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Theoretical Study of Surface Plasmons with Phase Singularities Generated by Evanescent Bessel Beams

(Kajian Teori Plasmon Permukaan dengan Singulariti Fasa yang dijana oleh Alur Bessel Evanesen)

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ABSTRACT

We give details about how a surface plasmons with phase singularities can be produced when a Bessel beam light is totally reflected internally at the planar surface of a dielectric on which an infinitesimally thin film has been deposited. The characteristic property of such a light is the exponential decay with distance in a vacuum which can basically provide a two-dimensional surface plasmons with phase singularities with attractive enhancements. Such a phenomenon is governable by altering the incident angle and the order of the Bessel beam.

Keywords: Metallic film; phase singularity; plasmon surface modes

ABSTRAK

Kertas ini membincangkan dengan terperinci bagaimana plasmon permukaan dengan singulariti fasa boleh dihasilkan apabila satu alur cahaya Bessel terpantul sepenuhnya di permukaan satah satu bahan dielektrik yang telah dimendapkan dengan satu lapisan filem yang sangat nipis. Sifat cahaya ini termasuklah pereputan eksponen dengan jarak dalam vakum yang secara asasnya boleh menghasilkan satu permukaan plasmon dua dimensi dengan singulariti fasa dan ini penambahbaikan yang menarik. Fenomena ini boleh dikawal dengan mengubah sudut tuju dan juga tertib alur Bessel tersebut.

Kata kunci: Filem logam; mod permukaan plasmon; singulariti fasa

INTRODUCTION

It is well known that laser light in general can mix with material excitations, producing a diversity of hybrid modes (Maier 2007; Raether 1988; Zayats & Smolyaninov 2003). A familiar type of hybrid excitation is when laser light appears in a vacuum at an interface with an electronically dense medium, leading to surface plasmons. These modes propagate along the planar surface, but their amplitudes decay exponentially with distance perpendicular to the surface and have an upper frequency limit that depends on the electronic density of the material.

In the last decade, plasmons surface modes have received considerable attention in theoretical and experimental studies, as well as in manufacturing in order to avoid or to use them in different applications (Barnes et al. 2003; Pillai et al. 2007; Tan et al. 2009). Surface plasmons are now so well characterized and their dispersion relations so accurately reproducible, that they can be used for calibration purposes (Schaadt et al. 2005). A surface plasmon mode can also be excited on the surface of a film of finite thickness deposited on a planar dielectric substrate (Lembessis et al. 2011). The main property of such a mode is the exponential decay with distance in the vacuum area. Such an excitation persists as a well defined entity, even when the film thickness is very small (Al-Awfi et al. 2011; Kirk et al. 2002) and can experimentally be represented as a two-dimensional sheet. Experience has shown that generating surface plasmons will lead to enhancement by as much as two orders of magnitude (Bennett et al. 2001).

This paper is concerned with the possibility of producing a surface plasmons with phase singularities which do not just have surface-type features (Al-Awfi et al. 2011; Andrews et al. 2010; Lembessis et al. 2011), but are also strongly plasmonic. The light fields so created can be very strong and so will couple strongly to matter localized in the vicinity of the surface.

The paper is organized as follows. Initially, we briefly outline the eclectic field distributions of Bessel beams propagated within a non-dispersive medium. After that, we describe the basic elements of the physical model that is employed to create plasmon surface modes, including a sheet in the form of an infinitesimally thin film deposited on a planar dielectric substrate. The procedure needed for the specification of the light mode bearing the evanescent component is then described, leading to appropriate field distributions in three regions of the structure. Then, we derive an expression for the Rabi frequency and optical dipole potential associated with an evanescent Bessel beam at the outer thin film surface. We explain when this product of the potential profile leads transparently to the presentation of a two-dimensional optical vortex. Lastly, we give a conclusion which contains the most important results.

