the surrounding surface-wave band-gap crystal.

surface-wave band-gap crystal, both the magnetic and the electric fields do not exhibit four-fold rotational symmetry.

The four-fold rotational symmetry of mode profile is the key to realize near-perfect guidance around a sharp corner. When propagating through the sharp bend, because of the presence of the preserved four-fold rotational symmetry, the surface waves feel as if it was still propagating along a straight waveguide. This is very similar to the previous coupled-resonator optical waveguides (CROW) [66] in which the rotational symmetry of each resonator causes the perfect transmission through a sharp bend, given that the mode profile of each resonator should not be perturbed by near-field coupling. On the other hand, without the four-fold rotational symmetry, serious scatterings occurs at the sharp corner in a typical domino-plasmon waveguide (Figs. 2.7c and 2.7f).

2.8 Near-field imaging of waveguiding through a tortuous line-defect waveguide with multiple sharp corners

Because the height of rods in the surface-wave band-gap crystal can be adjusted conveniently,

it is feasible to form an arbitrary path for routing surface-wave pulses. In the following, we constructed a tortuous line-defect waveguide with multiple sharp corners (Fig. 2.8a) and examine its performance. To reconstruct the surface pulse evolution in a tortuous line-defect waveguide with multiple sharp corners, we first mapped the near-field distributions of the vertical component of electric field (E_z) in a transverse plane 1 mm above the tortuous waveguide from frequencies 13 to 13.8 GHz with step 0.01 GHz. As an example, we show the measured near-field profiles of the tortuous line-defect waveguide at 13.1, 13.4 and 13.7 GHz in Fig. 2.8b-d. It can be seen that surface waves can be high efficiently transmitted through six sharp corners without perceptible reflections and radiations.

Figure. 2.8. Near-field imaging of waveguiding around a tortuous line-defect waveguide with multiple sharp corners. (a) Photograph of the tortuous line-defect waveguide. Inset shows the schematic of the pulse propagation. (b-d) Measured field profiles of the tortuous line-defect

waveguide at 13.1, 13.4 and 13.7 GHz in a transverse plane 1 mm above the waveguide.

2.9 Time-domain surface-wave pulse propagation around multiple sharp corners

Figure 2.9. Surface-wave pulse propagation around a tortuous line-defect waveguide. (a) Measured surface-wave pulse signals at the input and output of the tortuous waveguide. (b)

Reconstructed surface-wave pulse propagation in the tortuous waveguide at different time frames.

Then we launched the same surface-wave pulse (central frequency 13.4 GHz and band width 0.8 GHz) to examine its performance. We measured the input pulse (grey line) and the output pulse (red line), as shown in Fig. 2.9a. A time delay of 10.5 ns and only 15% amplitude decay can be observed after propagating through six sharp corners in the surface-wave band-gap crystal. To gain deeper insight of the dynamic evolution of the pulse in the tortuous waveguide, we reconstruct the images of the pulse at different time frames, as shown in Fig. 2.9b. It is clear that the pulse signal can smoothly propagate around all six sharp corners without perceptible reflection or scattering, confirming the high transmission of surface waves through multiple sharp corners along an arbitrary path in the surface-wave band-gap crystal.

2.10 Surface-wave band-gap crystal waveguide components including T-shaped splitter and square open ring resonator

In the following, we construct a T-shaped splitter by connecting two straight line-defect waveguides in the surface-wave photonic crystal (inset of Fig. 2.10a). An ideal splitter should transmit 50% of the input power (or -1.5 dB transmission of field amplitude) into each output. We can observe in the measured transmission spectra from the right (black line) and left (red line) outputs in Fig. 2.10a that the input surface waves are split almost ideally by the T-shaped splitter. In contrast, when the surrounding rods are all removed, the transmission through the

domino-plasmon T-shaped splitter is reduced by about 6 dB. Figures 2.10b-c show the measured field profiles at 13.7 GHz in a transverse plane 1 mm above the line-defect T-shaped splitter and above the domino-plasmon T-shaped splitter, respectively. Scattering at the junction is evident in domino-plasmon T-shaped splitter, but negligible in the line-defect T-shaped splitter. From the comparison between the performances of the line-defect T-shaped splitter and the domino-plasmon T-shaped splitter, we can see that the surrounding surface-wave photonic crystal is crucial to suppress the scattering at the junction of T-shaped splitter and the intrinsic momentum match before and after passing the T junction.

We then proceed to bend surface waves successively across four sharp corners to form a square open resonator. Compared to circular ring resonators, a square resonator is apparently more suitable for ultra-compact photonic circuits because of its compactness. Its longer coupling region with a bus waveguide can also lead to higher coupling efficiency which is desirable in many photonic applications. However, a square open resonator for surface waves is difficult to realize because of the serious scattering loss at the sharp corners.

Figure 2.10. T-shaped splitter of surface waves. (a) Measured transmission spectra from the left and right outputs of the line-defect splitter compared to those from the left and right outputs of the domino plasmon splitter. Inset shows the schematic of the T-shaped splitter. Black dashed line represents -1.5 dB transmission. (b-c) Measured field profiles at 13.7 GHz in a transverse plane 1 mm above the line-defect splitter (b) and the domino plasmon splitter (c).

By connecting four finite segments of line-defect waveguides with the same length, a square open resonator can be formed, as illustrated in the inset of Fig. 2.11a. The source and the probe are indicated with two black arrows. Figure 2.11a presents the measured near-field response spectrum of this square open resonator, in which several resonance peaks ranked from the first-order mode (M_1) to the tenth-order mode (M_{10}) can be clearly observed. The measured field profiles corresponding to these resonance modes are shown in Fig. 2.11b. It can be seen that besides the first two modes (between 12.6 GHz and 12.8 GHz) that are close to the lower edge of the band gap, there is no perceivable scattering of the other modes. In this sense, these resonance modes are analogous to the recently proposed localized spoof surface plasmons [32-34], whose resonance modes are formed by bending surface waves around a circular textured metal disk. Note that square-shaped localized spoof surface plasmons have not been realized because of the scattering loss at sharp corners. The quality factors of resonance modes in Fig. 2.11a can easily rise beyond 100 (e.g. 123 for M_5 and 283 for M_8), while the previously realized localized spoof surface plasmons have quality factors generally less than 100. For completeness, we also measured the near-field response spectrum (blue line in Fig. 2.11a) of the domino-plasmon square resonator by removing the surrounding surface-wave photonic crystal; no obvious resonance can be observed.

Figure 2.11. Square open resonator of surface waves. (a) Measured near-field response spectrum of the line-defect square resonator compared to that of the domino plasmon square resonator. Inset shows the schematic of the resonator. Two monopole antennas as the source and probe are indicated with two black arrows. (b) Measured field profiles of the resonance modes.

2.11 Conclusion

In summary, by implementing a surface-wave band-gap crystal on a single metal surface, we

provide a solution to route surface-wave pulses around sharp corners, as a step forward towards solving a long-lasting bottleneck in highly integrated subwavelength plasmonic circuits. We applied a time-resolved near-field imaging technique to reveal the dynamics of a surface-wave pulse in both space and time as the pulse propagates through multiple sharp corners along an arbitrary path without perceptible reflections and radiations. Functional devices such as a high-efficiency T-shaped splitter and a square surface-wave open resonator have also been realized, both of which cannot be achieved with traditional surface-wave waveguides. The concept and method are general and can be directly applied to the terahertz and far-infrared frequency ranges to realize ultra-compact terahertz and photonic circuits with high-performance and high integration density.

Chapter 3

Coupled-defect surface waveguide and multi-directional plasmonic surface-wave splitters with full bandwidth isolation

3.1 Point defect cavity in a surface-wave band-gap crystal

By merging the forbidden band gap of photonic crystal and the deep subwavelength nature of designer surface plasmons, in Chapter 2 we have demonstrated that concept of surface-wave photonic crystal with wide surface-wave band gap for surface guided mode and high-efficiency pulse routing around line-defect sharp corners in the surface-wave photonic crystal. Based on this platform, we will further explore various novel manipulations of designer surface plasmons such as coupled-defect surface waveguide and multi-directional plasmonic surface-wave splitters with full bandwidth isolation.

Photonic crystals are artificial periodic dielectric or metallic structures that can prohibit the propagation of electromagnetic (EM) waves in a certain frequency range [60]. When the translational symmetry of photonic crystal is locally broken, tightly localized point defect modes within the photonic bandgap can be created [69]. Therefore, only evanescent waves can penetrate into the photonic crystal from this defect location within a finite distance, offering the possibility for photons to hop between neighboring defect cavities. Based on this mechanism, waveguiding through coupled defect modes [66] is theoretically proposed and experimentally demonstrated [70-73] in the context of photonic crystals. However, since

photonic crystals originate from Bragg interference, the field confinement of these defect modes is generally diffraction-limited, i.e. being limited to the scale of about half a wavelength.

On the other hand, surface EM waves supported on structured metal surfaces, termed as designer (or spoof) surface plasmons [22-23], possess spatial scales typically much smaller than the operating wavelength. In particular, some structured metal surfaces possess a complete forbidden band gap for surface guided modes [58-59]. Recently, it is proposed theoretically [74-75] and demonstrated experimentally [64-65] that a tightly localized surface defect cavity mode exists within the forbidden band gap when an engineered defect is introduced on a perfectly structured metal surface.

In this chapter we propose and demonstrate a type of waveguiding mechanism of coupled defect modes by merging the subwavelength nature of designer surface plasmons and the forbidden band gap of photonic crystal. This new waveguiding mechanism offers an alternative way to manipulate surface waves at subwavelength scales on textured metal surfaces. Compared with conventional designer surface plasmon waveguides [22-31] which usually suffer from serious crosstalk when packed densely, this coupled-defect surface waveguide (CDSW) is achieved through weak coupling between otherwise tightly localized surface defect cavities, which allows molding the flow of surface waves at will with minimal crosstalk.

