Acoustic topological one-way waveguides with tunable widths using spinning components

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Abstract
We propose the topological one-way waveguide for acoustic waves whose width can be flexibly adjusted. The waveguide is constructed by a heterostructure where an ordinary phononic crystal is sandwiched by two time-reversal-symmetry-broken (TRS-broken) phononic crystals with their cylinders spinning in an opposite manner. The waveguide mode is confined to the ordinary phononic crystal and exhibits the gap-less and asymmetric dispersion. Therefore, we can tune the width of the waveguide by adjusting the thickness of the ordinary phononic crystal, and the waveguide mode is one-way transport which is robust against various types of local disorders and arbitrary bends. Owing to these, this acoustic topological one-way waveguide can meet the requirements of more applications compared with conventional waveguides and conventional one-way waveguides based on chiral surface waves.

1. Introduction
Phononic crystals [1–3] and acoustic metamaterials [4–6] provide versatile platforms for manipulating the acoustic waves in novel ways, such as super-lensing [7], cloaking [8], unidirectional radiation [9–11] and one-way waveguiding [12, 13]. Recently, there has been a surge of interest in topological physics. Besides abundant fundamental research [14–22], the characteristic of unidirectional boundary transport of topologically nontrivial surface states provides a robust route to the one-way waveguiding and inspires various applications [23–33]. The chiral surface state [34, 35] which is supported in quantum Hall systems is possibly the most fascinating one since its one-way transport is robust against almost all kinds of perturbations. In classical wave systems, the chiral surface state was first theoretically predicted [36–38] and experimentally realized [39, 40] in photonics by using two-dimensional magneto-optical photonic crystals, where the biased magnetic field breaks time-reversal symmetry (TRS) and induces a nonzero Chern number of the band. Later, using circulating fluid flows to break TRS in a phononic crystal, the acoustic counterparts were theoretically proposed and studied [41–43]. Based on this idea, the acoustic chiral surface state was experimentally observed recently [44].

Although the one-way waveguide is a direct application of the chiral surface state, in most cases, since the domain wall that separates the two topologically distinct materials cannot be very thick, the chiral surface wave has to be localized inside a narrow channel. Therefore, the carried energy usually cannot be so high, otherwise the nonlinear effect is unavoidable. On the other hand, limited by the thin waveguide channel, the one-way waveguides based on chiral surface waves will face many limitations in potential device applications, for example interfacing with conventional waveguides of much wider widths. Recently, M Wang et al theoretically and experimentally demonstrated that the width of the one-way waveguide channel can be greatly enlarged if an ordinary photonic crystal carrying Dirac points is used as the domain wall to separate two magneto-optical photonic crystals with gapped Dirac cones [45]. Later, large-area one-way surface magnetoplasmons [46] and photonic topological valley-locked waveguides with tunable widths [47] were also demonstrated. We are also aware that J-Q Wang et al experimentally realized the extended topological valley-locked surface acoustic waves very recently [48].
In this work, we proposed the acoustic topological one-way waveguide with a tunable width which overcomes the shortcomings of the conventional one-way waveguides based on the chiral surface modes, such as the excessive concentration of energy and the lack of width degree of freedom. Instead of connecting two TRS-broken phononic crystals directly \[41–44\], the waveguide is formed by a heterostructure where an ordinary phononic crystal is sandwiched by two TRS-broken phononic crystals. The ordinary phononic crystal is composed of cylinders arranged into a triangular lattice and carries Dirac points, while the two TRS-broken phononic crystals are the same as the ordinary one except that their cylinders are spinning in an opposite manner which lifts the Dirac points to generate band gaps with opposite nonzero gap-Chern numbers. As the ordinary phononic crystal acts as the domain wall to separate the two topologically distinct TRS-broken phononic crystals, two waveguide modes with their group velocities in the same direction are supported. The dispersion of the waveguide mode is gapless and very similar to that of the chiral surface mode supported on the interface formed by the two TRS-broken phononic crystals directly, irrespective of the thickness of the ordinary phononic crystal. Therefore, the one-way transport waveguide mode is robust against almost all kinds of local disorders and arbitrary bends. Moreover, as the acoustic wave is extended inside the ordinary phononic crystal, the width of the waveguide channel can be flexibly tuned by adjusting the thickness of the ordinary phononic crystal.