BESSEL LIGHT FORMULATION

Consider first the electric field of a Bessel beam traveling along a z – *direction* in a medium of constant refractive index n, characterized by the integer ℓ, angular frequency ω, and axial wavevector $k = nk_0$, where $k_0 = \omega/c$ is the wavevector in a vacuum. Such a beam with the plane polarized along the field vector can be written in cylindrical coordinates as (McGloin et al. 2003):

$$\mathbf{E}^{I}_{kl}(\mathbf{r}) = \hat{y}\mathbf{F}_{kl}(r_{\|}, z)\exp^{i(kz-\omega t)}\exp^{il\phi}, \tag{1}$$

where $Fkl(r_{\|}, z)$ is the standard envelope function:

$$\mathbf{F}_{k\ell}(r_{\|}, z) = \eta_{k0}\sqrt{2\pi k_r w_0}\left(\frac{z}{z_{\max}}\right)^{\ell+1/2}\exp\left(-\frac{z^2}{z^2_{\max}}\right) \times\exp\left[i\left(\ell\hat{\phi}-\frac{\ell\pi}{2}-\frac{\pi}{4}\right)\right]J_{\ell}(k_r r_{\|}), \tag{2}$$

where w_0 is the input beam waist, k_r is the radial wavevector and ϕ is the azimuthal coordinate. The factor η_{k0} is the amplitude of a corresponding plane wave of intensity I propagating in the dielectric medium of refractive index n:

$$\eta_{k0} = \sqrt{2I / n^2\varepsilon_0 c}, \tag{3}$$

while $z_{\max}$ is the typical ring spacing. The last term $J_l(k_r r)$ is the Bessel function of the order ℓ. The mathematical formula of the Bessel beam that is given in equation (2) is only valid for the central region of the Bessel beam that has been produced. The main condition for this validity is $2w(z)/w_0 << 1$ where $w(z) = 1/k_r$ is a measure of the width of the central lobe in a J_0 beam or the central dark fringe in a higher-order beam which can be called the beam size at axial coordinate z. This can be written as:

$$w(z) = w_0\sqrt{1+\left(\lambda z/\pi w_0^2\right)^2}. \tag{4}$$

Hence, the standard envelope function can simply be re-written as:

$$\mathbf{F}_{k\ell}(r_{\|}, z) = \eta_{k0}\sqrt{\frac{2\pi w_0}{w(z)}}\left(\frac{z}{z_{\max}}\right)^{\ell+1/2}\exp\left(-\frac{z^2}{z^2_{\max}}\right) \times\exp[i\Phi]J_{\ell}\left(\frac{r_{\|}}{w(z)}\right), \tag{5}$$

where the phase Φ can be written as:

$$\Phi = \ell\hat{\phi}-\frac{\ell\pi}{2}-\frac{\pi}{4} = \ell\tan^{-1}\left(\frac{x}{y}\right)-\frac{2\pi(2\ell-1)}{8}. \tag{6}$$

The plane $z = 0$ corresponds to the minimum beam waist $w(0) = w_0$ and on this plane the phase in equation (6) vanishes.

OBLIQUE INCIDENCE

The basic elements comprising of the physical model are shown schematically in Figure 1. Here a sheet in the form of an infinitesimally thin film is deposited on a planar dielectric substrate. A Bessel beam of a given order ℓ and frequency ω is incident at angle θ greater than the critical angle θ_c. From within the dielectric this is internally reflected at the interface $z = 0$ between the dielectric substrate and the thin film. Consequently a plasmon evanescent mode is produced in the vacuum, compliant with the requirement of the phase-matching condition at the interface, together with the usual electromagnetic boundary condition that the electric field vector which is tangential to the surface, does not suffer any change across the interface.

The relevant electric field vector at frequency ω can be written in term of incident (I), reflected (R) and evanescent (*evan.*) parts as follows:

$$E(\mathbf{r},t) = \{(E_I(\mathbf{k}_{\|},\mathbf{r},t) + E_R(\mathbf{k}_{\|},\mathbf{r},t))\Theta(-z) + E_{evan.}(\mathbf{k}_{\|},\mathbf{r},t)\Theta(-z)\}a + H.C., \tag{7}$$