Our purpose in this Chapter is to propose and demonstrate a waveguiding mechanism to manipulate the propagation of surface waves through coupled defect cavities. What is more, to realize a multidirectional plasmonic surface-wave splitter with full bandwidth isolation based on this waveguiding mechanism. We start our demonstration with an ideal surface-wave band-gap crystal which is the same with that of previous Chapter, as shown in Fig. 2.2a. First, we

construct a point defect cavity as the basic component. For this purpose, we partially shorten one single pillar from $h = 5$ mm to $h_d = 4.15$ mm at the center of the surface-wave band-gap crystal while keeping the surrounding pillars unchanged, as shown in the inset of Fig. 3.1a. Since the resonance frequency of the point defect cavity depends on the height of the shortened pillar, therefore, it falls exactly within the forbidden band gap of surrounding surface-wave band-gap crystal, and a defect cavity state is formed around the shortened pillar. The formation of the point defect cavity state in the surface-wave photonic crystal is mainly because of the following two reasons. First, the surface modes supported by the metallic rods is below the air light line, thus the resonance modes are surface states with evanescent decay in the vertical direction and tightly localized on the structured metal surface. Second, the resonance frequency of the surface state is located exactly within the forbidden band gap of surrounding surface-wave photonic crystals, thus the defect cavity state is highly confined in the transverse plane. To experimentally demonstrate this surface defect cavity, we adopt a microwave near-field imaging system to map its near-field profiles. The experimental setup is shown in the inset of Fig. 3.1a, which consists of a vector network analyzer (R&S ZVL-13) and a pair of monopole antennas to excite and probe the surface cavity modes. We first measure the near-field transmission of the perfect surface-wave band-gap crystal without any defect, as indicated with black line in Fig. 3.1a. We can observe no resonance peak in the transmission spectrum. Note that the origin of the random signals that exist at the 9-12.6 GHz frequency band come from the bulk states of the first band in the surface-wave photonic crystal, as shown in the calculated band structure of the surface-wave photonic crystal (red dotted line in Fig. 2.2b). The whole surface-wave photonic crystal structure is excited below the lower photonic band gap edge

Figure 3.1. (a) Measured transmission spectra of a surface-wave band-gap crystal with an isolated point defect cavity (red line) and without any point defect cavity (black line). Inset shows the experimental set up. (b) Measured field profile (E_z) of the point defect cavity at its resonance frequency 13.5 GHz. Inset shows the schematics of the point defect cavity. (c) Simulated field profile (E_z) of the point defect cavity at its resonance frequency 13.5 GHz in the transverse xy plane. (d) Simulated profile (E_z) of the point defect cavity in the vertical xz plane (black dashed line in Fig. 2(c)) at its resonance frequency 13.5 GHz.

(12.6 GHz) and all surface states are forbidden within the photonic band gap ranges from 12.6 GHz to 27 GHz.

Then we replace one single metallic pillar with a shorter one ($h_d = 4.15$ mm) and keep other experimental setup unchanged. The measured transmission spectrum of the point defect cavity is shown as a red curve in Fig. 3.1a. Clearly, a strong resonant peak emerges within the forbidden band gap of the surrounding surface-wave band-gap crystal, corresponding to the point defect cavity formed by shortening one single pillar on the structured metal surface. Using the microwave near-field imaging system, we map the field profile of the point defect cavity (E_z) at its resonant frequency 13.5 GHz, as shown in Fig. 3.1b. We can see that a tightly localized defect cavity mode is created on the textured metal surface. For comparison, we also show the simulated field profile of the point defect cavity on a transverse xy -plane and vertical xz -plane, as shown in Fig. 3.1c and Fig. 3.1d, respectively, further verify the existence of the point defect cavity mode, which are tightly localized in the defect site and decay exponentially in both horizontal and vertical directions. Note that due to the square crystal lattice symmetry,

the eigenmode of the point defect cavity exhibits four-fold rotational symmetry, which plays a critical role in the pulse routing around sharp corners in the previous Chapter and the coupled-defect surface waveguide in this Chapter. We can anticipate that if we change the crystal lattice of the surface-wave photonic crystals, the rotational symmetry of the eigenmode of the point defect cavity will be changed too. Therefore, the surface-wave photonic crystal provides a new degree freedom to manipulate designer surface plasmons using the crystal lattice.

3.2 Coupled-defect surface waveguide

We now construct a straight coupled defect surface waveguide to manipulate the propagation of surface waves by shortening a row of metallic pillars alternately, as schematically shown in the inset of Fig. 3.2a. The simulated dispersion relation of the coupled defect surface waveguide is shown in Fig. 3.2a. We can see that the dispersion of the coupled defect surface waveguide is a tight-binding-like dispersion, which is very different with the polaritonic dispersion of conventional designer surface plasmon waveguides [76] which starts from the light line and tends to zero group velocity at the band edge. This is because the propagation of surface waves in the coupled defect surface waveguide is through the weak coupling between neighboring defect cavities, which is exactly the classic wave analogue of tight-binding (TB) limit in solid state physics [77]. Therefore, similar to coupled resonator optical waveguide (CROW) [66] and coupled defect waveguide in photonic crystal [71], the dispersion relation of CDSW exhibits a tight-binding-like dispersion shape, as shown in Fig. 3.2a, rather than a polaritonic dispersion shape in conventional designer surface plasmon waveguides.

Figure 3.2 (a) Simulated dispersion of the coupled defect surface waveguide. Inset shows the schematics of the CDSW. (b) Measured near-field transmission spectrum of a perfect surface-wave band-gap crystal without any defects (black line) and a straight coupled defect surface waveguide (red line) (c) Mapped field profile (E_z) of a straight CDSW at 13.5 GHz. (d) Simulated field profile (E_z) of a straight CDSW at 13.5 GHz.

Then we measure the near-field transmission spectrum of a straight CDSW with twelve coupled point defect cavities, as shown in Fig. 3.2b. We observe a transmission band ranging from 13.35 GHz to 13.65 GHz within the forbidden band gap of the surrounding surface-wave

band-gap crystal. The measured near-field transmission spectrum of the perfect surface-wave band-gap crystal is also shown in Fig. 3.2b for comparison, which exhibits a forbidden band gap starting from 12.6 GHz. The measured and simulated mode profiles of CDSW are shown in Fig. 3.2c and Fig. 3.2d, respectively. We can see that surface waves are tightly guided along the CDSW.

3.3 Multi-directional plasmonic surface-wave splitters with full bandwidth isolation

In the past decades, several schemes have been proposed [78-84] to control the propagation direction of SPPs (or designer SPPs at low frequencies) in various plasmonic structures. For instance, the unidirectional SPPs coupler by constructing distinct SPPs waveguide structures on opposite side of slit enable slit-coupled light to be guided along different directions at different frequencies [78-81]. Recently, bidirectional SPPs splitters in the microwave frequencies were experimentally demonstrated [82-83], and then a multidirectional plasmonic surface wave splitter [84] followed. However, all the aforementioned plasmonic splitter structures have a major limitation associated with their intrinsic polaritonic dispersions, which make isolation between different splitter branches incomplete [85].

In the following, based on the tight-binding-like dispersions of CDSW, we propose and experimentally demonstrate a multidirectional plasmonic surface-wave splitter with full bandwidth isolation between different branches, which cannot be achieved using conventional plasmonic surface-wave splitters based on polaritonic dispersions [78-84]. Such structures can

work as a plasmonic wavelength-division multiplexer to split different frequency components of an EM pulse into different branches with full bandwidth isolation. Numerical simulations and near-field mapping at microwave frequencies are performed to verify our approach.

We shorten the metallic pillars alternately to create CDSWs with different heights (h) along different directions to construct a multidirectional plasmonic surface wave splitter. As shown in Fig. 3.3a, the four CDSW branches are indicated with black arrow ($h = 3.8$ mm), red arrow ($h = 4.1$ mm), blue arrow ($h = 4.4$ mm), and green arrow ($h = 4.7$ mm), respectively. And the periodicity of the CDSW is $D = 2d$. A monopole antenna, as indicated by a white arrow in Fig. 3.3a, is used to excite the coupled-defect surface modes in the four splitter branches at the same time.

The dispersion relations of these four different CDSWs that construct the multidirectional plasmonic surface-wave splitter are simulated by the eigenmode solver of a commercial software, CST Microwave Studio. In the simulation we adopt a super cell that consists of one and two half cylindrical metallic rods in y direction and 7 cylindrical metallic rod in x direction. The center pillar of the super cell is shortened from height $H = 5$ mm to h while other pillars keep unchanged. We apply electric boundary conditions in both x and z directions and periodic boundary condition in y direction. We then scan the phase shift between the two periodic boundaries from $kD = 0$ to $kD = \pi$ with 0.05π step, for each phase shift we obtain one eigenfrequency value for the eigenmode of the CDSW. Then the dispersion curves in the first Brillouin zone can be obtained. Fig. 3.3b shows the simulated dispersions of the four CDSWs with different heights of the shortened pillars (h). We can see that the four dispersions of CDSWs with different shortened pillar heights are completely separated without overlap.

Figure 3.3. (a) Photograph of the multidirectional plasmonic surface wave splitter realized on a structured metal surface (surface-wave photonic crystal) that consists of a square array of cylindrical aluminum pillars with radius $r = 1.25$ mm, height $H = 5$ mm, and periodicity $d = 5$ mm. The periodicity of CDSW is $D = 2d$. Four different coupled defect surface waveguides with various alternately shortened pillars height are indicated as black ($h = 3.8$ mm), red ($h = 4.1$ mm), blue ($h = 4.4$ mm), and green ($h = 4.7$ mm) arrows, respectively. A monopole source (white arrow) is adopted to excite the perfect multidirectional surface wave splitter. (b) Calculated dispersion relations for four different coupled defect surface waveguides that construct the multidirectional surface wave splitter. (c-e) Simulated E_z field distributions of the eigenmode of CDSW in xy , xz and yz plane at the edge of the first Brillion zone with shortened pillar height $h = 4.4$ mm.