Due to the flexible control of the width, this type of acoustic one-way waveguide has unique advantages in potential applications. For example, since the backward scattering is totally suppressed inside the one-way waveguide, when we abruptly reduce or increase the thickness of the ordinary phononic crystal, the guided wave will become more localized or more extended along the transverse direction due to energy conservation, thereby realizing the energy squeezing or decompressing. On the other hand, when interfacing with other waveguides of arbitrary widths, an optimized width can greatly reduce the insertion loss.

2. Results and discussions

2.1. The topological one-way waveguide mode
The phononic crystal waveguide investigated in this work is schematically shown in figure 1, which contains three domains marked as A (red), B (gray) and C (blue), respectively. The phononic crystals are composed of identical cylinders with the radius \(r\) arranged into a triangular lattice with the lattice constant \(a\) embedded in water. In domain A (C), the cylinders are spinning in the clockwise (anticlockwise) direction with the spinning circular frequency \(\Omega < 0\) (\(\Omega > 0\)), while in domain B, the cylinders are static. The operation frequency is within the band gaps of domains A and C (they share the same band gap), but on the passband of domain B. Therefore, the guided wave is almost uniformly extended inside domain B but decays into domains A and C. A thin, sound-permeable, and unspinning shell is coated on each cylinder to avoid direct contact between the cylinder and water, so that spinning will not drive water to move. The shell is so thin that the scattering of the shell can be safely ignored.

The bulk bands of domain B and domains A and C are plotted in figures 2(a) and (b), respectively. Throughout the paper, we employed the multiple scattering technique \[49, 50\] to calculate all band dispersions and simulate the propagation of the waveguide mode. For the spinning cylinder, taking both the Doppler effect and Coriolis force into account, the \(n\)th-order Mie coefficient for a time-harmonic \(e^{-i\omega t}\) incident acoustic wave is given by \[31, 51, 52\]

\[
D_n = -\frac{\lambda_n p_0 R_n h_n(k_0 r) - k_0 p J_n(\lambda_n r) J_n'(k_0 r)}{\lambda_n p_0 R_n H_n^{(1)}(k_0 r) - k_0 p J_n(\lambda_n r) H_n^{(1)*}(k_0 r)},
\]

Figure 1. Sketch of the topological one-way waveguide with tunable width. Domain A (red), domain B (gray) and domain C (blue) are phononic crystals composed of clockwise spinning, static and anticlockwise spinning cylinders embedded in water, respectively. The width of domain B can be tuned.
where \( J_n \) and \( H_n^{(1)} \) are the Bessel function and Hankel function of the first kind, \( k_0 \) is the wavenumber in the background medium, \( \rho_0 \) and \( \rho \) denote the mass densities of the background and the cylinder,

\[
\lambda_n = \sqrt{-(M^2 + 4\Omega^2)} / c, \quad \text{with} \quad M = -i(\omega - n\Omega)
\]

being the frequency correction for the nth vector cylindrical wave due to the Doppler rotational effect [32] and \( c \) being the sound speed inside a static cylinder, and the auxiliary function \( R_n \) is expressed as

\[
R_n = \frac{\omega^2(2\Omega^2 - M^2)f_n(\lambda_n r) - 3nM\Omega J_1(\lambda_n r)}{(4\Omega^2 + M^2)(\Omega^2 + M^2)}.
\]

Equation (1) reduces to the expression for static cylinders when \( \Omega = 0 \). When \( \Omega \neq 0 \), it is easily found that \( D_n = D_{-n} \) for \( n \neq 0 \), indicating that TRS of the system is broken.