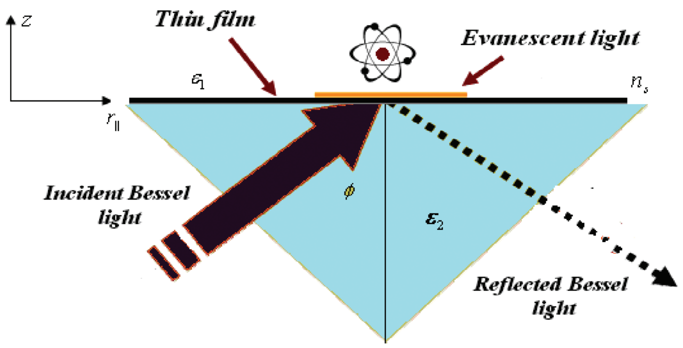


FIGURE 1. Schematic total internal reflection of a single Bessel beam at a planar dielectric interface with a thin film, producing an evanescent mode

where Θ is the unit step function and the fields are given by:

$$\mathbf{E}_I(\mathbf{k}_\parallel,\mathbf{r},t) = N_I\left(1,0,\frac{k_\parallel}{\mathrm{k}_{z2}}\right)\exp(ik_{z2}z)$$
$$\exp i(\mathbf{k}_\parallel.\mathbf{r}_\parallel - \omega t), \tag{8}$$

$$\mathbf{E}_R(\mathbf{k}_\parallel,\mathbf{r},t) = \mathrm{N}_R\left(1,0,\frac{k_\parallel}{\mathrm{k}_{z2}}\right)\exp(-ik_{z2}z)$$
$$\exp i(\mathbf{k}_\parallel.\mathbf{r}_\parallel - \omega t), \tag{9}$$

$$\mathbf{E}_{van.}(\mathbf{k}_\parallel,\mathbf{r},t) = \mathrm{N}_{evan.}\left(1,0,\frac{ik_\parallel}{\mathrm{k}_{z1}}\right)\exp(-k_{z1}z)$$
$$\exp i(\mathbf{k}_\parallel.\mathbf{r}_\parallel - \omega t) \tag{10}$$

where carets denote unit vectors and the coefficient N_I represents the in-plane incident mode form of the Bessel beam. The quantities N_R and $\mathrm{N}_{evan.}$ are to be determined in terms of N_I by the application of boundary conditions. The incident field amplitude N_I has the standard form:

$$\mathrm{N}_I = \frac{\sqrt{2\pi}\eta_{k0}}{\left(1+\left(\lambda z/\pi w_0^2\right)^2\right)^{1/4}}\left(\frac{z}{z_{\max}}\right)^{\ell+1/2}\exp\left(-\frac{z^2}{z_{\max}^2}\right)$$
$$\times \exp i\left[\ell\tan^{-1}\left(\frac{x}{y}\right)\right]-\left(\frac{2\pi\left(2\ell-1\right)}{8}\right)$$
$$\times J_\ell\left(\frac{r_\parallel}{w_0\sqrt{1+\left(\lambda z/\pi w_0^2\right)^2}}\right), \tag{11}$$

and $\mathbf{k}_\parallel$ is the wavevector parallel to the surface. Its magnitude $\mathbf{k}_\parallel$ is given by $c^2k_\parallel^2 = \omega^2\varepsilon_2\sin^2\theta$. The three quantities between the brackets in equations (8) to (10) stand for the vector components parallel to $\mathbf{k}_\parallel$, perpendicular to it on the surface plane and along the *z – direction*, respectively. k_{z1} and k_{z2} (both real) are defined by:

$$k_{z1}^2 = k_\parallel^2 - \varepsilon_1\omega^2/c^2 > 0. \tag{12}$$

$$k_{z2}^2 = \varepsilon_2\omega^2/c^2 - k_\parallel^2 > 0. \tag{13}$$

The notation is such that the parameters associated with the substrate are labeled by the subscript 2, while for the outer region (the vacuum) the label is 1. The dielectric function ε_2 is assumed, in general, to be frequency-dependent (i.e. $\varepsilon^2 = n^2$) while we assume that $\varepsilon_1 = 1$, as appropriate for a vacuum. The position vector is written as $\mathbf{r} = (\mathbf{r}_\parallel, z)$ in terms of an in-plane position vector $\mathbf{r}_\parallel$ and a *z–coordinate* relative to the thin film. The role of the thin film is primarily to provide a two-dimensional charge density n_s and so, an electric conductivity $in_se^2/m^*(\omega+i\gamma)$, where m^* and e are the electronic effective mass and charge and $\gamma \ll \omega$ accounts for film plasma loss effects.