This indicates full bandwidth isolation among different branches of the proposed

multidirectional plasmonic surface-wave splitter. Thus when the frequencies of surface waves lie within the waveguiding band of one CDSW branch of the multidirectional splitter, surface waves will not be supported by other CDSW branches and only propagate along this CDSW branch. Thus surface waves with different frequencies will split into different branches of the splitter and propagate along different directions without bandwidth overlap. The field profiles of the eigenmode of CDSW with shortened pillar height $h = 4.4$ mm in the xy , xz and yz planes are shown in Fig. 3.3c-e, respectively. We can see that the coupled defect surface modes are tightly localized around the shortened pillar (white dashed circle/square) and propagate along y direction through weak coupling between neighboring defect cavities.

To study the performance of the multidirectional plasmonic surface wave splitter, we first simulate the near-field transmission spectra and field intensity distributions of the plasmonic splitter using a commercial software CST Microwave Studio. Fig. 3.4a shows the simulated transmission spectra of the four CDSW branches of the multi-directional splitter, which are consistent with the dispersions of the four CDSW branches shown in Fig. 3.3b. We can see that surface waves with different frequencies are completely separated into different waveguiding bands without overlap. For comparison, we also show the simulated near-field transmission spectrum (grey solid line) of a perfect surface-wave band-gap crystal without defects in Fig. 3.4a. The simulated field intensity profiles of the multidirectional plasmonic splitter at 12.65 GHz, 13.15 GHz, 13.68 GHz and 14.28 GHz are shown in Fig. 3.4b-e, respectively. We can see that surface waves with different frequencies are completely split into four branches of the multidirectional splitter without complete bandwidth isolation.

Figure 3.4. (a) Simulated near-field transmission spectra of the four branches of the multidirectional splitter. The near-field transmission spectrum of a perfect surface wave band-gap crystal (grey solid line) is also shown for comparison. (b-e) Simulated field intensity distributions of the multidirectional splitter at four different frequencies 12.65 GHz, 13.15 GHz, 13.68 GHz and 14.28 GHz within different CDSW waveguiding bands.

Now we start to experimentally demonstrate this multidirectional plasmonic surface-wave splitter. A square array of (25×25) aluminum pillars with radius $r = 1.25$ mm and height $H = 5.0$ mm are fabricated. All metallic pillars stand on a flat aluminum plate with 125 mm × 125 mm in size and 5.0 mm thick. We use a pair of monopole antennas as the point source and probe. Both the monopole source and probe are connected to a vector network analyzer (R&S ZVL-13) to measure the near-field transmission spectra of the multidirectional splitter. Fig. 3.5a shows the measured transmission spectra for the four different branches of the multidirectional splitter, confirming the full bandwidth isolation of the 4-way plasmonic surface wave splitter.

Fig. 3.5b-e show the measured E_z field intensity distributions of the multidirectional plasmonic splitter at frequency 12.85 GHz, 13.35 GHz, 13.88 GHz, 14.48 GHz, respectively. We can observe that tightly confined surface waves with different frequencies propagate along different directions with full bandwidth isolation among different waveguide branches of the multidirectional splitter. The experimental results (Fig. 3.5b-e) match well with the simulation results (Fig. 3.4b-e), confirming that the multidirectional plasmonic surface-wave splitters can split surface waves with different frequencies into different branches without bandwidth overlap. This excellent performance originates from the tight-binding dispersion relation of the CDSW in the surface-wave photonic crystal which cannot be achieved by conventional designer surface plasmon waveguides.

Figure 3.5. (a) Measured near-field transmission spectra of the four branches of the multidirectional surface wave splitter. (b-e) Measured E_z field intensity distributions of the splitter for four different frequencies 12.85 GHz, 13.35 GHz, 13.88 GHz and 14.48 GHz within different CDSW waveguiding bands, respectively.

3.4 Conclusion

In conclusion, in chapter 3 we have experimentally demonstrated a waveguiding mechanism to manipulate surface waves at a subwavelength scale. The guiding of surface waves through near-field weak coupling between neighboring defect cavities is fundamentally different with conventional designer surface plasmon waveguides. Based on the CDSW, we proposed and experimentally demonstrated a multidirectional plasmonic surface wave splitters with complete bandwidth isolation among different waveguide branches at the microwave frequencies. This device employs coupled defect surface waveguiding with tight-binding-link dispersion rather than conventional designer surface plasmon waveguiding with polaritonic dispersion relations. Thus the coupled defect waveguiding bands of the four different branches of the splitter can be completely separated without overlap. Four different CDSWs are introduced by alternately shortening pillar heights on four opposite sides of the monopole source to guide surface waves at different frequencies along different directions with perfect isolation. Transmission spectra and near-field imaging measurements are conducted at microwave frequencies to verify this concept. This coupled-defect surface waveguide and multidirectional plasmonic surface wave splitter can be useful for the integrated plasmonic circuits from microwave to terahertz frequencies.

Chapter 4

Ultrathin surface-wave band-gap crystal

4.1 Ultrathin surface-wave band-gap crystal

In Chapter 2 and 3, the concepts of surface-wave band-gap crystals [58-59] were developed with superior deep-subwavelength wave manipulations of designer surface plasmons including broadband high transmission around sharp corners [58] and multidirectional plasmonic surface wave splitters with complete bandwidth isolation [59]. However, the aforementioned surface-wave band-gap crystals with square array of metallic rods standing on a metal surface have a major limitation associated with their inherent bulk geometry with thickness of at least the order of the wavelength.

Moreover, a specific type of plasmonic metamaterials, termed “spoof” or “designer” surface plasmons [22-23], has also led to subwavelength guiding propositions. Later, thin metamaterials with microwave surface-plasmon-like modes [87] and ultrathin flexible films with conform surface plasmons [31], as well as a chain of ultrathin conducting spirals with super-sub-wavelength surface modes [88] were proposed with a thickness much smaller than the operating wavelength. Yet those designs suffer markedly from serious scattering loss when encountering sharp corners because of the dramatic momentum mismatch of surface electromagnetic waves before and after passing the sharp corners in an extremely compact space. Though an ultrathin 90-degree sharp bend has been numerically proposed for designer

surface plasmons [89], it stands in air and the momentum mismatch is intrinsic before and after passing the sharp corner. Thus it only achieve high transmission at the lower frequency part of the dispersion relation (close to the light line), while at the higher frequency part of the dispersion relation (apart from the light line), the transmission is rather low due to the momentum mismatch at the sharp corner. Therefore, achieving high-efficiency transmission of surface electromagnetic waves around ultrathin sharp bends, which is vital for the miniaturization of integrated circuits, still remains a challenge.

In this Chapter, based on a well-known Sievenpiper “mushroom” structure [90] that was first proposed as a resonant high-impedance surface to suppress surface waves with thickness only a fraction of the wavelength, we show that we can tailor unit cells (metallic patches) locally to shape the flow of surface waves at deep subwavelength scales within its forbidden band gaps. We validate our approaches in experiments demonstrating various components including point defect cavities, line defect waveguide, sharp corners and T-shaped splitters. Transmission spectra and near-field imaging measurements have also been implemented in the microwave frequencies to examine the performance of these components.

We start our experimental demonstration using a medium, as schematically shown in Fig. 4.1a, that is composed of square $18\ \mu\text{m}$ thick copper patch of side length $a = 9\ \text{mm}$ arranged in a square array of pitch $d = 10\ \text{mm}$. The pitches are separated from a copper ground plane by a low loss $1.508\ \text{mm}$ thick dielectric layer [Rogers RT/duroid 5880 (Permittivity $\varepsilon = 2.2 \pm 0.02$; Loss tangent $\Delta\delta = 0.0009$)]; However, each individual copper pitch is connected to the ground plane by a hollow copper via of radius $r = 0.25\ \text{mm}$. We use COMSOL Multiphysics 5.1 to calculate the band structure of the ultrathin surface-wave photonic crystal. In band

structure simulation, a single metallic patch (with side length $a = 9$ mm) which is connected to a ground plane (with periodicity $d = 10$ mm) by a hollow copper via of radius $r = 0.25$ mm is adopted as the unit cell. The periodic boundary conditions are imposed with phase shifts in both x and y directions. And the electric boundaries are adopted in the z direction. By scanning the phase shifts in the first Brillouin zone of a square lattice as shown in the inset of Fig. 4.1b, the complete band structure of the perfect ultrathin surface-wave photonic crystal can be obtained. The numerical calculation of the band structure of this perfect high-impedance surface is shown in Fig. 4.1b. The fundamental band is TM band and the second-order band is TE band. The TM band does not reach the TE band low edge, but stops below it, forming a forbidden bandgap for surface waves from 5.1 GHz to 6.5 GHz (yellow region). The TE band slopes upward upon crossing the light line. Thus, a complete surface-wave band gap spans from the edge of the TM band to the point where the TE band crosses the light line, as indicated with yellow color in Fig. 4.1b. The ultrathin surface-wave photonic crystal investigated in this Chapter is fabricated using printed circuit board (PCB) technology. First, an array of holes is drilled completely through a 1.507 mm thick, 150×150 mm panel of Rogers RT 5880, clad either side with an 18 μm thick copper layer. Electroless deposition is subsequently used to deposit copper on the inner surface of the holes, where the outer radius of these hollow, metallic ‘pins’ is 250 μm . For the final step, the desired pattern of patches is etched into the top copper layer, so that a pin is located at the center of each patch, providing electrical connection to the continuous lower copper layer (‘ground plane’).