When the cylinders are static, there is a pair of Dirac points formed at the K and K’ points due to the lattice symmetry. While the cylinders are spinning in the clockwise (counterclockwise) direction, because of the breaking of TRS, the Dirac points are lifted to forms at K and K’ points due to the lattice symmetry. The inset shows the unit cell.

Figure 2 (a) The band structure (dark color lines) of domain B where the cylinders are in the static case. A Dirac point forms at K (K’) point due to the lattice symmetry. The inset shows the unit cell. (b) The band structure (dark color lines) of domain A or C for the spinning frequency \( \Omega = 1.01c_0/a \). The cyan region denotes the band gap which covers the frequency range 0.588 \( \lesssim \omega/a_c \lesssim 0.640 \). The mass density and sound speed of the background (water) are \( \rho_0 = 10^3 \) kg m\(^{-3}\) and \( c_0 = 1489 \) m s\(^{-1}\). The radius and the mass density of the cylinder are \( r = 0.12a \) and \( \rho = 800 \) kg m\(^{-3}\). In the static, the sound speed of the cylinder is \( c = 400 \) m s\(^{-1}\).
velocities (marked by the black circles). The other two surface bands are also for the interface when domain B is above domain A. Because the frequencies are on the passband of domain B, all the cylinders inside the domain B are excited, see the inset of figure 3(b).

It is interesting that the dispersions of the surface band for AB0C and the waveguide band for AB6C are very similar. This similarity is due to the fact that the effective Hamiltonians of the three domains around the K and K’ valleys process the same Fermi velocity \[ v_f \] [49]. According to the k.p method [48, 54], near the K valley, the effective Hamiltonians are expressed as

\[ H_{\text{eff}} = v_f (q_x \sigma_x + q_y \sigma_y) + m \sigma_z, \]  

where \( \sigma_x, \sigma_y, \sigma_z \) are Pauli matrices, \( (q_x, q_y)^T \) is the displacement of the wavevector from K valley in the momentum space, \( v_f \) is the Fermi velocity, \( m > 0 \) for domain A, \( m = 0 \) for domain B and \( m < 0 \) for domain C characterizing the effective mass term. Writing the momentum operators in terms of a spatial derivative \( (i \nabla_{x,y}) \), the wavefunction of the waveguide mode corresponding to the eigenvalue \( \varepsilon = v_f q_x \) is calculated as [17]

\[ \psi(y) = e^{i q_x y} \exp \left( -\frac{1}{v_f} \int_0^y m(y') dy' \right) \left( \begin{array}{c} 1 \\ i \end{array} \right), \]  

where \( y = 0 \) is the bottom of domain B. Therefore, the group velocity of the waveguide mode is \( v_g = \frac{\partial \varepsilon}{\partial q_x} = v_f \), equaling to the Fermi velocity of the effective Hamiltonian irrespective of the thickness of domain B. Moreover, from equation (4) we can see that the wavefunction of the waveguide mode is just regarded as the combination of the chiral surface mode for AB0C and the bulk mode in domain B [45]. Therefore, the distributions of the scattering cross sections of cylinders in domains A and C for both supercells are almost the same, see the insets of figure 3. The waveguide shown by figure 1 is a one-way waveguide since no leftward propagating waveguide mode is supported. To achieve the waveguide that supports only the leftward propagating modes, we just need to swap domain A and domain C, see figure 5. We note that although only two typical cases (AB0C and AB6C) are shown, the conclusions are the same for AB\( _n \)C of an arbitrary and non-negative integer \( n \).