However, the thin film only enters the formalism via the electromagnetic boundary conditions involving the tangential components of the magnetic fields corresponding to equations (8) to (10). This can be calculated using Maxwell's equation $\mathbf{H} = -(i\varepsilon_0c^2/\omega)\nabla \times \mathrm{E}$. The application of the first boundary condition, namely the continuity of the tangential component of the electric field vector at $z = 0$, yields:

$$\mathrm{N}_I + \mathrm{N}_R + \mathrm{N}_{evan.} \tag{14}$$

The second electromagnetic boundary condition is that the tangential component of the magnetic field vector experiences a discontinuity at $z = 0$, arising from the surface current induced by the in-plane component of the electric field at the thin film. We have:

$$H_\parallel(0_-) - H_\parallel(0_+) = \frac{in_se^2}{m^*\left(\omega+i\gamma\right)}E_\parallel\left(0\right), \tag{15}$$

where $0_\pm$ are the limit as $\zeta \to 0$ of $(0 \pm \zeta)$. The application of the boundary condition leads to a second relation connecting the field amplitudes:

$$\frac{\varepsilon_2}{k_{z2}}\left(\mathrm{N}_I - \mathrm{N}_R\right)+\frac{i\varepsilon_1}{k_{z1}}\mathrm{N}_{even.} = \frac{in_se^2}{\varepsilon_0m^*\omega\left(\omega+i\gamma\right)}. \tag{16}$$

The elimination of N_R between equations (14) and (16) yields straightforwardly:

$$\mathrm{N}_{evan.} = 2\mathrm{N}_I\left[\frac{ik_{z2}}{\varepsilon_2}\left(\frac{n_se^2}{\varepsilon_0m^*\omega\left(\omega+i\gamma\right)}-\frac{\varepsilon_1}{k_{z1}}\right)+1\right]^{-1}. \tag{17}$$

The next step is to fix the value for the amplitude N_I. This is the amplitude of the incident field in the unbounded bulk of material 2 and the value of N_I. It is such that the Hamiltonian H_I reduces to the canonical form:

$$H_\mathrm{I} = \frac{\varepsilon_0}{2}\int d\mathbf{r}\left\{\varepsilon_0E_I^2+\frac{1}{\mu_0}B_I^2\right\} = \frac{1}{2}\hbar\omega\left(aa^\dagger + a^\dagger a\right). \tag{18}$$

The incident of the Bessel beam is such that its axis is inclined at an angle of incidence θ relative to the *z – axis*. This means that the form of the coefficient N_I represents the amplitude distribution of the incident Bessel beam, and can be obtained from the standard functional form of the Bessel mode profile by the transformation equations (Al-Awfi et al. 2011):

$$x \to -x\cos\theta,\ y \to y,\ z \to -x\sin\theta. \tag{19}$$

Thus we find:

$$N_l = \frac{\sqrt{2\pi}\eta_{k0}}{\left(1+\left(\lambda z/\pi w_0^2\right)^2\right)^{1/4}}\left(\frac{-x\sin\theta}{z_{\max}}\right)^{\ell+1/2}\exp\left(-\frac{\left(-x\sin\theta\right)^2}{z_{\max}^2}\right)$$
$$\times\exp i\left[\ell\tan^{-1}\left(\frac{x\cos\theta}{y}\right)-\left(\frac{2\pi\left(2\ell-1\right)}{8}\right)\right]$$
$$\times J_\ell\left(\frac{\left(x\cos\theta\right)^2+y^2)^{1/2}}{w_0\sqrt{1+\left(\lambda z/\pi w_0^2\right)^2}}\right). \quad (20)$$