Figure 4.1. (a) Schematic depiction of a small section of the “mushroom” structure that consists of $18\text{ }\mu\text{m}$ thick copper patches of size $a = 9\text{ mm}$, arranged in a square lattice of periodicity $d = 10\text{ mm}$, where each patch is separated from a ground plane by a $h = 1.507$ thick dielectric ($\epsilon = 2.2 + 0.0009i$) layer, but electrically connected to it by a centrally located via of radius $r = 0.25\text{ mm}$ that is located at the center of each copper patch. (b) Calculated band structure of the ultrathin surface-wave band-gap crystal. Yellow region represents the surface modes forbidden band gap of the ultrathin surface-wave band-gap crystal. (c) Measured near-field transmission spectrum of the ultrathin surface-wave band-gap crystal.

Using a vector network analyser (R&S®ZVL13) and two homemade monopole antennas placed at opposite sides of the experimental sample, we measured the near-field transmission spectrum of the surface modes supported by the ultrathin surface-wave band-gap crystal, as shown in Fig. 4.1c. The measured spectral transmission exhibits a substantial drop from 5.1 GHz to 6.4 GHz, corresponding to the forbidden band gap of the ultrathin surface-wave band-gap crystal in Fig. 4.1b. This forbidden band gap for surface waves finds its origin in the local resonance of the metallic patch, and is therefore radically different from the bandgaps obtain in conventional photonic crystals [60], which are governed by the Bragg scattering of waves within the unit cells.

4.2 Point defect cavity in an ultrathin surface-wave band-gap crystal

Our goal of the chapter is to use this ultrathin high-impedance surface to control the propagation of surface electromagnetic waves at the deep subwavelength scales within the surface-wave forbidden band gap, rather than to suppress the surface waves to realize low-profile antennas [90]. To achieve this goal, we start by designing the most basic component that can be used to control waves at the deep subwavelength scale: a point defect cavity. Owing to the deep subwavelength periodicity of the ultrathin surface-wave band-gap crystal under study, removing a metallic patch has no effect, in contrary to what would happen in conventional photonic crystals in which the defect will support a tightly confined cavity mode [69]. However, if we replace one copper patch ($a = 9$ mm) with a slightly smaller one ($a = 8$ mm), its own resonance frequency will rise due to the decrease of the capacitive coupling

between neighboring copper patches and falls exactly within the forbidden band gap of the high-impedance surface supported by the surrounding copper patches, and a point defect is formed. As shown in Fig. 4.2a, we use a monopole antennas as the point source to excite the point defect cavity, and using another monopole antenna as the probe to measure the near-field response spectrum of the point defect cavity, as shown in Fig. 4.2b. It can be seen that an obvious resonance peak is emerged at 5.7 GHz within the surface-wave forbidden band gap of the high-impedance surface. To visualize the point defect cavity mode, we mapped its spatial distribution of the electromagnetic field inside and around the cavity using a 5 mm-long monopole antenna which was able to scan in the horizontal plane 1 mm above the sample, measuring both the amplitude and phase of the resonance wave fields. By stepping the probe with 0.5 mm increment in both x and y directions, a full 2D spatial field map can be obtained experimentally. The measured E_z field distribution of the point defect cavity is shown in Fig. 4.2c, which matched well with the simulation results shown in Fig. 4.2d. It is clear that the surface electromagnetic waves are tightly confined around the defect cavity.

Figure 4.2. (a) Fabricated sample of the point defect cavity and its experimental set up. (b) Measured near-field transmission spectra of an ultrathin locally resonant metamaterial with a single point defect (red line) and without any defect (black line). (c) Observed field pattern (E_z) when the source is near the point defect cavity at its resonance frequency 5.7 GHz. (d) Simulated field pattern (E_z) when the source is inside the point defect at its resonance frequency 5.7 GHz.

4.3 Line-defect waveguide in an ultrathin surface-wave band-gap crystal

Using the same experimental method, we now move on to more complex components to manipulate surface waves on this ultrathin locally resonant metamaterial. We first construct a straight line-defect waveguide, as shown in Fig. 4.3a, replacing a row of copper patches by smaller ones akin to our point defect cavity experiment. We first measure its near-field transmission spectrum, as shown in Fig. 4.3b (red line), a narrow waveguiding band appears from 5.5 GHz to 5.8 GHz within the surface-wave forbidden band gap of the high-impedance surface. We also plot the transmission spectrum of the high-impedance surface without defect for comparison (black line), which shows a wide forbidden band that spans from 5.1 GHz to 6.4 GHz, being consistent with the band structure shown in Fig. 4.1b. Measured and simulated maps of the spatial distributions of electric field 1 mm above the sample at 5.7 GHz are presented in Fig. 4.3c and Fig. 4.3d, respectively. Evidently, surface waves tunnel from resonator to resonator, forming a very efficient and subwavelength confined propagation.

Figure 4.3. (a) Fabricated sample of the straight line-defect waveguide constructed in the ultrathin locally resonant metamaterials. (b) Measured normalized transmission spectrum for a straight line-defect waveguide (red line) and an ideal ultrathin locally resonant metamaterial without any defect (black line). (c) Measured field pattern (E_z) of a straight line-defect waveguide at 5.7 GHz. (d) Simulated field pattern (E_z) of a straight line-defect waveguide at 5.7 GHz.

4.4 Sharp bend and T-shaped splitter in an ultrathin surface-wave band-gap crystal

Owning to the resonant nature of the unit cells in this ultrathin surface-wave band-gap crystal, we can mold the flow of surface waves in any given path through tailoring the unit cell locally at deep subwavelength scales. Thus we can construct various components to manipulate surface waves, i.e., a sharp corner and a T-shaped splitter, as shown in Fig. 4.4a and Fig. 4.4d. We first consider the sharp bend waveguide with zero bending radius. The measured

Figure 4.4. (a) Fabricated sample of a sharp bend in the ultrathin locally resonant metamaterials. (b) Observed field pattern (E_z) of a sharp corner at 5.7 GHz. (c) Simulated field pattern (E_z) of a sharp corner at 5.7 GHz. (d) Fabricated sample of a T-shaped splitter in the ultrathin locally

resonant metamaterials. (b) Observed field pattern (E_z) of a T-shaped splitter at 5.7 GHz. (c) Simulated field pattern (E_z) of a T-shaped splitter at 5.7 GHz.

near-field distributions of the sharp corner is shown in Fig. 4.4b, matching well with the simulation results shown in Fig. 4.4c. In both the experiment and simulation results, we can see that surface waves can be high-efficiently transmitted through the sharp corner without perceptible radiation and reflection. The high-efficiency guidance of surface waves around sharp corners is mainly because of the following two reasons. First, the wide surface-wave band gap of the surrounding locally resonant metamaterials suppresses scattering into the ambient space. This feature is similar to previous wavelength-scale photonic circuits in a photonic crystal [61] with full forbidden band gap. Second, because each copper patch exhibits an eigenmode with four-fold rotational symmetry, there is no difference between the 90° sharp bend and the straight line-defect waveguide. This mechanism of resonance transmission through sharp corners is similar to previous wave manipulation along a line defect in a 3D wire medium metamaterial [67] and the coupled-resonator optical waveguide [66].

In the following, we construct a T-shaped splitter by connecting two straight line-defect waveguides in the ultrathin surface-wave band-gap crystal, as shown in Fig. 4.4c. An ideal splitter should transmit 50% of the input power into each output. We can observe in the measured near-field distribution of the T-shaped splitter in Fig. 4.4e and the simulation results in Fig. 4.4f that the input surface waves are split almost equally by the T-shaped splitter with negligible scattering at the junction of the T-shaped splitter.

4.5 Conclusion

In conclusion, in chapter 4 we have proposed and experimentally demonstrated an ultrathin surface-wave band-gap crystal to control the surface wave propagation at the deep subwavelength scales. Compared with the previous bulk locally resonant metamaterials and surface-wave band-gap crystal, this structure is much thinner with thickness only a fraction of the wavelength (about $1/40\lambda$) and compatible to standard printed circuit board (PCB) technology. Functional devices such as a high-efficiency sharp bend and T-shaped splitter have also been realized, both of which cannot be achieved with traditional surface-wave waveguides.

Part II

Designer localized surface plasmons and its applications

Chapter 5

Forward/backward switching of plasmonic wave propagation using sign-reversal coupling

5.1 Forward and backward waves of surface plasmons

In general, designer surface plasmons can be classified into two different types: propagating designer surface plasmons (SPPs) which are propagating surface modes on the extended interface between metal and dielectric [21-31], or designer localized surface plasmon (LSPs) in structures with closed surfaces such as plasmonic particles [32-34]. After we have studied the various novel manipulations of propagating designer surface plasmons in the precious three Chapters, in this Chapter we focus on the novel wave manipulation of designer localized surface plasmons: forward/backward switching of plasmonic wave propagation using sign-reversal coupling in coupled designer surface plasmon resonator waveguide.

In negative refractive index materials, the phase velocity of electromagnetic (EM) waves is opposite to the direction of the power flow. This phenomenon is widely known as the backward wave propagation. This unusual class of EM waves forms the foundation of negative refraction [91] and has attracted considerable interests in the past decade in the context of photonic metamaterials [92-93]. Recently, the concept of backward waves has been further extended to surface plasmon polaritons (SPPs) [94-95]. However, in most previous approaches, the backward waves only exist in a single frequency band, and lack the capability of switching

between forward and backward waves in the same frequency band. Moreover, in previous approaches of achieving backward waves for SPPs, the backward propagation generally occurs at the anomalous dispersion region within a single narrow band, being unable to cover multiple frequency bands.