### 2.2. Waveguide of tunable width

Since the guided wave is almost confined to domain B, the thickness of domain B can be regarded as the equivalent width of the one-way waveguide. The thickness of domain B is not fixed but flexibly controlled. Here, we assume that the spinning velocity of each cylinder is controlled by each individual motor. When the cylinder is spinning clockwise, anticlockwise or static, it belongs to domain A, C or B, respectively. For instance, when the layer of cylinders in domain A or C adjacent to domain B become static, the cylinders fall into domain B, enlarging the thickness of domain B. To the contrary, when the layer of cylinders in domain B adjacent to domain A (C) start to spin clockwise (counter-clockwise), they fall into domain A (C), and the thickness of domain B is reduced. Therefore, the waveguide of tunable widths can be achieved by precisely controlling the spinning of cylinders.
every cylinder in the phononic crystals. Our work provides a new type of mechanical control method for active
tunable phononic crystals and metamaterials [55, 56]. Due to the flexibility of the waveguide channel width, we can adjust the widths of the left and right ends individually, and connect the two ends to other waveguides with different cross sections. Therefore, this topological one-way waveguide with tunable width can be used as a connector for waveguides with different cross sections. On the other hand, because of the backscattering immunity, the energy of the guided wave can be squeezed by abruptly reducing the width of the waveguide channel or decompressed by abruptly enlarging the width. Figure 4 shows that the acoustic wave propagates inside a topological one-way waveguide which has varying width along the propagation direction. The line source is located at $x = 10a$. Before $x > 20a$, the layer number of domain B is 4, and it is abruptly changed to 8 at $x = 20a$ and then suddenly reduced to 0 at $x = 38a$. As we can see, the acoustic wave cannot propagate leftward and no backward scattering takes place even though there are large geometry mismatches. In figure 4(b), we showed the normalized energy fluxes of the leftward and rightward propagating waves as functions of the dimensionless frequency by the black and red symbol lines, respectively. It is clearly that within the band gap of domain A or C as shown by the shaded region, the leftward propagation is totally suppressed. We also depicted the pressure fields at the lines $l_1$ and $l_2$ as marked by the white dashed lines in figure 4(a). The pressure fields at the line $l_1$ are vanishing small, in stark contrast to the pressure fields at the line $l_2$, which again confirms that there is no leftward propagating mode supported. The vanishing small pressure fields at the line $l_1$ is attributed to the evanescent waves.

When the acoustic wave propagates from a relatively narrower channel into a relatively broader channel, the field intensities are increased. On the contrary, the field intensities are remarkably enhanced when the acoustic wave goes from a broad channel into a very narrow channel. This flexible control of the field intensities is beneficial to the adjustment of the interaction between acoustic wave and matter. When the domain A and

![Figure 4. (a) Energy squeezing and decompressing. The black disks represent the cylinders and the green lines divide the three domains. The absolute values of pressure fields are shown. The frequency of the acoustic wave is $f_0/c_0 = 0.624$. (b) The leftward and rightward propagating acoustic energy fluxes normalized to the total energy flux as functions of the reduced frequency. The shaded region denotes the band gap of domain A or C. The leftward (rightward) propagating energy flux is probed at the line $l_1$ ($l_2$) as marked by the white dashed line in (a). (c) At $f_0/c_0 = 0.624$, the pressure fields at the two probing lines are depicted. The line source is located at $x = 10a$, $y = \frac{1}{\sqrt{3}}a$.](image-url)
domain C are swapped, as we have discussed previously, the acoustic wave can only propagate leftward instead, see figure 5.

2.3. Robustness against disorders and bends
Due to the topological protection, the one-way transport is robust against local disorders and in principle any kinds of bends. In figure 6(a), we showed the propagation of the acoustic wave in the waveguide with a 90° bend. Domain B of the waveguide has 4 layers. It is seen that the acoustic wave can pass through the bend without any reflection. We also showed that the acoustic wave propagates inside a waveguide with disorders in figure 6(b). Without disorders, domain B of the waveguide has 6 layers. Within the range 20a < x < 30a, we considered that each cylinder in domain B is displaced from its original position by 1.5δ along the x direction and δ\sqrt{3}/2 along the y direction, where -0.5 ≤ δ ≤ 0.5 is a number randomly chosen for each cylinder. And we also removed 6 cylinders during the range 40a < x < 45a. We can see that the acoustic wave is still focused in domain B and the backscattering is totally suppressed.

Figure 5. The waveguide mode when domain A and domain C are swapped. The black disks represent the cylinders and the green lines divide the three domains. The absolute values of pressure fields are shown. The frequency of the acoustic wave is \( f_a/c_0 = 0.624 \).