The explicit form of the evanescent electric field that displays the mode characteristics then emerges as follows:

$$\mathbf{E}_{evan.}\left(k_\parallel,\mathbf{r}\right)=\frac{2\mathrm{N}_l\left(1+k_\parallel^2/k_{z1}^2\right)\exp\left\{-zk_0\sqrt{n^2\sin^2\phi-1}-ik_0nx\sin\theta\right\}}{\left[\frac{ik_{z2}}{n^2}\left(\frac{n_se^2}{\varepsilon_0m^*\omega\left(\omega+i\gamma\right)}-\frac{\varepsilon_1}{k_{z1}}\right)+1\right]\exp\left(k_{z1}z\right)}. \quad (21)$$

Surface Plasmons with Phase Singularities

In the presence of any inhomogeneous fields such as an evanescent plasmon mode due to the Bessel beams, the atom becomes subject to light-induced forces. These forces include the dissipative force $F_{dissp.}$ that has been exploited in the heating and cooling of atomic motion and the dipole force $F_{dip.}$ (or gradient of dipole optical potential $F_{dip.} = -\vec{\nabla}U$) that is used to generate the trapping potential. The potential U acting on an adsorbed atom of transition ω_0 moving in the vacuum region at velocity **v** due to the evanescent light of frequency ω is given by the well known form (Babiker & Al-Awfi 1999):

$$U(\mathbf{R},\mathbf{v}) = \frac{\hbar\delta}{2}\ln\left[1+\frac{2\Omega^2}{\delta^2+\Gamma^2}\right], \quad (22)$$

where Γ is the decay emission rate of the atom and δ the dynamic detuning given by:

$$\delta = \delta_0 - \mathbf{v}.\vec{\nabla}\theta. \quad (23)$$

Here $\delta_0 = (\omega - \omega_0)$ is static detuning and Ω is the Rabi frequency which is characterized by the interaction of an atom of electric dipole moment $\boldsymbol{\mu}$ approaching the thin film surface from the vacuum region ($z > 0$) and so interacting with the evanescent light. The Rabi frequency is defined as (Al-Awfi et al. 2010):

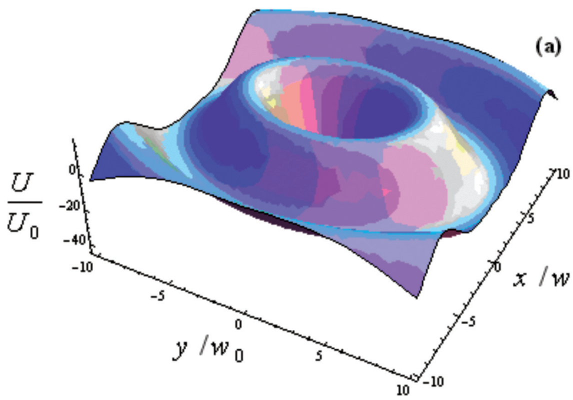
$$\Omega_{knm}(z \geq 0) = |\mathbf{u}.\mathbf{E}_{even.}(\mathbf{k}_\parallel,\mathbf{r})/\hbar|. \quad (24)$$

Using equation (24), we have the Rabi frequency as:

$$\Omega_{k\ell}\left(z\geq 0\right)=\frac{\sqrt{8\pi}\mu\eta_{k0}\left(1+k_\parallel^2/k_{z1}^2\right)^{1/2}}{h\left(1+\left(\lambda z/\pi w_0^2\right)^2\right)^{1/4}\left[ik_{z2}n^{-2}\rho+1\right]}\left(\frac{-x\sin\theta}{z_{\max}}\right)^{\ell+1/2}$$
$$\times\exp\left(-\frac{\left(-x\sin\theta\right)^2}{z_{\max}^2}\right)\times J_\ell\left(\frac{\left(x\cos\theta\right)^2+y^2)^{1/2}}{w_0\sqrt{1+\left(\lambda z/\pi w_0^2\right)^2}}\right)\times\exp\left(-k_{z1}z\right), \quad (25)$$

where $\rho = [n_se^2[\varepsilon_0m^*\omega+i\gamma)]^{-1} - \varepsilon_1(k_{z1})^{-1}]$.