Apart from typical SPPs that propagate on a continuous metal surface [94-95], plasmonic waves can also propagate through near-field coupling between plasmonic particles supporting tightly localized photonic modes [96]. In fact, this kind of near-field coupling between tightly localized photonic modes, as an analogue of the tight-binding atomic wave coupling in solid-state systems, is of fundamental importance in many modern photonic applications ranging from slow-light engineering [97-102], quantum simulation [103-104], to topological photonics [35-41] and $\mathcal{PT}$ -symmetric photonics [105-107]. The wave propagation direction along a chain of plasmonic particles with tightly localized photonic modes is determined by the sign of the coupling factor [66,96], which stems from the overlap of the evanescent tails of tight-binding photonic modes. Although the coupling strength can be easily tuned by adjusting the inter-particle distance, it remains difficult to flip the sign of coupling factor [108-113] and consequently switch between forward and backward wave propagations.

In this chapter, by directly measuring the tight-binding Bloch waves on a periodic array of coupled designer SPP resonators (“meta-particles”) [32-33] in the microwave regime, we demonstrate that the coupling factor between tightly localized photonic modes can reverse its sign by switching between two different coupling configurations. As a consequence, two surface modes with positive and negative group velocities arise in the same frequency range, forming a Dirac-point-like crossing of dispersions [114]. We experimentally measured the

Bloch wavevectors of these modes as a function of frequency and achieved the functionality of switching between forward and backward surface Bloch waves using selective excitation. We also show that this sign-reversal coupling approach to realize forward/backward switching is applicable to any multipolar mode that is achievable in multiple frequency bands.

5.2 Sign-reversal coupling between two tightly localized resonators

We first consider two identical two-dimensional (2D) resonators hosting tightly localized photonic modes, as shown in Fig. 5.1a. Without loss of generality, we assume that the 2D resonators are made of a homogeneous and isotropic material with a constant permittivity. In a typical coupling configuration [66], the two resonators couple through facing two “poles,” or intensity maxima, of resonant wave patterns that are tightly confined around the two resonators. Defining the x -axis as in Fig. 5.1a, and setting the radii of two identical resonators as R , the

Figure 5.1. Coupling configurations between tightly confined multipolar photonic modes. (a) “Pole-pole” coupling configuration by facing intensity maxima between resonators. (b) “Node-node” coupling configuration by facing intensity minima between resonators.

multipolar resonance mode profile of each resonator can be expressed as $E_z =$

$B_m(kr)\cos(m\varphi)$, where m is the order of multipolar resonance mode, and (r, φ) are the polar coordinates with the center of each resonator as the origin. $B_m(kr) = \begin{cases} J_m(k_d r), & r < R \\ \frac{J_m(k_d R)}{H_m^1(ik_0 R)} H_m^1(ik_0 r), & r > R \end{cases}$, where $J_m(\cdot)$ is Bessel function, $H_m^1(\cdot)$ is first-kind Hankel function, and k_d and k_0 are wavevectors in resonators and the environment, respectively. The coupling factor (κ) is a mathematical expression [66] that depends on the integration of overlapped resonance wave fields of the two resonators. We can easily understand that this kind of “pole-pole” coupling is associated with a strong interaction between the two resonators.

Interestingly, there is another configuration to couple these two resonators through facing two “nodes,” or intensity minima, of resonance wave patterns between the two resonators, as shown in Fig. 5.1b. In this case, the multipolar resonance mode profile of each resonator can be expressed as $E_z = B_m(kr)\sin(m\varphi)$. Intuitively, the strength of coupling, or the magnitude of coupling factor κ , in this “node-node” coupling configuration should be relatively small as compared with that in the “pole-pole” coupling configuration. However, a rigorous derivation shows that the coupling factor κ keeps its magnitude unchanged, and simply reverses its sign. In other words, this sign-reversal coupling reverses the phase relation between binding/anti-binding photonic modes of two coupled resonators, as well as the Bloch-wave dispersion slope of a periodic resonator array, as we will demonstrate later.

To facilitate the experimental retrieval of coupling factor κ , we first describe the two coupled resonators with the coupled mode theory [115] as follows:

$$\begin{cases} -\frac{da_1}{dt} = i\omega_0 a_1 + i\kappa\omega_0 a_2 \\ -\frac{da_2}{dt} = i\omega_0 a_2 + i\kappa\omega_0 a_1 \end{cases} \quad (1)$$

where a_1 and a_2 represent resonance wave fields in the two resonators, and ω_0 denotes the

intrinsic resonance frequency of a multipolar resonance mode in a single resonator. By solving this eigenvalue problem in Eq. (1), we obtain two orthogonal eigen solutions of $[a_1 \ a_2]^T$ as $[1 \ -1]^T$ and $[1 \ 1]^T$, in which the former corresponds to the out-of-phase (a_1 and a_2 differ by a phase of π) coupled mode with eigen frequency $\omega_{1,-1} = \omega_0 - \kappa\omega_0$, and the latter to the in-phase (a_1 and a_2 have the same phase) coupled mode with eigen frequency $\omega_{1,1} = \omega_0 + \kappa\omega_0$. This relation indicates that the magnitude of coupling factor κ determines the frequency shift in mode splitting, and the sign of κ corresponds to the phase relation of binding/anti-binding modes, i.e. whether it is out-of-phase or in-phase.

5.3 Experimental demonstration of mode splitting in a coupled-resonator dimer

Figure 5.2. Experimental setup for measuring coupled resonance waves. (a) Photo of a designer SPP resonator with radius $R=12$ mm. The depth of grooves is $r=9$ mm. The width and periodicity of grooves are $a=0.625$ mm and $d=1.255$ mm, respectively. (b) Schematic of a periodic array of designer-surface-plasmon resonators. A network analyzer records the Bloch waves by scanning the near-field distribution above the resonator array.

We then proceed to experimental demonstration. We adopt recently proposed designer SPP resonators as shown in Fig. 5.3a, which are constructed by decorating periodic subwavelength grooves on circular metallic disks. Using the experimental setup in Fig. 5.2b, we first measured the near-field response spectrum of an isolated resonator, as demonstrated in Fig. 5.3b, the marked M_1 - M_6 peaks indicate the excited designer-LSPs of dipole, quadrupole, hexapole, octopole, decapole, dodecapole modes, respectively, which are located at 3.88, 5.30, 6.08, 6.45, 6.66, and 6.76 GHz. The insets show the measured E_z -field distribution (with phase information included) of the multipolar designer-LSPs on a plane 1 mm above the isolated resonator, corresponding to the marked resonance peaks in the near-field response spectrum.

Figure. 5.3 One single designer localized surface plasmon resonator. (a) The photo of one single resonator with two monopole antenna as the source and prober indicated with one pair of red dots. (b) The near-field response spectrum of a single resonator with six resonance peaks indicated with arrows. Inset shows the measured near-field (E_z) distribution of the lowest six resonance modes on a plane 1 mm above the resonator.

Designer SPPs are tightly confined EM modes on patterned metallic structures with

subwavelength corrugations, whose dispersion properties and spatial confinements are akin to those of natural SPPs at optical frequencies [32]. Recently, designer SPP resonators are shown to support localized designer SPPs [32-33], opening an opportunity to image directly the near-field coupling mechanism. Because of the analogy between localized EM resonance modes and atomic wave functions, these designer SPP resonators are termed as a type of “meta-atoms” [33] in the field of metamaterials. Another benefit of this designer-SPP-resonator metamaterial platform is that it allows selective excitation of multipolar resonance modes with opposite parity, meaning that the analogous “atomic wave functions” can be selectively controlled. We construct a periodic array of designer SPP resonators, as shown in Fig. 5.2b. A monopole source is employed at one side of the array to selectively excite a multipolar resonance mode. The Bloch waves on the array is recorded by a near-field scanning probe, connected to a microwave network analyzer.

Figure 5.4. Mode splitting of a coupled-resonator dimer. (a) The measured near-field transmission spectra through a coupled-resonator dimer with end and side excitations. Inset: photo of the two designer-surface-plasmon resonators with inter-resonator distance $\Lambda = 30$ mm.

(b) The measured resonance wave patterns. The position of source is indicated by a red arrow.

We now consider coupling between two designer SPP resonators. Fig. 5.4a shows the measured near-field transmission spectra through this coupled-resonator dimer. We adopt two excitation configurations. Firstly, we locate the source and the probe in symmetric positions at two ends of the dimer, as indicated by a pair of red dots in the inset of Fig. 5.4a. Secondly, the source and the probe are both placed by one side of two resonators, as indicated by a pair of blue dots in the inset of Fig. 5.4a. We call the first configuration as “end excitation” and the second as “side excitation.” Each of these excitation configurations will selectively excite one of the multipolar resonance modes with opposite parity. The transmission spectrum for a single resonator is also measured for comparison by placing the source and probe at opposite sides of the resonator. Three resonance modes can be observed in a single resonator at 5.34 GHz, 6.07 GHz, and 6.44 GHz as shown in Fig. 5.4a. They correspond to the quadrupole (labeled as “Q”), hexapole (labeled as “H”), and octopole (labeled as “O”) modes, respectively, as will be demonstrated later.