Figure 6. Robustness against the bend and disorders. The black disks denote the cylinders. (a) The acoustic wave propagates inside the topological one-way waveguide with a 90° bend. The layer number of domain B is 4. (b) Disorders are introduced to domain B. From \( x = 20a \) to \( x = 30a \), the cylinders in domain B are randomly displaced from their original positions. During \( 40a < x < 45a \), 6 cylinders in domain B are removed. The frequency of the acoustic wave is \( f_a/c_0 = 0.624 \).
3. Conclusions

In summary, we have proposed the use of a heterostructure composed by an ordinary phononic crystal sandwiched between two TRS-broken phononic crystals to achieve the acoustic topological one-way waveguide with the width of the waveguide channel can be flexibly tuned by adjusting the thickness of the ordinary phononic crystal. The one-way transport of the guided acoustic wave is protected by the gap-Chern numbers of the two TRS-broken phononic crystals, and thus robust against various kinds of local disorders and arbitrary bends. The tunability of the width enables the topological one-way waveguide to find some unique applications, such as interfacing with waveguides of arbitrary widths and squeezing or decompressing the acoustic energy flexibly.

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Data availability statement

The data generated and/or analysed during the current study are not publicly available for legal/ethical reasons but are available from the corresponding author on reasonable request.

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References

[1] Deymier P A 2013 Acoustic Metamaterials and Phononic Crystals, Springer Series in Solid State Sciences vol 173 (Berlin, Heidelberg: Springer)
[2] Martinez-Sala R, Sancho J, Sanchez J, Gomez V, Linares J and Meseguer F 1995 Sound attenuation by sculpture Nature (London) 378 241
[3] Cheng W, Wang J, Jonas U, Fytas G and Stefanou N 2006 Observation and tuning of hypersonic bandgaps in colloidal crystals Nat. Mater. 5 830
[4] Liu Z, Zhang X, Mao Y, Zhu Y, Yang Z, Chan C and Sheng P 2000 Locally resonant sonic materials Science 289 1734
[5] Fang N, Xi D, Xu J, Ambati M, Srituravanich W, Sun C and Zhang X 2006 Ultrasonic metamaterials with negative modulus Nat. Mater. 5 452
[6] Cheng Y, Xu J and Liu X 2008 One-dimensional structured ultrasonic metamaterials with simultaneously negative dynamic density and modulus Phys. Rev. B 77 045134
[7] Li J, Fok L, Yin X, Bartal G and Zhang X 2009 Experimental demonstration of an acoustic magnifying hyperlens Nat. Mater. 8 931–4
[8] Cummer S and Schurig D 2007 One path to acoustic cloaking New J. Phys. 9 45
[9] Tong S, Ren C and Tao J 2022 Compact topological waveguide for acoustic enhanced directional radiation Appl. Phys. Lett. 120 063504
[10] Temelkuran B, Bayindir M, Ozbay E, Biswas R, Sigalas M M, Tuttle G and Ho K M 2014 Photonic crystal–based resonant antenna with a very high directivity J. Appl. Phys. 87 603
[11] Qiu C, Liu Z, Shi J and Chan C T 2005 Directional acoustic source based on the resonant cavity of two-dimensional phononic crystals Appl. Phys. Lett. 86 224105
[12] Zanjanii M B, Davoyan A R, Mahmoud A M, Engheta N and Lukes J R 2014 One-way phonon isolation in acoustic waveguides Appl. Phys. Lett. 104 081905