In view of the fact that adsorption process is due to a potential that depends on the coordinate z vertical to the thin film surface, we will only be concerned with the motion of the atom in the parallel adsorption plane close to the thin film surface. The potential distribution (in unit $U_0 = (1/2)\mathrm{h}\Gamma$) corresponds to the evanescent light produced by an incident Bessel beam for which $\ell = 0$ is shown in Figure 2. This potential distribution is plotted at a fixed value of $z > 0$ close to the thin film surface. The parameters are such that the decay emission rate Γ is taken to be the free space value; $\Gamma_0 = 6.1 \times 10^7\mathrm{s}^{-1}$ for *Na* atom with transition wavelength $\lambda_0 = 589$ nm. This is, in fact, a very good approximation of a dynamics region which is adequately far from the thin film surface. The static detuning is taken to be $\delta_0 = 5.0 \times 10^2\Gamma_0$ and the intensity of the incident Bessel light is assumed to be $I = 2.0 \times 10^5$ Wm^{-2}. Finally the areal density n_s of the thin film is fixed at $n_s = 215n_s^{silver}$ (where $n_s^{silver} = 5.573 \times 1018$ m^{-2}). It can be deduced from this figure that the product potential possesses distinct maxima and minima values that could be



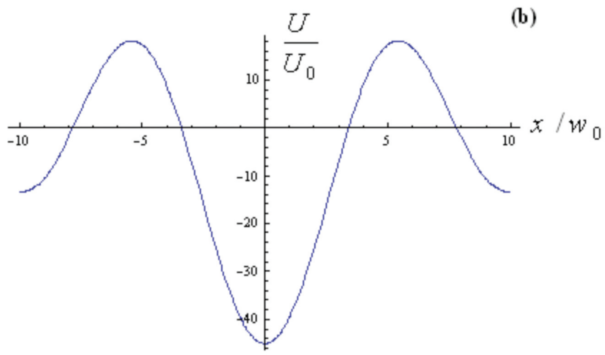


FIGURE 2. The optical dipole potential distribution due to the plasmonic surface mode with $\ell = 0$. The angle of incidence is taken as 42.00. (a) Three-dimension profile and (b) two-dimension profile

used to trap (reject) atoms that have transition frequencies suitably detuned from the frequency ω of the Bessel beam. Additionally, it can be seen that the profile of the potential distribution is actually not as an ideal back bowl-shaped quantum well, but has a slightly elliptical profile, because the light strikes the surface at the angle of incidence θ.

The potential distributions corresponding to the evanescent light produced by an incident Bessel beam of order $\ell = 1$ and $\ell = 2$ are respectively shown in Figures 3 and 4, with the same parameters as in the earlier figure. We can easily see that the potential with $\ell = 0$ can be exactly worked as an atomic mirror system, hence will push any moving atom away from the thin film surface (Cooke & Hill 1982). However, the optical potential with $\ell > 0$ presents a more complex potential distribution which will give rise to additional effect in terms of atom dynamics which could not have been realized with $\ell = 0$. This is due to the orbital angular momentum property associated with the azimuthal dependence of the Bessel field structure. The interaction of the atom with an evanescent Bessel mode of order $\ell > 0$ gives rise to a number of rotational effects which will make the atomic motion drastically different from that associated with $\ell = 0$. These rotational effects produces the phenomenon referred to as a surface plasmons with phase singularities. Clearly the potential distribution for any value of the parameter ℓ can be obtained in a similar manner. In addition, these depths can be controlled by changing the intensity and incident angle of the Bessel beam.