Let us consider the end excitation. When the coupled-resonator dimer is excited with the end excitation, each of the three resonance modes in a single resonator splits into binding and anti-binding modes: from one mode at $Q=5.34$ GHz to two modes at $Q_{\pi}=5.19$ GHz (binding mode) and $Q_0=5.49$ GHz (anti-binding mode), from one at $H=6.07$ GHz to two at $H_0=6.03$ GHz (binding mode) and $H_{\pi}=6.11$ GHz (anti-binding mode), and from one at $O=6.44$ GHz to two at $O_{\pi}=6.42$ GHz (binding mode) and $O_0=6.47$ GHz (anti-binding mode). Here the subscript “ π ” or “0” denotes out-of-phase or in-phase phase relation for the two resonators after mode

splitting. To explicitly show the phase relation, we measure the resonance wave pattern for each resonance mode by scanning the probe in a transverse plane 1 mm above the coupled-resonator dimer, as shown in Fig. 5.4b. The source position is indicated by a red arrow. We observe that for the coupled quadrupole and octopole resonance modes, the two resonators are out-of-phase at lower resonance frequencies (Q_π and O_π), and in phase at higher resonance frequencies (Q_0 and O_0). However, the situation is reversed for the hexapole mode: the two resonators are in-phase at the lower resonance frequency (H_0), and out-of-phase at the higher resonance frequency (H_π). The different phase relation of two coupled resonators for the coupled hexapole resonance mode is simply because the hexapole mode is an odd mode (mode order $m=3$) while both the quadrupole mode (mode order $m=2$) and the octopole mode (mode order $m=4$) are even modes.

Using the eigen-frequency analysis from Eq. (1), we can experimentally retrieve the coupling factor κ from the measured split resonance frequencies and the phase relation between the coupled two resonators. We obtain that with end excitation, $\kappa_Q = 0.0281$, $\kappa_H = -0.0066$, and $\kappa_O = 0.0039$ for the coupled quadrupole, hexapole, and octopole resonance modes, respectively. It is apparent that a higher-order multipolar resonance mode has a smaller magnitude of coupling factor due to its tighter field confinement. The negative sign of κ_H for the hexapole resonance mode arises from the phase relation of binding/anti-binding modes, which is opposite to that in the quadrupole and octopole resonance modes.

Now we consider the side excitation. It can be seen in Fig. 5.4a that the split resonance frequencies match well with those in the end excitation. For the quadrupole and octopole modes, because of their four-fold rotational symmetry, the resonance wave patterns for their split

modes are almost the same with those in the end excitation. We thus do not discuss them here. However, it is not the same situation when it comes to the hexapole mode. As shown in Fig. 5.4b, we measure the resonance wave patterns for the two split hexapole modes ($H_{\pi}^*=6.03$ GHz and $H_0^*=6.11$ GHz) in the coupled-resonator dimer, where the source position is indicated by a red arrow. While all previous resonance wave patterns correspond to the “pole-pole” coupling configuration as in Fig. 5.1a, here these two resonance wave patterns correspond to the “node-node” coupling configuration as in Fig. 5.1b. Although frequencies of the two split hexapole modes are consistent with those in the “pole-pole” coupling configuration, the phase of these two split modes has been exchanged: the in-phase mode (H_0^*) now has a higher resonant frequency than the out-of-phase mode (H_{π}^*). Following the process above to retrieve the coupling factor, we can get $\kappa_H^* = 0.0066$ for the hexapole mode in the “node-node” coupling configuration. This confirms the sign reversal of coupling factor compared to the previously retrieved $\kappa_H = -0.0066$ for the hexapole mode in the “pole-pole” coupling configuration.

This mode splitting is analogous to that of the electron orbitals in diatomic molecules such as H_2^+ , in which when two atoms (cavity) are brought close together, the interaction and field hybridization will produce two new energy levels which lie slightly below (*bonding* mode) and above (*antibonding* mode) the frequency of the original level (isolated cavity). In these electronic molecules the bonding strength is determined by the Coulomb force between the nuclei and the electrons. Instead, the electromagnetic coupling between neighboring plasmonic molecules (cavities) is determined by the geometry of the molecules, thus can be altered by changing the inter-cavity distance, the surrounding medium or the sizes of the coupled resonators. We restrict ourselves here to the influence of the inter-cavity distance D for the two

coupled cavities in the case of front excitation. The measured transmission spectrums are shown in Fig. 5.5a with varying the inter-cavity distance D . Each spectrum consists of several resonance peaks corresponding to different splitting resonant modes. With decreasing the inter-cavity distance D the separation between the split resonance increases, as expected from the tight-binding formalism that the two cavities couple through the overlap of their evanescent modal tails. To highlight the main phenomena, we extract the spectral positions of the transmission peaks of the splitting quadrupole, hexapole and octopole modes and present in Fig. 5.5b. It is apparent that the two splitting resonant modes demonstrate an opposite spectral behavior that the symmetric (antisymmetric) modes are shifting to lower (higher) frequencies when increasing the inter-cavity distance D . Consequently the spectral splitting decreases as the inter-cavity distance increase, which implies a weaker EM coupling between the two resonators. Our results are the first direct experimental imaging of the *bounding/antibonding* mechanism in designer plasmonic resonators which has only been studied in optical frequencies.

Figure. 5.5 Mode splitting with varying inter-cavity distance. (a) The measured near-field response spectrums with three different inter-cavity distance D . (b) The relationship between

the two splitting resonance frequencies and the inter-cavity distance D for quadrupole, hexapole and octopole mode, respectively.

5.4 Forward/backward switching of plasmonic wave on a straight CROW

We can proceed to construct a coupled-resonator optical waveguide (CROW) [66] with a periodic array of designer SPP resonators, as shown in Fig. 5.6a. Each multipolar resonance mode in a single resonator now spreads continuously to form a transmission band on the CROW. We describe the CROW that consists of an infinite chain of designer SPP resonators with coupled mode theory as follows:

$$-\frac{da_n}{dt} = i\omega_0 a_n + i\kappa\omega_0 a_{n-1} + i\kappa\omega_0 a_{n+1} \quad (2)$$

where a_n , a_{n-1} and a_{n+1} denote resonance fields in the n -th, $(n-1)$ -th and $(n+1)$ -th resonators, respectively. The periodicity of this structure allows the application of Bloch theorem which gives $a_{n+1} = a_n \cdot e^{iK\Lambda}$, where K represents the wavevector of Bloch waves on the CROW, and Λ is the periodicity of designer SPP resonators. Then the intrinsic dispersion relation of the Bloch waves on the CROW can be obtained as $\omega = \omega_0 [1 + 2\kappa \cos(K\Lambda)]$, being consistent with Yariv's original derivation [66]. Apparently, the sign of coupling factor κ affects the relation between Bloch wavevectors and frequencies.

Figure 5.6. Bloch waves on a periodic resonator array. (a) Photo of the designer SPP array with periodicity $\Lambda = 30$ mm. Positions of the source and probe for the end and side excitations are indicated as a pair of red dots and blue dots, respectively. (b) Dispersion curves calculated with experimentally retrieved coupling factors for the quadrupole, hexapole, and octopole modes. Black solid lines correspond to the “pole-pole” coupling configuration. The red solid line

corresponds to the “node-node” coupling configuration for the hexapole mode. The black dashed lines correspond to the “node-node” coupling configuration for the quadrupole and octopole modes. (c) The measured transmission spectra through the array. The transmission spectrum without the array is also measured for comparison. (d-g) Measured Bloch waves of the quadrupole (d), octopole (e) and hexapole (f, g) modes with three characteristic Bloch wavevectors as indicated by three dots in each dispersion curve in (b). The red arrow indicates the position of source.

Given the experimentally retrieved coupling factors, we plot with solid black lines in Fig. 5.6b the dispersion curves for the quadrupole, hexapole, and octopole resonance modes along the CROW in the “pole-pole” coupling configuration. This coupling configuration can be achieved with the end excitation, by locating the source and probe at positions as indicated by a pair of red dots in Fig. 5.6a. Working bandwidth gets narrower as the mode order gets higher, because the tighter field confinement in a higher-order mode leads to a smaller magnitude of κ . Because of the positive sign of κ_Q and κ_O , the dispersion curves of Bloch waves for the quadrupole and octopole resonance modes exhibit properties of “backward waves,” in which the phase velocity and the group velocity are in opposite directions. Regarding the hexapole resonance mode, because of its negative coupling factor κ_H , its Bloch waves are “forward waves,” whose phase velocity and group velocity are in the same direction.

Similar to previous discussions on the coupled-resonator dimer, the side excitation in the CROW will turn on the “node-node” coupling configuration for the hexapole resonance mode with a sign-reversal coupling factor, i.e. $\kappa_H^* = -\kappa_H$. The setup of end excitation is indicated by

a pair of blue dots in Fig. 5.6a. We plot the dispersion curve for the hexapole resonance mode in the “node-node” coupling configuration in Fig. 5.6b with a red solid line. This dispersion curve shares the same bandwidth with that of the hexapole resonance mode in the “pole-pole” coupling configuration. The crossing of the two dispersions gives rise to a Dirac point [114]. Because of the sign reversal of coupling factor, Bloch waves of this dispersion curve in the “node-node” coupling configuration exhibit properties of “backward waves,” where the phase velocity and group velocity are in opposite directions.

We adopt two steps in experiment to demonstrate the above predictions. In the first step, we measure transmission spectra to verify transmission bandwidth. The results are shown in Fig. 5.6c, being consistent with dispersion curves in Fig. 5.6b. In the second step, we directly image Bloch waves propagating on the CROW. We select three characteristic Bloch wavevectors ($K\Lambda = 0, \frac{\pi}{2}, \pi$) as indicated by three dots in each dispersion curve in Fig. 5.6b. For the quadrupole mode in the “pole-pole” coupling configuration, as shown in Fig. 5.6b, the frequency of Bloch waves decreases as the Bloch wavevector increases. A similar situation applies to the octopole mode in the “pole-pole” coupling configuration, as shown in Fig. 5.6e. These results confirm their “backward-wave” properties.