[13] Ouyang S, He H, He Z, Deng K and Zhao H 2016 Acoustic one-way mode conversion and transmission by sonic crystal waveguides J. Appl. Phys. 120 104504
[14] Chin C-K, Teo J C Y, Schnyder A P and Ryu S 2016 Classification of topological quantum matter with symmetries Rev. Mod. Phys. 88 035005
[15] Bergholtz E J, Carl Budich J and Kunst F K 2021 Exceptional topology of non-Hermitian systems Rev. Mod. Phys. 93 015005
[16] Lu L, Joannopoulos J D and Soljačić M 2014 Topological photonics Nat. Photon. 8 821
[17] Ozawa H M et al 2019 Topological photonics Rev. Mod. Phys. 91 015006
[18] Rechtsman M C, Zeuner J M, Plotnik Y, Lumer Y, Podolsky D, Dreisow F, Nolte S, Segev M and Szameit A 2013 Photonic Floquet topological insulator Nature 496 196
[19] Ma G, Xiao M and Chan C T 2019 Topological phases in acoustic and mechanical systems Nat. Rev. Phys. 1 281
[20] Lu L, Wang Z, Ye D, Ran L, Fu L, Joannopoulos J D and Soljačić M 2015 Experimental observation of Weyl points Science 349 622
[21] Yang B et al 2018 Ideal Weyl points and helicoid surface states in artificial photonic crystal structures Science 359 1013
[22] Xue H, Yang B, Gao F, Chong Y and Zhang X 2019 Acoustic higher-order topological insulator on a kagome lattice Nat. Mater. 18 108
[23] Harari G et al 2018 Topological insulator laser: theory Science 359 eaar4003
[24] Bandres M et al 2018 Topological insulator laser: experiments Science 359 eaar4005
[25] Dong J-W, Chen X-D, Zhu H, Wang Y and Zhang X 2017 Valley photonic crystals for control spin and topology Nat. Mater. 16 298
[26] He H, Qiu C, Ye L, Cai X, Fan X, Ke M, Zhang F and Liu Z 2018 Topological negative refraction of surface acoustic waves in a Weyl phononic crystal Nature 560 61
[27] Liu, J., Qiu, C., Ye, L., Fan, X., Ke, M., Zhang, F. and Liu, Z. 2017 Observation of topological valley transport of sound in sonic crystals Nat. Phys. 13 369
[28] Yang, Y., Yamagami, Y., Xu, Y., Pitchappa, P., Webber, J., Zhang, B., Fujita, M., Nagatsuma, T. and Singh, R. 2020 Terahertz topological photonics for on-chip communication Nat. Photon. 14 146
[29] Yu, S.-Y., He, C., Wang, Z., Liu, F.-K., Sun, X.-C., Li, Z., Lu, H.-Z., Lu, M.-H., Liu, X.-P. and Chen, Y.-F. 2018 Elastic pseudospin transport for integratable topological phononic circuits Nat. Commun. 9 3072
[30] Nassar, H., Yousefzadeh, B., Fleury, R., Ruzzene, M., Ali, A., Darar, C., Norris, A. N., Huang, G. and Haberman, M. R. 2020 Nonreciprocity in acoustic and elastic materials Nat. Rev. Mater. 5 567
[31] Wang, N., Zhang, R.-Y. and Chan, C. T. 2021 Robust acoustic pulling using chiral surface waves Phys. Rev. Appl. 15 024034
[32] Wang, N., Zhang, R.-Y., Guo, Q., Wang, S., Wang, G. P. and Chan, C. T. 2022 Optical pulling using topologically protected one-way transport surface-arc waves Phys. Rev. B 105 014104
[33] He, C., Ni, X., Ge, H., Sun, X.-C., Chen, Y.-B., Lu, M.-H., Liu, X.-P. and Chen, Y.-F. 2016 Acoustic topological insulator and robust one-way sound transport Nat. Phys. 12 1524
[34] Kitzing, K. V. 1986 The quantized Hall effect Rev. Mod. Phys. 58 519
[35] Hatsugai, Y. 1993 Chern number and edge states in the integer quantum Hall effect Phys. Rev. Lett. 71 3697