CONCLUSION

In the above analysis, we have outlined a rigorous theory of the surface plasmons with phase singularities based on one of simplest methods for exciting surface plasmon modes. This is especially the case when an infinitesimally thin film is deposited on a planar dielectric substrate to produce surface plasmon modes by the total internal reflection process of a Bessel light. We have shown that this configuration can be used as a significant element for the realization of a two-dimensional surface plasmons with phase singularities for cold and trapped atoms. This is an aspect which has attracted considerable attention in the last few years. The process of the surface plasmons with phase singularities is well recognized and can be readily carried out and duplicated in the laboratory. The set-up can be extended to include a system of incident counter-

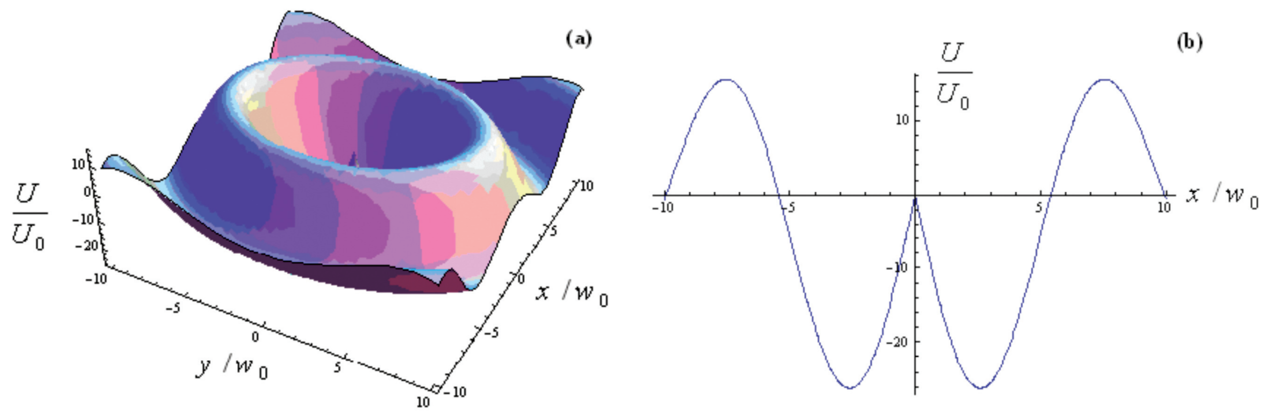


FIGURE 3. The optical dipole potential distribution due to the plasmonic surface mode with $\ell = 1$. The angle of incidence is taken as 42.00. (a) Three-dimension profile and (b) two-dimension profile

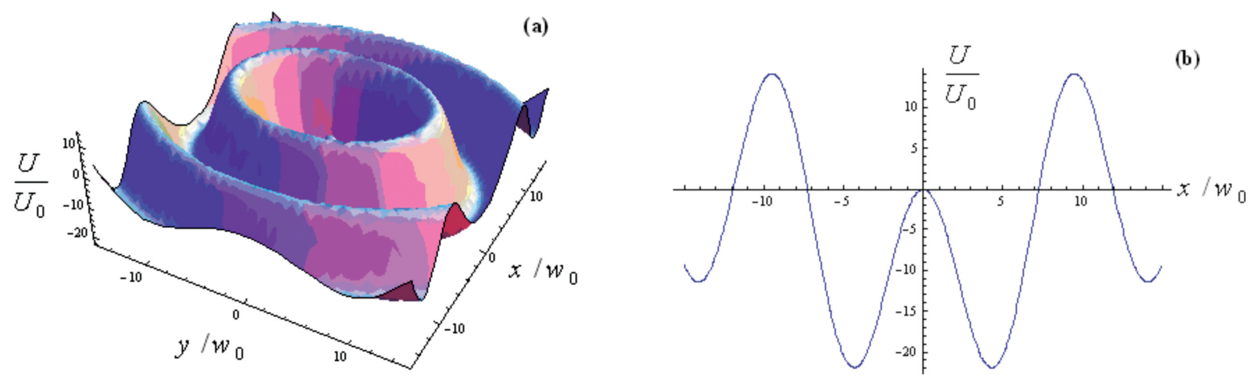


FIGURE 4. The optical dipole potential distribution due to the plasmonic surface mode with $\ell = 2$. The angle of incidence is taken as 42.00. (a) Three-dimension profile and (b) two-dimension profile

propagating or co-propagating Bessel beams which will have a diversity of evanescent Bessel intensity patterns, whose interference effects are expected to create novel sets of potential profiles.

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