What is more interesting is the hexapole mode that can be excited with both the “pole-pole” and “node-node” coupling configurations, where sign reversal of the coupling factor is expected. Fig. 5.6f shows the measured Bloch waves of the hexapole mode in the “pole-pole” coupling configuration. It can be seen that the frequency of Bloch waves increases as the Bloch wavevector increases. This confirms the “forward-wave” properties of these Bloch waves. In Fig. 5.6g, the measured Bloch waves are plotted for the hexapole mode in the “node-node”

coupling configuration. Instead, the frequency of Bloch waves decreases as the Bloch wavevector increases. Apparently, they are “backward waves.” The change of relationship between the frequency of the Bloch waves and the Bloch wavevectors determines the surface waves are “forward waves” or “backward waves”. Through near-field imaging we directly determined the three characteristic Bloch wavevectors ($K\Lambda = 0, \frac{\pi}{2}, \pi$) and their frequencies, thus we determined the change of the relationship between the Bloch waves and the Bloch wavevectors (“forward waves” or “backward waves”) in the first Brillouin zone, as indicated by three dots in each dispersion curve in Fig. 5.6b.

Note that the sign-reversal coupling for the quadrupole (octopole) modes can also be achieved by locating the excitation position at the middle (quarter) between the current positions of end excitation and side excitation. We also plot in Figure 4b the dispersion relations of quadrupole (octopole) mode (black dashed lines) for the “node-node” coupling configuration, when the excitation is relocated at proper positions. We also plot in Fig. 5.6b the dispersion relations of quadrupole and octopole mode (black dashed lines) for the “node-node” coupling configuration, when the excitation is relocated at the middle between the current positions of end excitation and side excitation. We can see that the surface modes will also switch from backward waves to forward waves in the same frequency range (experimental data now shown for brevity). Therefore, the sign-reversal coupling approach to realize the switching between forward and backward waves is general, being applicable to any multipolar resonance mode in multiple frequency ranges.

5.5 Sign-reversal coupling in a bent CROW

In the following, we will show that the sign-reversal coupling can occur in a single CROW, in which the Bloch waves with both positive and negative coupling factors can co-exist. We construct a right-angle bent CROW as shown in Fig. 5.7a. Similar to previous demonstrations, the end (side) excitation is achieved by placing the source and probe at positions indicated by a pair of red (blue) dots. Because of the four-fold rotational symmetry of mode profiles, the guidance of Bloch waves for the quadrupole and octopole modes will not be affected by the bending corner. In other words, their coupling factors will not change over the bending corner.

Figure 5.7. Sign-reversal coupling in a bent resonator array. (a) Photo of the bent designer SPP array. Positions of the source and probe for the end and side excitations are indicated as a pair

of red dots and blue dots, respectively. (b) Measured transmission spectra through the bent resonator array. The transmission spectrum without the bent resonator array is also measured for comparison. (c-f) Measured Bloch waves of the quadrupole (c), octopole (d) and hexapole (e,f) modes. The red arrow indicates the position of source.

On the other hand, the Bloch waves of the hexapole mode, when they propagate across the bending corner, will switch from the “pole-pole” coupling configuration to the “node-node” coupling configuration, or vice versa, and thus induce the sign reversal of coupling factor at the bending corner. We first measure the transmission spectra as shown in Fig. 5.7b, whose results are consistent with dispersion curves in Fig. 5.7b in the main text. We then image the Bloch waves. Both the quadrupole mode (Fig. 5.7c) and the octopole mode (Fig. 5.7d) are excited in the “pole-pole” coupling configuration. The sign-reversal coupling does not occur at the bending corner. For the hexapole mode excited with end excitation (Fig. 5.7e), the “pole-pole” (“node-node”) coupling configuration applies to the horizontal (vertical) arm. The coupling factor changes its sign from negative to positive at the bending corner. The hexapole mode can also be excited with side excitation (Fig. 5.7f), where the “node-node” (“pole-pole”) coupling configuration applies to the horizontal (vertical) arm. This time the coupling factor changes its sign from positive to negative at the bending corner.

5.6 Conclusion

In summary, in this chapter we have revealed a new way to flip the propagation directions of plasmonic waves along a chain of plasmonic particles by using the sign-reversal coupling between tightly localized photonic modes. On the platform of designer SPP resonators, we confirm the sign reversal of coupling factors by experimentally monitoring the phase of resonance wave fields in the coupled two photonic modes. By monitoring Bloch waves in a periodic array of designer SPP resonators, we show that the sign-reversal coupling will induce two different waveguiding dispersion curves in the same frequency range with opposite group velocities. The forward wave and backward wave can be switched within the same frequency range through changing the excitation position. This sign-reversal coupling approach can be applied to any multipolar resonance modes in multiple frequency ranges. In view of the demanding coupling tuning and backward waves, we expect that our study will find novel use in designs of functional photonic systems and metamaterials.

Chapter 6

Flexible photonic topological insulator

6.1. Photonic topological insulator

In the previous four Chapters, we have studied both the propagating designer surface plasmons and the localized designer surface plasmons with novel wave manipulations and applications. Due to its highly tenability and ease to directly image the near-field mode profiles, designer surface plasmon structure is an excellent platform to explore various physical phenomena and mechanisms. Owing to its vertical confinement in free space, designer surface plasmons have been successfully employed in detecting topological phenomena of classical waves in recent works. In this Chapter, we will further extend the concept of recently emerged topological photonics to designer surface plasmons and demonstrate a flexible photonic topological insulator.

A topological insulator (TI) is an emerging electronic material that is insulating in the bulk of a sample, but conductive on the surfaces via a family of “topological edge states” [116-117]. It has spurred tremendous research interests in constructing analogous topological effects in classical wave systems [118-123]. Photonic TIs, as an analogue of TIs in photonic systems [37-47] have been under investigation on various platforms. However, currently proposed photonic TIs usually require covers to confine the waves in the vertical direction to prevent radiative leakage into free space. In experiments, the covers will largely hamper the probing inside the

photonic crystals and the experimental mapping of the fields inside the structure becomes a difficult challenge. As a result, there are few direct experimental observations of topological edge states. In order to achieve direct observations, the edge states need to be vertically confined in free space to eliminate covers, and the elimination of covers is also beneficial for both manufacture and applications.

Moreover, all of these proposed photonic TIs are inherently rigid without elastic flexibility. On the technical level, flexibility is extremely desirable in flexible electronics [126-128] and flexible photonics [129-134], which aim at achieving electronic and photonic circuits, devices, and systems on flexible, stretchable, and biocompatible curved substrates. Promising applications include personal health monitors and biomedical devices [126], electronic eyeball camera [127] paper-like display [128], flexible nano-membrane photonic crystal cavities [129-131] stretchable silicon photonic waveguides and filters [132] and flexible nanoantennas array [133]. To date, however, a flexible (electronic or photonic) TI still remains unrealized.

On the scientific level, a flexible (electronic or photonic) TI that can be bent and folded allows involvement of spatial topology in manipulating topologically protected edge states. A typical nontrivial spatial topology is the Mobius strip whose upper and lower edges are indistinguishable. Although the topology of a Mobius strip has been implemented in a topological electronic circuit which does not rely on the mechanical flexibility [134] the photonic version of a similar nontrivial topology has not been developed.

In this Chapter, we experimentally realize a paper-like flexible photonic TI and demonstrate the measured near-field maps of the topological edge states in the microwave regime. The flexible photonic TI is constructed using an ultrathin designer surface plasmon structure, which

is essentially a periodic metallic structure supporting non-leaky electromagnetic (EM) surface modes, similar to the famous surface plasmons at optical and infrared regimes but arise at much lower frequencies.

Topologically protected photonic states can be formed on the flexible photonic TI, regardless of it being flat or curved. Owing to the extraordinary flexibility, we further fold the paper-like photonic TI and connect two edges, in order to simulate different spatial topologies. The photonic pseudospin will flip or preserve its spin direction, depending on the folding manner, and will subsequently propagate in different directions along different edges. Our work shows that the propagation of topological edge states can be manipulated by introducing spatial topologies to a flexible topological system.

6.2. Flexible photonic topological insulator

To experimentally demonstrate the flexible topological system, we apply conformal surface plasmons (CSP) [31] and standard printed circuit board (PCB) technology. CSP are surface electromagnetic waves that can propagate on ultrathin and flexible films over a long distance in a wide broadband range, covering from far-infrared to microwave frequencies. CSP hold considerable promises of controlling the propagation of surface waves on curved surfaces conformably to produce advanced plasmonic functional devices. The proposed flexible photonic TI is shown in Fig. 6.1a. It consists of a square 5×5 lattice of 18- μm -thick CSP ring resonators (structured copper disk) printed on a 0.254-mm-thick paper-like dielectric film (Rogers RT5880). The ring resonators sitting on the 5×5 sites are called “lattice rings” and

those that couple neighboring lattice rings are called “coupling rings”. Surface waves can circulate clockwise or anti-clockwise in the lattice rings, serving as two components of a pseudo-spin, that is, spin-down (clockwise circulation) and spin-up (anti-clockwise circulation) states [124-125]. Previous studies have demonstrated that such a periodic system can host a topologically nontrivial “anomalous Floquet TI phase” [135], which cannot be predicted with the conventional Chern number topological invariant.

A photo of the fabricated flexible photonic TI is shown in the inset of Fig. 6.1a. Each ring resonator is formed by 48 gratings. The periodicity, width and height of the gratings are $d = 4.0$ mm, $a = 2.0$ mm and $h = 4.0$ mm, respectively. The radius of each ring resonator is $R = 30.57$ mm. The inter-ring distance is $g = 2.0$ mm. By taking the phase delay ($\emptyset$) of surface waves along a quarter of a lattice ring resonator as the quasi-energy, it can be shown that a topological phase transition will occur when the inter-ring coupling strength increases beyond a critical value ($\theta = 0.25\pi$ in Ref. 124). Owing to its ultra-thin thickness, the fabricated photonic TI can be bent, folded, or even twisted to mold the flow of the topological edge states on curved surfaces, as schematically shown in Fig. 6.1b.

Figure 6.1. (a) Schematic of the flexible photonic topological insulator. Inset is a photo of the

CSP ring resonators printed on a flexible paper-like dielectric film. (b) Schematic of the folded flexible photonic topological insulator.