[36] Haldane, F. D. M. and Raghu, S. 2008 Possible realization of directional optical waveguides in photonic crystals with broken time-reversal symmetry Phys. Rev. Lett. 100 013904
[37] Raghunath, S. and Haldane, F. D. M. 2008 Analogs of quantum-Hall-effect edge states in photonic crystals Phys. Rev. A 78 033834
[38] Wang, Z., Chong, Y. D., Joannopoulos, J. D. and Soljacic, M. 2008 Reflection-free one-way edge modes in a gyromagnetic photonic crystal Phys. Rev. Lett. 100 013905
[39] Wang, Z., Chong, Y., Joannopoulos, J. D. and Soljacic, M. 2009 Observation of unidirectional backscattering-immune topological electromagnetic states Nature 461 772
[40] Poo, Y., Wu, R. X., Lin, Z. F., Yang, Y. and Chan, C. T. 2011 Experimental realization of self-guiding unidirectional electromagnetic edge states Phys. Rev. Lett. 106 093903
[41] Yang, Z., Gao, F., Shi, X., Lin, X., Gao, Z., Chong, Y. and Zhang, B. 2015 Topological acoustics Phys. Rev. Lett. 114 114301
[42] Khaniyev, A. B., Fleury, R., Mousavi, H. and Alu, A. 2015 Topologically robust sound propagation in an angular momentum-biased graphene-like resonator lattice Nat. Commun. 6 8260
[43] Ni, X., He, C., Sun, X.-C., Liu, X., Lu, M.-H., Feng, L. and Chen, Y.-F. 2015 Topologically protected one-way edge mode in networks of acoustic resonators with circulating air flow New J. Phys. 17 035016
[44] Ding, Y., Peng, Y., Zhu, Y., Fan, X., Yang, J., Liang, B., Zhu, X., Wan, X. and Cheng, J. 2019 Experimental demonstration of acoustic chern insulators Phys. Rev. Lett. 122 014302
[45] Wang, M., Zhang, R.-Y., Zhang, L., Wang, D., Guo, Q., Zhang, Z.-Q. and Chan, C. T. 2021 Topological one-way large-area waveguide states in magnetic photonic crystals Phys. Rev. Lett. 126 067401
[46] Shen, Q., Zheng, X., Zhang, H., You, Y. and Shen, L. 2021 Large-area unidirectional surface magnetoplasmons using uniaxial mu-near-zero material Opt. Lett. 46 5978
[47] Chen, Q., Zhang, L., Chen, F., Yan, Q., Xi, R., Chen, H. and Yang, Y. 2021 Photonic topological valley-locked waveguides ACS Photon 8 1400
[48] Wang, J.-Q. et al. 2022 Extended topological valley-locked surface acoustic waves Nat. Commun. 13 1324
[49] Faulkner, J. S. 1979 Multiple-scattering approach to band theory Phys. Rev. B 19 6186
[50] Chen, S. K., Nicorovici, N. A. and McPhedran, R. C. 1994 Green’s function and lattice sums for electromagnetic scattering by a square array of cylinders Phys. Rev. E 49 4590
[51] Censor, D. and Aboudi, J. 1971 Scattering of sound waves by rotating cylinders and spheres J. Sound Vib. 19 437
[52] Zhao, D., Wang, Y.-T., Fung, K.-H., Zhang, Z.-Q. and Chan, C. T. 2020 Acoustic metamaterials with spinning components Phys. Rev. B 101 054107
[53] Zhang, X., Li, L.-M., Zhang, Z.-Q. and Chan, C. T. 2001 Surface states in two-dimensional metaldielectric photonic crystals studied by a multiple-scattering method Phys. Rev. B 63 125114
[54] Mei, J., Wu, Y. and Chan, C. T. 2012 First-principle study of Dirac and Dirac-like cones in phononic and photonic crystals Phys. Rev. B 86 035141
[55] Chen, S., Fan, Y., Fu, Q., Wu, H., Jin, Y., Zheng, J. and Zhang, F. 2018 A review of tunable acoustic metamaterials Appl. Sci. 8 1480
[56] Wang, Y.-F., Wang, Y.-Z., Wu, B., Chen, W. and Wang, Y.-S. 2020 Tunable and active phononic crystals and metamaterials Appl. Mech. Rev. 72 040801