6.3. Dispersion and eigenmode field profile of conformal surface plasmons

Fig. 6.2a shows the schematic of conformal surface plasmon (CSP) waveguide we use to construct the flexible photonic topological insulator, which consists of one dimensional periodic gratings corrugated on a metal strip with thickness $t = 0.018$ mm. The periodicity, width, and height of the gratings are $d = 4$ mm, $a = 2$ mm, and $h = 4$ mm, respectively. We assume the metal is perfect electric conductor (PEC) and use a finite element method (FEM) to calculate the dispersion relation of conformal surface plasmon propagating along y-direction with wave vector k . The simulated dispersion relation is shown in Fig. 6.2b, which deviates significantly from the light line (black line), indicating that the corrugated metal stripe supports highly confined surface waves. This property is illustrated in Fig. 6.2c, which shows the $|E_z|$ field evaluated in a xz plane perpendicular to the CSPs propagation direction (y). Clearly, the electromagnetic wave is highly localized around the comb-shaped metal structure.

Figure 6.2. Design of conformal surface plasmon (CSP) structure. (a) Schematic of the conformal surface plasmon structure, with periodicity $d = 4$ mm, grating height $h = 4$ mm, grating width $a = 2$ mm, and stripe thickness $t = 0.018$ mm. (b) Calculated dispersion

relation of the fundamental CSP mode. The black solid line is the light line. (c) The eigenmode profile ($|E_z|$) at the center xz plane of the grating.

6.4. Calculated band structure of flexible photonic topological insulator

Fig. 6.3a shows the simulated band structure for a semi-infinite strip of the photonic TI, which has five lattice rings in the y direction, and is infinite in the x direction. Here we only consider a single spin (spin-down) and assume no spin mixing. The band diagram reveals a gap between 10.4 and 12.4 GHz, spanned by unidirectional states localized to opposite edges of the flexible TI. Note that within such a narrow bandwidth, the inter-ring coupling strength is mainly determined by the inter-ring spacing g . We estimate that for the current setting with $g = 2.0$ mm, the inter-ring coupling strength is about $\theta = 0.42\pi$ within the bandgap of interest [124]. According to calculation from the network model, this bandgap is topologically nontrivial. In the measurement, the spin-down (up) state can be excited via the input port 1 (2) of a U-shape waveguide, as illustrated in Fig. 6.3a. Fig. 6.3b shows the measured transmission spectrum of topological edge states from the input port 1 to another U-shaped waveguide at the right side of the photonic TI (as shown in Fig. 6.1a). It can be seen that the transmission matches well with the calculated band structure in Fig. 6.3a.

Figure 6.3. (a) Simulated band diagram of the flexible photonic topological insulator realized with CSP ring resonators. (b) Measured transmission spectrum of the topological edge states.

6.5. Experimental demonstration of flexible photonic topological insulator

We then adopt a microwave near-field imaging system [136] to map the near-field distributions on this topological structure. The near-field imaging system consists of a vector network analyzer (R&S ZVL-13) and two homemade monopole antennas as the source and probe. The probe, placed 1 mm above the sample, can move freely to map the electric near-field distributions. First, we excite a central site in the bulk of the planar topological structure, as schematically shown in Fig. 6.4a. The operating frequency is fixed at 11.6 GHz. The insulating nature of the bulk is revealed by the localized resonant mode emerged at the vicinity of the source, as shown in Fig. 6.4b. Then we wrap the flexible photonic TI on a foam cylinder with radius of 12.5 cm and excite the same central site. Similar field pattern can still be observed on the curved cylindrical surface, as shown in Fig. 6.4c.

Figure 6.4. (a) Schematic of the flexible photonic TI with the excitation at a central lattice site. (b) Observed field pattern when the excitation is inside the bulk on a planar surface. (c) Observed field pattern when the excitation is inside the bulk on a curved cylindrical surface. (d) Schematic of the flexible photonic TI with the excitation at the edge. (e) Observed topological edge state on a planar surface. (f) Observed topological edge state on a curved cylindrical surface. (g) Schematic of the flexible photonic TI with an introduced defect (one lattice ring and three surrounding coupling rings removed) on the edge. (h) Observed

topological protection on a planar surface. (i) Observed topological protection on a curved cylindrical surface.

Next, we excite the flexible topological system via input port 1 indicated by a red arrow, as schematically shown in Fig. 6.4d. The frequency is at 11.6 GHz. Fig. 6.4e shows that in the planar photonic TI, surface waves remain confined to the edges, propagating around the corner. Then we wrap the flexible topological structure on a foam cylinder and map the field pattern of the edge state again, as shown in Fig. 6.4f. It can be seen that the topological edge state on a curved surface does not change significantly.

To probe the robustness of the topological edge states, we fabricate a new sample with a defect, which is formed by removing a lattice ring and its three coupling rings on one edge, as shown in Fig. 6.4g. The frequency remains at 11.6 GHz. As shown in Fig. 6.4h, in the planar photonic TI, because of the topological protection, the edge states detour around the defect and continue to propagate along the original edge. Then we wrap the structure on the cylindrical foam to test the robustness of the topological protection on a curved surface. The measured results (Fig. 6.4i) show that the topological protection survives even on a curved surface.

6.6. Implementing global topologies to flexible photonic topological insulator

Due to the flexibility of the fabricated photonic TI, spatial topologies can be constructed via folding and connecting. Here we implemented two spatial topologies by overlapping the last

ring resonator at the top edge to the first ring resonator at the bottom edge [red disks in Fig. 6.5a and 5d] with different folding manners. The first topology is realized by twisting the paper-like photonic TI before connecting the corresponding ring resonators, as shown in Fig. 6.5b. This kind of Mobius-like topology can flip the spin. To facilitate experiment, we only probe one point along each ring resonator and plot the electric field distribution at 11.6 GHz, as shown in Fig. 6.5c. It can be seen that the original spin-down state on the upper edge is now converted to the spin-up state on the lower edge, both propagating in the same direction. Then we implement another spatial topology, as shown in Fig. 6.5e, which is realized by overlapping the corresponding ring resonators without twisting the photonic TI. The measured field pattern in Fig. 6.5f shows that the spin-down state on the upper edge still remains its spin state when being coupled to the overlapped ring resonators. This spin-down state keeps propagating upwards on the left edge of the photonic TI.

Figure 6.5. (a-b) Schematic of folding the flexible photonic TI with twisting. The red (green) line indicates the path of spin-down (spin-up) states. (c) Observed field pattern of topological edge state for the folding in (b). (d-e) Schematic of folding the flexible photonic TI without twisting. The red line indicates the path of spin-down state. (f) Observed field pattern of topological edge state for the folding in (e).

6.7. Conclusion

In summary, we have realized experimentally a mechanically flexible photonic TI that supports robust topological edge states on a curved surface. Two spatial topologies are implemented with different folding manners to manipulate the propagation of topological edge states which are inaccessible in conventional bulky and rigid topological systems. Such a flexible platform provides a versatile approach towards controlling and manipulating topological edge states on curved surfaces.

Summary and outlook

In this thesis, we have experimentally studied the novel wave manipulation of designer surface plasmons and its various applications at microwave frequencies.

First, we proposed and experimentally demonstrated a surface-wave band-gap crystal implemented on a single metal surface, which can be used to route surface-wave pulse around multiple sharp corners without perceptible scattering. The remarkable high-efficiency time-domain surface-wave pulse routing around multiple sharp corners in the surface-wave photonic crystals is mainly because of the following two reasons. First, the wide forbidden band gap of the surface-wave photonic crystal suppresses scattering of surface waves into the ambient space in a broad bandwidth. Second, the eigenmode supported by the line-defect waveguide of the surface-wave photonic crystal is essentially coupled toroidal resonance mode with four-fold rotational symmetry, the momentum of the surface EM waves before and after passing the sharp corners in an extremely compact space is inherently matched.

Second, we studied a new type of waveguiding mechanism for coupled surface defect modes in a surface-wave band-gap crystal to manipulate the flow of surface waves in deep subwavelength scales. Based on this new waveguiding mechanism, we realize a multi-directional plasmonic splitter with full bandwidth isolation, which cannot not be achieved using conventional designer surface plasmons with polaritonic dispersions.

Third, using the well-known Sievenpiper “mushroom” structure, we proposed and experimentally demonstrated the deep-subwavelength manipulation of surface waves on an ultrathin surface-wave band-gap crystal with a thickness very much smaller than the

wavelength. Functional components including T-shaped splitter and sharp bend are also demonstrated with good performance.

Fourth, on a platform of coupled designer surface plasmon resonators in the microwave regime, we demonstrated multi-band forward/backward switching of plasmonic wave propagation using the sign-reversal coupling that occurs when switching the coupling configuration between tightly confined photonic modes. We also demonstrate that this sign-reversal coupling mechanism is general and can be applied to any multipolar resonance modes in multi frequency bands.

Finally, using conformal surface plasmons we proposed and experimentally demonstrated a flexible photonic topological insulator which can be bended, folded, or even twisted to mold the flow of topological edge states and introduce exotic global topologies to manipulate the topological edge states, which cannot be achieved in conventional bulk and rigid photonic topological systems.

Due to its high tunability and ease of fabrication and characterization with direct mode profile imaging, designer surface plasmon structures provide an excellent platform to explore various physical mechanisms and phenomena. For example, in this thesis we have demonstrated the photonic band gap nature of designer surface plasmons to mimic the conventional photonic crystals and the topological protected edge states in the recently emerged concept of topological photonics. We anticipate that, in the near future, new emerged physics such as graphene, valleytronics (quantum valley Hall effect), spintronic (quantum spin Hall effect) can be realized using designer surface plasmon structures with various crystal lattices to mimic the novel concepts and properties in the condensed matter research fields.

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